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Cosmological Constraints on Perturbative Supersymmetry Breaking

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Abstract

We discuss the cosmology of string models with perturbative supersymmetry breaking at a scale of $\mathcal{O}$(TeV). Such models exhibit Kaluza-Klein like spectra and contain unstable massive gravitinos/gravitons. We find that considerations of primordial nucleosynthesis constrain the maximum temperature following inflation to be not much larger than the supersymmetry breaking scale. This imposes conflicting requirements on the scalar field driving inflation, making it rather difficult to construct a consistent cosmological history for such models.

There has recently been much interest in the possibility of realistic string theories with spontaneous *perturbative* supersymmetry breaking [1, 2]. These theories are very predictive in that they yield, in addition to the supersymmetric standard model, an entirely new phenomenon with a striking experimental signature [1]-[5]. Specifically they contain a repeating spectrum of Kaluza-Klein (KK) modes all the way up to the Planck scale, whose spacing ($\epsilon \approx 1/2R$, where R is the radius of compactification) is comparable to the supersymmetry breaking scale which is of $\mathcal{O}$(TeV). These modes can be excited at forthcoming accelerators such as the LHC [3], hence such models should be of immediate interest to experimentalists. However there are open questions concerning the cosmological viability of these models which need to be addressed first. In this paper we will investigate whether the Kaluza-Klein spectrum is consistent with cosmological constraints on massive unstable relic particles [6]. (Note that the more commonly discussed models with dynamical supersymmetry breaking in a 'hidden' sector are also constrained by similar cosmological considerations [7, 8].)

The only previous relevant discussion on the cosmology of Kaluza-Klein theories concentrated on the prospect that they may include absolutely stable massive particles referred to as 'pyrgons' [9]. Such particles reside on the first rung of the ladder of Kaluza-Klein states, and are unable to decay because they carry a charge which is not exhibited by any of the massless particles. We shall not consider such models since, as we shall see, 'pyrgons' do not exist in string theories with spontaneously broken supersymmetry. Although motivated primarily by string theory, our discussion will apply to all Kaluza-Klein theories in which the couplings are approximately independent of winding number.

For such a theory, the thermal history of the universe is radically altered in the following way. We assume, as is usual, that there was an inflationary DeSitter phase, followed by reheating to a temperature $T_{\rm R}$ [10]. In conventional supergravity, reheating results in the production of gravitinos with number density proportional to $T_{\rm R}$; the subsequent decays of the gravitinos can adversely affect primordial nucleosynthesis and requiring that they not do so results in an upper bound on the reheat temperature of $\sim 10^5$ TeV [11]-[13]. After reheating, the entropy, which we shall assume is subsequently conserved, is evenly spread out amongst the strongly (as opposed to gravitationally) interacting KK modes and the massless matter multiplets. At a temperature much higher than the KK level-spacing ($T \gg \epsilon$), nearly all the entropy is in the KK modes and almost none in the matter multiplets. Until the temperature drops below the first KK level, the evolution of the universe is therefore governed by the KK modes, whose contribution to the entropy is continually decreasing as the temperature drops. During this period there is production of massive gravitons and gravitinos which can only decay to the massless (twisted) particles since their decays to untwisted KK modes is kinematically suppressed. Under these circumstances one might suspect that there is very severe bound on $T_{\rm R}$ and this indeed turns out to be the case.

Let us first present the 'conventional' picture. Gravitinos are generated at high temperatures by two-body scatterings and the equation governing their number density is

$$\dot{n}_{3/2} + 3\,H\,n_{3/2} = \langle\sigma v\rangle n_{\rm rad}^2 - \frac{n_{3/2}}{\tau_{3/2}}\,, \tag{1}$$

where H is the Hubble expansion rate and $\langle\sigma v\rangle$ is the thermally averaged cross-section for gravitino production in the radiation bath of number density $n_{\rm rad}$. The gravitino lifetime

$\tau_{3/2}$ is given at rest by [14]

$$\tau_{3/2} \sim M_{\rm Pl}^2/m_{3/2}^3 \simeq 10^5 \ {\rm sec} \left(\frac{m_{3/2}}{\rm TeV}\right)^{-3} , \tag{2}$$

where $M_{\rm Pl} \equiv G_{\rm N}^{-1/2} \simeq 1.22 \times 10^{16}$ TeV. Given the effective $N = 1$ supergravity couplings,

$$\begin{aligned} \delta L &= \frac{\sqrt{2\pi}}{2M_{\rm Pl}} \bar{\lambda}_a \gamma^\rho \sigma^{\mu\nu} \psi_\rho F^a_{\mu\nu} + {\rm h.c.} \\ &\quad + \frac{\sqrt{2\pi}}{M_{\rm Pl}} \bar{\psi}_\rho \gamma^\mu \partial_\mu z^i \gamma^\rho \psi_i + {\rm h.c.} \ , \end{aligned} \tag{3}$$

it can be shown that $\langle \sigma v \rangle \sim (8\pi/M_{\rm Pl}^2)$ at temperatures $T \ll M_{\rm Pl}$ [12]. The radiation density is given by

$$n_{\rm rad} = g(T) \ \frac{\zeta(3) T^3}{\pi^2} , \tag{4}$$

where $g(T)$ counts the relativistic degrees of freedom contributing to the total number density and is constant above temperatures of $\mathcal{O}$(TeV) for the minimal supersymmetric standard model (MSSM), with

$$g(T_{\rm R}) = \hat{g} = 427/2 \ . \tag{5}$$

Assuming the canonical radiation-dominated evolution at this time, we can solve eq.(1) to obtain

$$Y_{3/2}(T) \equiv \frac{n_{3/2}}{n_{\rm rad}} = \frac{g_s(T)}{g_s(T_{\rm R})} \ \frac{n_{\rm rad}(T_{\rm R}) \langle \sigma v \rangle}{H(T_{\rm R})} \ \exp(-t/\tau_{3/2}) \ , \tag{6}$$

where time is related to temperature as

$$t = 2.42 \times 10^{-12} \ [g_s(T)]^{-1/2} \ {\rm sec} \left(\frac{T}{\rm TeV}\right)^{-2} , \tag{7}$$

and the factor $g_s(T)/g_s(T_{\rm R})$ takes into account the decrease in the number of relativistic degrees of freedom, given constant total entropy

$$s R^3 = g_s(T) \frac{\pi^2 T^3}{30} R^3 = {\rm constant} \ . \tag{8}$$

Note that we have taken $g_s(T)$ to also be the number of degrees of freedom determining the total energy density, as is appropriate at temperatures above a few MeV (when the neutrinos decouple). For the MSSM, one has

$$g_s(T_{\rm R}) = \hat{g}_s = 915/4 \ , \tag{9}$$

at high temperatures when all particles are relativistic.

Now let us consider the cosmological evolution when KK modes are present. Above the supersymmetry breaking scale the number of relativistic degrees of freedom is now no longer constant. The KK modes are labelled by quantum numbers of internal momenta/charges which are of the form

$$P_{\substack{\rm L \\ \rm R}} = \frac{n}{R} \pm \frac{mR}{2} , \tag{10}$$

where R represents some internal radius of compactification. The winding modes ($m \neq 0$) have masses of $\mathcal{O}(M_{\rm Pl})$ and need not be considered further, while the particles in the nth KK mode have masses $m_n \sim n\epsilon$. Roughly speaking, whenever the temperature is raised by ϵ, two new levels of (gauge interacting) KK excitations becomes relativistic, so that the number of degrees of freedom increases *linearly* with temperature. We can allow for this by writing

$$g(T) = (\hat{g} + \frac{T}{\epsilon} g_{\rm K}) \ , \qquad g_s(T) = (\hat{g}_s + \frac{T}{\epsilon} g_{s\rm K}) \ . \tag{11}$$

The constants $g_{\rm K}$, $g_{s\rm K}$ are determined by evaluating the number density and entropy density, respectively, of the plasma. For example consider the number density of KK modes in equilibrium:

$$n_{\rm eq}(T) = \sum_n \frac{g_i}{(2\pi)^3} \int {\rm d}^3 p \ \frac{1}{\exp(\sqrt{p^2 + n^2\epsilon^2/T}) \pm 1} \ , \tag{12}$$

where g_i is the total number of interacting degrees of freedom in any KK level. Using various redefinitions, this becomes

$$n_{\rm eq}(T) \sim \frac{g_i T^4}{\pi^2 \epsilon} \int_0^\infty {\rm d}x \int_{\epsilon/T}^\infty {\rm d}y \ \frac{x^2}{\exp\sqrt{x^2+y^2} \pm 1} \ , \tag{13}$$

where we have approximated the sum at small ϵ/T by an integral, and included a factor of two for positive and negative values of internal momentum (defined above as n). The integral becomes temperature independent for $T \gg \epsilon$, and we find that it deviates by less than 5% from the T^4 behaviour for high values of T/ϵ ($\gg 2$). The limiting values (when $\epsilon/T \to 0$) may be determined analytically to be

$$g_{\rm K} = g_i(\chi_{\rm F} + \chi_{\rm B}) \ , \tag{14}$$

where

$$\chi_{\rm F} = 7\pi^5/480\,\zeta(3) = 3.71 \ , \qquad \chi_{\rm B} = \pi^5/60\,\zeta(3) = 4.24 \ . \tag{15}$$

In a similar fashion, the contribution of the KK modes to the total entropy may be used to determine $g_{s\rm K}$. Using $s \equiv (\rho + p)/T$, we find that

$$s_{\rm eq}(T) \approx \frac{g_i T^4}{\pi^2 \epsilon} \int_0^\infty {\rm d}x \int_{\epsilon/T}^\infty {\rm d}y \ \frac{x^2}{\exp\sqrt{x^2+y^2} \pm 1} \ \frac{(4x^2/3 + 3y^2)}{\sqrt{x^2+y^2}} \tag{16}$$

which gives limiting values of

$$g_{s\rm K} = g_i(\chi_{s\rm F} + \chi_{s\rm B}) \ , \tag{17}$$

where

$$\chi_{\rm F} = 10125\,\zeta(5)/64\pi^3 = 5.29 \ , \qquad \chi_{\rm B} = 675\,\zeta(5)/4\pi^3 = 5.64 \ . \tag{18}$$

In the spontaneously broken string theories, each KK level comes in $N = 4$ multiplets, so that KK gauge bosons contribute 8 bosonic and 8 fermionic degrees of freedom in the vector and fermionic representations of SO(8) respectively. In the minimal case in which the KK excitations are in SU(3) $\otimes$ SU(3)$_{\rm c}$ multiplets [4], this gives

$$g_{\rm K} = 128(\chi_{\rm F} + \chi_{\rm B}) = 1018, \qquad g_{s\rm K} = 128(\chi_{sF} + \chi_{sB}) = 1400 \ . \tag{19}$$

Of course there are additional contributions from higgs multiplets which are also expected to have KK excitations, so we shall consider these values to be a lower limit. To find the time-temperature relation, we assume that after inflation the metric is of the usual FRW form, with all the relativistic degrees of freedom in chemical equilibrium and therefore at the same temperature. Entropy conservation then gives

$$sR^3 = \frac{g_{sK(T)}}{\epsilon}\frac{\pi^2 T^4}{30}R^3 = \text{constant}, \tag{20}$$

and in particular

$$H = -\frac{4}{3}\frac{\dot{T}}{T} \, . \tag{21}$$

Thus the Hubble parameter is

$$H_{(T)} = 1.66\sqrt{\frac{g_{sK}}{\epsilon}}\frac{T^{5/2}}{M_{\text{Pl}}} \, . \tag{22}$$

Differentiating with respect to time and substituting eq.(21), we find

$$t_{(T)} = \frac{8}{15 H_{(T)}} \, . \tag{23}$$

For the minimal value of g_{sK} above, this becomes

$$t = 6.9 \times 10^{-14} \ \text{sec} \left(\frac{\epsilon}{\text{TeV}}\right)^{1/2}\left(\frac{T}{\text{TeV}}\right)^{-5/2} , \tag{24}$$

at temperatures $T \gg \epsilon$.

In order to ascertain the abundance of massive gravitons/gravitinos, we need to identify the processes which can contribute to their manufacture and decay. Vertices between KK modes (which come from untwisted sectors of the string theory) must satisfy the condition that their internal momenta/charges are conserved, and are simply proportional to the string coupling constant. In addition vertices can exist between untwisted modes and twisted (massless) matter multiplets, with couplings $g_n \propto g\delta^{-m_n^2/M_{\text{Pl}}^2} \approx g$, where δ is some constant depending on the type of compactification [15]. For masses much less than M_{Pl} these are clearly unsuppressed, so that we can write down effective terms for the coupling of the KK modes to each other, and to the massless multiplets. The creation of KK gravitons and gravitinos goes via effective four particle interactions. For example the cross section for

$$A_k^a + A_l^b \rightarrow \lambda_m^c + \psi_n^\mu \tag{25}$$

goes as

$$\sigma \sim \frac{g^2}{64\pi} f_{abc} f^{abc} \frac{8\pi}{M_{\text{Pl}}^2}\delta(n+m-k-l) \, , \tag{26}$$

where the integers k, l, m, n label KK modes, and we have omitted numerical factors of $\mathcal{O}(1)$ coming from the trace over solutions to the Rarita-Schwinger equation and phase-space integrations. We therefore write the total cross section for n-gravitino production from k plus l interactions as

$$\sigma_{klmn} = \hat{\sigma}\ \delta(n+m-k-l) \, . \tag{27}$$

where $\hat{\sigma}$ is a factor of $\mathcal{O}(8\pi/M_{\rm Pl}^2)$, which is dependent on the details of the model, but independent of the KK-modes. In addition there is a contribution coming from the massless sector which we shall neglect since there is only one such sector.

The massive n-gravitinos/gravitons may decay into either two massless twisted states, or to untwisted l plus $n-l$-states. In the first case the decay rate is found to be,

$$\Gamma_{\rm twisted} \approx \left(\frac{m_n^3}{M_{\rm Pl}^2}\right) . \tag{28}$$

The decay to untwisted states is kinematically suppressed however. Consider a positive-n state decaying to an l-state plus an $n-l$-state. If l is negative then the masses of the products is $|n|\epsilon + 2|l|\epsilon$ which is larger than the mass, $|n|\epsilon$, of the decaying particle. If l is also positive, then sum of the product masses is equal to that of the decaying particle. The decay rate is proportional to the momentum of the decay products in the centre of mass frame:

$$\Gamma_{\rm untwisted} \propto |p| = \sqrt{1 + \frac{[l^2-(n-l)^2]^2}{n^4} - 2\frac{[l^2+(n-l)^2]}{n^2}}\,\frac{\epsilon|n|}{2} = 0 . \tag{29}$$

As discussed in ref.[5], there may be a mass splitting of $\mathcal{O}(\epsilon)$ in any KK level, so that the decay rate to untwisted modes will generally be suppressed by a factor $1/|n|$. When the sum over l is taken this decay mode may be of the same order as the decay to twisted states. We shall therefore take the lifetime of n-gravitons/gravitinos to be

$$\tau_n \sim 9.8 \times 10^4 \;\sec \left(\frac{\epsilon}{\rm TeV}\right)^{-3} |n|^{-3} , \tag{30}$$

corresponding to the lifetime for decay of standard gravitinos into photons and photinos [14]. Inclusion of all the strongly interacting, twisted, final states would speed up the decays by a small factor. For example, if the twisted products consisted of all the matter and Higgs particles in the MSSM, then the above would be reduced by a factor 12/49.

With these estimates we are ready to tackle the evolution of the n-graviton/gravitino number density n_n. This is governed by the equation [16]

$$\begin{aligned} \dot{n}_n + 3Hn_n + \frac{n_n}{\tau_n} &= \sum_{mkl}\sum_{spins} \int\!\!\int\!\!\int\!\!\int \frac{d^3q_k}{(2\pi)^3}\frac{d^3q_l}{(2\pi)^3}\frac{d^3q_m}{(2\pi)^3}\frac{d^3q_n}{(2\pi)^3}(2\pi)^4\delta^4(q_k+q_l-q_m-q_n) \\ & f_k f_l (1-f_m)\,|M_{kl\to mn}|^2 \end{aligned} \tag{31}$$

where the f_i are the occupation numbers and we have assumed that the massive particles are non-relativistic when they decay. This equation describes the creation of a gravitationally interacting n-state plus a gauge interacting m-state, from gauge interacting k plus l-states. We have taken the occupation number of the n-state to be negligible (since it is never in equilibrium), so that the reverse process does not occur. We may reasonably adopt equilibrium distributions for the three remaining (k,l,m) states, omit the Pauli blocking factor for the m-state, and rewrite eq.(31) as

$$\dot{n}_n + 3Hn_n + \frac{n_n}{\tau_n} = \sum_{mkl} \sigma_{klmn}\, n_{{\rm eq}(k,T)}\, n_{{\rm eq}(l,T)}, \tag{32}$$

where, since we are concerned with the longest lived states, we have taken $v \approx c$. Although approximate, this expression has the correct temperature dependence and we shall bury our approximations in the parameter $\hat{\sigma}$ which is still of $\mathcal{O}(8\pi/M_{\rm Pl}^2)$. Defining

$$n_n = \hat{n}_n \ \exp(-t/\tau_n), \tag{33}$$

and using eq.(21), we find

$$T^4 \frac{\mathrm{d}}{\mathrm{d}t} \left(\frac{\hat{n}_n}{T^4}\right) = \sum_{mkl} \sigma_{klmn} \ n_{\mathrm{eq}(k,T)} \ n_{\mathrm{eq}(l,T)} \ \exp(-t/\tau_n) \ . \tag{34}$$

The number density of the individual k and l states follow some equilibrium curve which obeys, e.g.

$$n_{\mathrm{rad}} \approx n_{\mathrm{eq}(T)} = \sum_{k=-\infty}^{+\infty} n_{\mathrm{eq}(k,T)} \ . \tag{35}$$

Since the epoch of nucleosynthesis is much later than the time at which these particles are created, we take the exponential factor to be unity for the purposes of calculating the initial abundance. Performing the summations in eq.(34) and integrating from $T_{\rm R}$ to $T = \epsilon$ gives

$$Y_{n(T)} = \frac{8}{9} \ \frac{g_{s(T)}\, g_{\rm K}}{\hat{g}_s\, \hat{g}} \ \frac{\hat{\sigma}\, n_{\mathrm{rad}}(T_{\rm R})}{H(T_{\rm R})} \ \exp(-t/\tau_n), \tag{36}$$

where, again, there is a factor to account for the change in photon number density between ϵ and the final temperature $T < \epsilon$. Note that the final density of n-gravitons/gravitinos is proportional to $T_{\rm R}^{3/2}$. For typical parameter values ($\hat{\sigma} = 8\pi/M_{\rm Pl}^2$, $g_{\rm K} = 1018$, $g_{s\rm K} = 1400$), this is

$$Y_{n(T)} \sim 3 \times 10^{-16} \ \left(\frac{\epsilon}{\rm TeV}\right)^{1/2} \left(\frac{T_{\rm R}}{\rm TeV}\right)^{3/2} \ \exp(-t/\tau_n) \tag{37}$$

at $T \lesssim \epsilon$; the relic abundance during nucleosynthesis would be smaller by a factor $g_{s(T)}/g_{s(T_0)}$ where $g_{s(T_0)} = 43/11$ is the effective number of entropic degrees of freedom at $T \ll m_e$ (taking into account the three decoupled relativistic neutrino species). This is close to the value of 3.36 for the effective number of degrees of freedom contributing to the total energy density, so for convenience we have ignored the small difference.

The energy density in the relic KK modes decreases as $T^{17/3}$ for $T > \epsilon$, i.e. *faster* than the energy density in 'radiation' (including KK excitations) which goes as T^5. (At $T < \epsilon$ the former decreases as $T^{16/3}$, while the latter does so as T^4.) Therefore the universe will become 'radiation' dominated when

$$\sum_n |n| \ \epsilon \ Y_{n(T)} \ n_{\mathrm{rad}(T)} < g_{s(T)} \ \frac{\pi^2 T^4}{30} \ . \tag{38}$$

where $g_{s(T)}$ is taken from eq.(11). Anticipating the bounds which we will find on $T_{\rm R}$, the temperature at equality will be higher than ϵ, therefore the appropriate time-temperature relationship is eq.(24). Since at late times, the number density of KK modes is dominated by the lighter (slowly decaying) particles which have a small exponential suppression factor in Y_n, we may approximate the sum on the left above by an integral,

$$\begin{aligned} \sum_{n=-\infty}^{\infty} |n| \ \exp(-tn^3/\tau_1) \ &\approx \ \int_0^{\infty} \mathrm{d}^2 x \ x \ \exp(-x^3) \left(\frac{t}{\tau_1}\right)^{-2/3} \\ &= \ \frac{2}{3}\,\Gamma\left(\frac{2}{3}\right) \left(\frac{t}{\tau_1}\right)^{-2/3} . \end{aligned} \tag{39}$$

This expression is valid for $t \lesssim \tau_1$; at later times, the KK modes have all decayed, i.e. their number is exponentially suppressed. Therefore 'radiation' domination will occur at

$$T \gtrsim 1.1 \times 10^6 \text{ TeV} \left(\frac{\epsilon}{\text{TeV}}\right)^{5/2} \left(\frac{T_\text{R}}{\text{TeV}}\right)^{-1} \tag{40}$$

corresponding to a time

$$t \lesssim 5.3 \times 10^{-29} \text{ sec} \left(\frac{\epsilon}{\text{TeV}}\right)^{1/2} \left(\frac{T_\text{R}}{\text{TeV}}\right)^{5/2} . \tag{41}$$

It may appear surprising that the higher the reheat temperature, the *less* likely the gravitinos are to matter dominate at a late time. This can be explained as follows; there is only a limited amount of entropy, and when T_R is high, more of it is initially distributed in heavier modes with higher KK number. Since these modes decay more rapidly, they are able to matter dominate only at very early times.

We can now examine the effect of the decaying particles on the abundances of the light elements. The effect of hadronic decays occuring at the beginning of the nucleosynthesis era on neutron-proton interconversions has been studied in detail in ref.[17]. In the time-interval $1 \lesssim t \lesssim 10^2$ sec, the requirement that the ^{4}He mass fraction not be increased above 25% translates into the requirement (see figure 3 in ref.[6]):

$$\sum_n |n| \left(\frac{\epsilon}{\text{TeV}}\right) Y_{n(T)}\, n_{\text{rad}(T)} \lesssim 2 \times 10^{-11}, \tag{42}$$

which corresponds to the bound

$$T_\text{R} \lesssim 21 \text{ TeV} \left(\frac{\epsilon}{\text{TeV}}\right)^{1/3} . \tag{43}$$

For $t \gtrsim 10^4$ sec, the photodissociation of ^{2}H due to the radiation cascades triggered by the decaying particles impose the constraint [6]

$$\sum_n |n| \frac{\epsilon}{\text{TeV}} Y_{n(T)}\, n_{\text{rad}(T)} \lesssim 10^{-13}, \tag{44}$$

or

$$T_\text{R} \lesssim 2.7 \text{ TeV} \left(\frac{\epsilon}{\text{TeV}}\right)^{1/3} . \tag{45}$$

One should however bear in mind that the approximation we used to calculate the KK number density breaks down when T_R becomes comparable to ϵ since the number density is then suppressed by a Boltzmann factor. We can however justifiably assert that the maximum temperature which the universe reached cannot have significantly exceeded the supersymmetry breaking scale.

A precise formulation of the early history of the universe is lacking for these models, however it is clearly essential that there be an inflationary phase [10], followed, as we have shown, by reheating to a temperature no greater than the supersymmetry breaking scale. We now argue that this imposes a conflicting set of requirements on the scalar field Φ which is presumed to drive inflation. In order to account for the amplitude of the scale-invariant density fluctuations probed by *COBE* [18], m_Φ is required to be of

$\mathcal{O}(10^8)$ TeV [19]. For the reheat temperature to be low, the inflaton must then be very weakly coupled to matter fields. However even assuming only gravitational couplings, i.e. a decay rate $\Gamma_\Phi \sim m_\Phi^3/M_{\rm Pl}^2$, the reheat temperature cannot be reduced below

$$T_{\rm R} \sim \left[\frac{g_{s{\rm K}}(T)}{\epsilon}\right]^{-1/4} (\Gamma_\Phi M_{\rm Pl})^{1/2} \sim 10^3 \ {\rm TeV} \left(\frac{\epsilon}{\rm TeV}\right)^{1/4} \left(\frac{m_\Phi}{10^8 {\rm TeV}}\right)^{3/2}, \qquad (46)$$

where $T_{\rm R}$ has been obtained by equating Γ_Φ to the Hubble rate H. In fact, the only mass scale in these models, apart from the Planck scale, is the supersymmetry breaking scale. One could then envisage a scenario with two epochs of inflation [8]; the first stage to create the correct level of density fluctuations, and the second, with $m_\Phi \sim \epsilon$, to remove the KK states. However assuming gravitational couplings as above, the reheat temperature is then of $\mathcal{O}$(keV), i.e. too low for even primordial nucleosynthesis to occur. On the other hand if the inflaton is coupled directly, the reheat temperature would be too high, unless the gauge coupling is unnaturally small, of $\mathcal{O}(10^{-4})$ [8]. A way out may be to introduce an intermediate scale of unification, with a low-mass 'flaton' singlet remaining after symmetry breaking [20]. However the reheat temperature is then of $\mathcal{O}$(MeV), i.e. barely high enough for nucleosynthesis, and further, the Affleck-Dine mechanism must be invoked in order for the baryons to be synthesized just beforehand.

In conclusion, it appears difficult to construct a consistent cosmological history for models with spontaneous supersymmetry breaking. In view of their many other advantageous features [2], possible resolutions to this problem should be pursued.

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