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AFTER PRIMORDIAL INFLATION

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ABSTRACT

We consider the history of the early Universe in the locally supersymmetric model we have previously discussed. We pay particular attention to the requirement of converting the quanta of the field which drives primordial inflation (inflatons) to ordinary particles which can produce the cosmological baryon asymmetry without producing too many gravitinos. An inflaton mass of about 10^{13} GeV (a natural value in our model) produces a completely acceptable scenario.

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An inflationary Universe[1)-8)] certainly presents the most promising solutions to the isotropy, flatness, homogeneity, ..., problems. In general, such a scenario is realized via the effects of a supercooled phase transition at or before the grand unified theory (GUT) era. If the Universe becomes dominated by the vacuum energy density due to a scalar field for a sufficiently long time, the Universe enters an approximate De Sitter state and begins to expand exponentially. We will not discuss here many of the problems which plagued both the original[1),2)] and new[3)-5)] inflationary models [see, e.g., Refs. 6), 7) and 9)]. Instead we will concentrate on primordial inflation in $N = 1$ supergravity[8)].

Primordial supersymmetric inflation[7)] was originally introduced to cure difficulties with the new inflationary scenario[3)-5)]. In general, supersymmetric models allow for considerably flatter potentials than in standard SU(5), thus allowing a longer inflationary period. Inflation at the GUT scale of $\sim 10^{15}$ GeV still required precise adjustments of scalar self-couplings. These fine-tunings seem to disappear when the inflation scale is moved up towards the Planck mass, $M_P \approx 10^{19}$ GeV. To compensate for first order gravitational (FOG) effects the present authors and Tamvakis have recently considered primordial inflation in $N = 1$ supergravity[8)].

In order to achieve a sufficient amount of inflation, we had to rely upon non-renormalizable interactions in the superpotential. As $N = 1$ supergravity is non-renormalizable anyway, there is no reason to exclude such terms from the superpotential. This model must then be viewed as an effective theory which correctly describes interactions at scales below the Planck scale M_P. It is precisely the non-renormalizable interactions which we have exploited further to discuss the SU(5) to SU(3) × SU(2) × U(1) phase transition[10)] at a scale $\Lambda_5 \simeq 10^9$ GeV where Λ_5 is the scale of confinement in the SU(5) phase. These considerations cleared a number of difficulties in realizing a supersymmetric cosmological scenario[11)] in which the Universe was not trapped in the SU(5) phase. This led to a prediction[10)] of the GUT scale $M_X^2 \approx \mu M_P$ where μ is the supersymmetry breaking scale. Non-renormalizable couplings were also found[10)] to be responsible for providing for the cosmological baryon asymmetry. In the present letter, we would like to take a closer look at what happens between inflation and the breakdown of SU(5), in particular, the reheating of the Universe and the solution of the gravitino mass problem[12),13)].

Let us begin therefore by writing down all the essential ingredients. Our "complete" superpotential can be divided into three parts

$$f = f_I + f_5 + f_B \tag{1}$$

$$f_I = m^3\left[\sum_{n=0}^{\infty} \frac{\lambda_n}{n+1}\left(\frac{\varphi}{M}\right)^{n+1}\right] \tag{2}$$

$$f_5 = \frac{a_1}{M} X^4 + \frac{a_2}{M^2} X^2 T_R(\Sigma^3) + \mu^2(z+\Delta) \tag{3}$$

$$f_B = \frac{a_3}{M} X Y \bar{H} H + a_4 M \bar{\theta}\theta + \frac{a_5}{M}\bar{\theta}\Sigma^2 H + \frac{a_6}{M}\bar{H}\Sigma^2\theta + a_7\mu Y^2 + a_8 Y^3 + \text{Yukawa couplings} \tag{4}$$

ϕ is the inflaton field, an SU(5) singlet which acquires a vacuum expectation value (VEV) $v = \langle 0|\phi|0\rangle \simeq M = M_P/\sqrt{8\pi} \approx 2.4 \times 10^{18}$ GeV and drives inflation. X and z are also SU(5) singlets while Σ is an SU(5) $\underset{\sim}{24}$. f_5 is responsible for breaking SU(5) as well as local supersymmetry. f_B is responsible for the cosmological baryon asymmetry. H and $\bar{H}$ are a 5 and $\bar{5}$ of Higgs supermultiplets, θ and $\bar{\theta}$ are a 50 and $\overline{50}$, and Y is a singlet. We employ the Higgs doublet-triplet splitting mechanism of Ref. 14), which can give masses of 10^{10} GeV to the colour triplets contained in the H and $\bar{H}$ fields, while leaving their weak doublet partners light. The H and $\bar{H}$ fields also couple to quark and lepton superfields.

We now turn to a discussion of the parameters in the partial superpotentials (2)-(4). The mass scale m is an overall scale for f_I and is determined by requiring sufficient inflation and reasonable density perturbations; in Ref. 8) it was found that $m \sim 10^{-2}M$. The couplings λ_i are also selected by the same requirements and a viable solution has $\lambda_0 \sim \lambda_3 \sim 10^{-1}$, $\lambda_1 = 0$, $\lambda_2 \sim 10^{-3}$ while the rest are adjusted to keep the final cosmological constant at that scale equal to zero as well as preserving supersymmetry. This requires

$$\sum_n \lambda_n\left(\frac{v}{M}\right)^n = \sum_n \frac{\lambda_n}{n+1}\left(\frac{v}{M}\right)^{n+1} = 0;\ \sum_n n\lambda_n\left(\frac{v}{M}\right)^{n-1} \neq 0 \tag{5}$$

Thus $v \simeq M$ at the global minimum of $V(\phi)$ in order to avoid fine-tunings of parameters in (2). We take $v = M$ exactly for convenience. The inflaton field, like all other chiral fields in our model, is expected to be a bound state of some more fundamental theory such as N = 8 supergravity.

In the remaining parts of the superpotential we will assume that the couplings a_i are all of order 10^{-1} - 1. The Polonyi term $\mu^2(z + \Delta)$ is used to break local supersymmetry at a scale $\mu \sim O(10^{11})$ GeV (corresponding to a gravitino mass $m_{3/2}$ of about 100 GeV). Δ is adjusted to cancel the cosmological constant at the scale μ. Finally, the chiral superfield Y is included to facilitate the generation of a baryon asymmetry through the decays of the Higgs fields H and $\bar{H}$. [In practice two sets of H and $\bar{H}$'s are needed[15].]

Let us now carefully run through the sequence of events which begin at the Planck time. The effective potential is given in terms of the superpotential by[16]

$$V = e^{\sum_\alpha |\xi_\alpha|^2/M^2}\left[\sum_\alpha |f_{\xi_\alpha}|^2 - \frac{3}{M^2}|f|^2\right] \tag{6}$$

where

$$f_\xi = \frac{\partial f}{\partial \xi} + \xi^* f/M^2 \tag{7}$$

for all fields ξ. The inflationary superpotential f_I (2) yields a vacuum energy density $V(\phi = 0) = \lambda_0 m^6/M^2 \approx 10^{-14}M^4$. During the roll-over of the field ϕ to its global minimum at $\phi = M$ the Universe enters an approximate De Sitter stage characterized by a Hubble parameter $H = (1/\sqrt{3})\lambda_0 m^3/M^2 \approx 10^{11}$ GeV. The Universe feels an effective temperature $T_H = H/2\pi \approx 10^{10}$ GeV at the beginning of the inflationary epoch. The vacuum energy density at the completion of the roll-over is converted into the inflatons ϕ. Thus at this stage, the Universe is inflaton dominated. All other particle species have been diluted by inflation and only ϕ's can be produced because no direct couplings of ϕ to other fields have been introduced. The inflaton mass is given by

$$m_\phi = \frac{\partial^2 f}{\partial \phi^2} = \frac{m^3}{M^2}\sum_n n\lambda_n \approx 10^{13}\,\mathrm{GeV} \tag{8}$$

[assuming that the higher order couplings are $O(1)$] so that the inflatons will be relativistic with a temperature

$$T_\phi \approx 3\times 10^{-4} M \approx 7\times 10^{14}\,\mathrm{GeV} \tag{9}$$

The inflatons will eventually decay through their gravitational interactions induced by the term included in (6),

$$\left| \frac{\partial f}{\partial \varphi} + \varphi^* f/M^2 \right|^2 \tag{10}$$

Remembering that both f_I and $\partial f/\partial\phi$ are zero at the minimum $\phi = M$, we can expand about the minimum and obtain the following interaction for the shifted field ϕ

$$\frac{m_\varphi \varphi}{M} \langle\varphi\rangle f = \frac{m_\varphi}{M} \varphi \left(f_5 + f_B\right) \tag{11}$$

In particular, we see from (3) and (4) that the dominant decay modes will be $\phi \rightarrow 3Y$, HT_3T_3 and $\bar{H}T_3\bar{F}_3$ (where T_3 and $\bar{F}_3$ are the third generation of quark and lepton fields), with a rate given by

$$\Gamma_\varphi \sim m_\varphi^3/M^2 \tag{12}$$

for superpotential couplings of $O(1)$. [The first and second generations have only non-renormalizable couplings[17] and so will not be produced directly.]

Because this decay rate (12) is relatively small, the Universe becomes matter (ϕ) dominated when T_ϕ falls below m_ϕ. The ϕ's finally decay when

$$\Gamma_\varphi \sim H \tag{13}$$

or

$$m_\varphi^3/M^2 \sim m_\varphi^{1/2} T_{D\varphi}^{3/2}/M \tag{14}$$

so that the decay temperature

$$T_{D\varphi} \sim m_\varphi^{5/3}/M^{2/3} \sim 10^{-8} M \tag{15}$$

After ϕ decay, the Universe will heat up marginally to a temperature

$$T_R \sim \left(m_\varphi T_{D\varphi}^3\right)^{1/4} \sim m_\varphi^{3/2}/M^{1/2} \sim O(10^{10}-10^{11})\,\mathrm{GeV} \tag{16}$$

SU(5) breaking will have occurred at a temperature $T \sim \Lambda_5 \sim O(10^9 - 10^{10})$ GeV through f_5 (3) as discussed in detail in Ref. 10)[*]. We note at this point that if m_ϕ was much smaller than (8) the Universe will never have reheated and generated a baryon asymmetry. As we will also see if m_ϕ was much larger than (8) the Universe would have reheated too much thereby causing a problem with gravitino production[12].

As was discussed in Ref. 10), the superpotential f_5 (3) yields an effective potential which at $T = 0$ has as its global minimum the $SU(3) \times SU(2) \times U(1)$ symmetric state. This differs significantly from the situation in global supersymmetry[11] in which the lowest states were all degenerate and supersymmetric. It was here once again that the use of non-renormalizable interactions played a key role to lift the degeneracy as well as lower the barrier between the SU(5) and $SU(3) \times SU(2) \times U(1)$ phases. In global supersymmetry the barrier could only be lowered from $O(10^{16}\ \text{GeV})^4$ to $O(10^9\ \text{GeV})^4$ by introducing couplings of $O(10^{-12})$! The non-renormalizable terms in (3) produce the same effect with all couplings O(1). In addition the superpotential (3) predicts the GUT scale M_X. The $SU(3) \times SU(2) \times U(1)$ minimum occurs when Σ picks up a VEV; this VEV can be shown[10] to yield

$$M_X^2 \simeq \mu M \simeq (10^{15}\ \text{GeV})^2 \tag{17}$$

The cosmological baryon asymmetry will be generated through the series of decays

$$Y \to H_3,\ \bar{H}_3 \to \text{matter} \tag{18}$$

Through the non-renormalizable couplings in f_B (4) the fields Y decay only to H and $\bar{H}$, i.e., through the term

$$a_3 \langle X \rangle Y H \bar{H} / M \tag{19}$$

where $\langle X \rangle$ takes the value

$$\langle X \rangle / M = (\mu/M)^{3/16} \sim O(10^{-6}) \tag{20}$$

This leads to a decay rate for the field Y

$$\Gamma_Y \sim a_3^2 \langle X \rangle^2 m_Y / M^2 \tag{21}$$

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*) See note added to that paper.

and hence the Y's will decay when

$$\frac{\langle X \rangle^2}{M^2} m_Y \sim m_Y^{1/2} T_{DY}^{3/2}/M \tag{22}$$

or

$$T_{DY} \sim 10^{-8} m_Y^{1/3} \sim 10^7 \, GeV \tag{23}$$

(for a mass $m_Y \sim O(10^{10})$ GeV the Universe is again matter dominated, this time by Y's). Remember also that the ratio $n_Y/n_\gamma \gtrsim 1$ where n_γ is the number density of all relativistic species. Thus, when the Higgs triplets H and $\bar{H}$ are produced, they will also have $n_H/n_\gamma \gtrsim 1$. The mass of these triplets will also be $O(10^{10})$ GeV. Needless to say that with $O(10^{10})$ GeV Higgs bosons, protons will decay predominantly into μK and/or νK[11),15)] with the nominal lifetime $10^{31\pm1}$ years. In equilibrium one would expect that at 10^7 GeV, $n_H/n_\gamma \ll 1$ foregoing any possibility of generating a baryon asymmetry. In this scenario, the H's are wildly out of equilibrium and an asymmetry of order $n_B/n_\gamma \sim 10^{-2}\Delta B$ (where ΔB is the net baryon number per decay) can be produced. We note also that because SU(5) is broken at $10^9 - 10^{10}$ GeV, this scenario offers both the possibility of solving the magnetic monopole problem as well as leaving a possibility of a non-zero (and perhaps observable) flux which is not an option of standard scenarios incorporating inflation.

Finally, we would like to comment on the solution of the gravitino problem[12),13)]. Gravitinos will be produced either in the decays of the inflaton or through scatterings with other supersymmetric particles such as the photino. A large abundance of gravitinos with a mass of order 100 GeV can be a problem after their decay. In particular, one has to worry that their decay does not produce too much entropy so as to drown out the baryon asymmetry[12)]. Let us define the relative abundance of gravitinos

$$Y_{3/2} = n_{3/2}/n_\gamma \tag{24}$$

at the time of their production. The temperature at which they decay will be determined by $Y_{3/2}$. If $Y_{3/2}$ is large, the Universe will become matter dominated by gravitinos which decay when

$$\Gamma_{3/2} \sim m_{3/2}^3/M^2 \sim Y_{3/2}^{1/2} m_{3/2}^{1/2} T_{D3/2}^{3/2}/M \tag{25}$$

or

$$T_{D3/2} \sim m_{3/2}^{5/3} / Y_{3/2}^{1/3} M^{2/3} \tag{26}$$

whereas for small $Y_{3/2}$ the Universe remains radiation dominated and gravitinos decay when

$$T_{D3/2} \sim m_{3/2}^{3/2} / M^{1/2} \tag{27}$$

For gravitino masses $\sim$100 GeV, gravitinos decay after nucleosynthesis. Thus, we must require that the entropy does not increase by a factor of at most O(10) *) This translates to the requirement that

$$\rho_{3/2}/\rho_\gamma = m_{3/2} T_{D3/2}^3 Y_{3/2} / T_\gamma^4 < O(10) \tag{28}$$

or

$$Y_{3/2} < \frac{T_{D3/2} N(T_P)}{m_{3/2}} \cdot 10 \sim 10 N(T_P) \left(\frac{m_{3/2}}{M}\right)^{1/2} \sim O(10^{-5}) \tag{29}$$

where $N(T_P)$ is the total number of degrees of freedom when the gravitinos were produced, so that $T_\gamma^3/T_{3/2}^3 \sim N(T_P)$ due to annihilations. Also we have used the radiation dominated value of $T_{D3/2}$ (27).

*) Any entropy increases larger than this would require a missing mass problem for the solar neighbourhood in order to remain consistent with the successful prediction of big bang nucleosynthesis[18].

Let us now see what values of $Y_{3/2}$ arise in this scenario. We assume that the Polonyi field z has a VEV of order M at all times. There are several reasons for this. First of all, $z \simeq M$ is the only local minimum of the z potential (in contrast to the ϕ potential, which has a local minimum at $\phi = 0$), and finite temperature contributions to the z potential (which are never larger than order $m_{3/2}^2 T^2$) have no reason to pick out any other value of z before the Planck time. If this is the case ($z \simeq M$ at all times), then primordial gravitinos and z particles will be depleted during the inflationary era. They can be recreated by either inflaton decays or scatterings when the Universe is ϕ dominated. The decays $\phi \to \psi\tilde{G}$ or $\psi \to \phi\tilde{G}$ (where ψ is the fermion partner of the ϕ and $\tilde{G}$ is the gravitino) can be forbidden kinematically by requiring $|m_\phi - m_\psi| < m_{3/2}$. Gravitino production can still occur in a ϕ decay such as $\phi \to$ light scalar + light fermion + gravitino, but the decay rate for this process will be suppressed relative to a process without a final state gravitino by a factor $(m_\phi/M)^2$, yielding the safe value

$$Y_{3/2} \simeq \left(\frac{m_\varphi}{M}\right)^2 \simeq 10^{-10} \tag{30}$$

If, however, they are produced by scatterings $X + \hat{\gamma} \to X + \tilde{G}$, where $\hat{\gamma}$ is a gauge fermion, their production rate could be as large as*)

$$\Gamma_p \sim \alpha N(T)\, T^3/M^2 \tag{31}$$

where α is a gauge coupling. In this case

$$Y_{3/2} = \frac{\Gamma_p}{H} = \frac{\alpha\sqrt{N}\, T_R}{M} \tag{32}$$

where T_R is the reheated temperature (16) after inflation. Thus, in this model $Y_{3/2} \sim O(10^{-8})$ which is again a safe value. We note in passing that standard inflationary scenarios which reheat to a temperature $T > 10^{13}$ GeV (normally necessary for baryon generation) will still be plagued with a gravitino mass problem.

In conclusion, we would like to make the following interesting observation. Recall that in order to create a baryon asymmetry we needed $T_R \gtrsim O(10^{10})$ GeV and the gravitino problem required $T_R < O(10^{13})$ GeV. This translates to a mass range for the inflaton,

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*) We would like to thank J. Ellis for communicating this observation by S. Weinberg.

$$10^{-5} \lesssim m_\phi/M \lesssim 10^{-3} \tag{33}$$

It is surprising that if the inflaton is to yield a density fluctuation spectrum compatible with galaxy formation[19], we expect $m_\phi/M \approx 10^{-5}$ [see Eq. (8)]. Thus, three different constraints - galaxy formation, baryon production and gravitino suppression - all point to the same range of parameters. We do not believe that this is merely a coincidence. We feel therefore, that there exists at least a prototype for a supersymmetric cosmological model which inflates naturally without severe parameter adjustments and simultaneously creates a baryon asymmetry and density fluctuations suitable for galaxy formation without overproducing magnetic monopoles or gravitinos. Included also are natural Higgs doublet-triplet splittings[14], as well as predictions of the GUT scale and proton decay modes (predominantly to μK and/or νK with a lifetime of $10^{31\pm1}$ years). We stress again that all of these results can be traced back directly to non-renormalizable gravitational effects.

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