COSMOLOGICAL PROBLEMS FOR THE POLONYI POTENTIAL

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Received 5 July 1983

We study the cosmological implications of $N = 1$ supergravity with the Polonyi potential. We find that for typical values of the gravitino mass (10^2–10^3 GeV) the universe goes through a late period of reheating (i.e., from a temperature of about 10^{-7} MeV to 10^{-2} MeV). Any baryon-to-photon ratio is thus diluted by an unacceptable 15 orders of magnitude, with no hope of regeneration.

Recently Albrecht et al. [1] have studied the implications of the Geometric Hierarchy (GH) model [2] for the early universe. The model is (for the purpose of this discussion) characterized by a naturally shallow scalar potential $V(\phi)$ with an intermediate scale energy density of order μ^4 ($\mu \sim 10^{12}$ GeV) at the origin and by a vacuum expectation value ϕ_0 of order the Planck mass, $M_{Pl} \sim 10^{19}$ GeV. The scalar field ϕ is a linear superposition of an SU(5) adjoint and singlet. As a result $\phi_0 \sim M_{Pl}$ is also the grand unification scale in this model and several fields obtain masses proportional to ϕ.

The high temperature ($T \gg \mu$) behavior of the free energy $V^T(\phi)$ was shown [3] to be dominated by the free particle term $V^T(\phi) \sim -N(\phi)T^4$ where $N(\phi)$ is the number of states with mass much less than T for a given value of the scalar field ϕ. [Note $N(0) > N(\delta > T)$.] The free energy for $T \gg \mu$ is thus minimized for $\phi \simeq 0$. For $T < \mu$ the free energy obtains its zero temperature form with the minimum at ϕ_0, and the scalar field evolves from $\phi \sim 0$ to $\phi = \phi_0$ by "rolling down the potential." It was shown that the universe expands exponentially during the evolution of ϕ and may inflate many orders of magnitude [1]. However, as a result of the decoupling which occurs in this class of supersymmetric theories [2,4] the universe does not reheat *sufficiently* to solve any of the well-known problems of cosmology [5]. The situation is in fact worse than this. The energy stored in the ϕ field is eventually released and increases the entropy by about 15 orders of magnitude. The baryon-to-photon ratio is thus suppressed

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by the same unacceptable factor. The maximum reheat temperature is about 10^{-2} MeV, which is too low to regenerate the baryon-to-photon ratio by known mechanisms [6]. Thus it was concluded that either (i) the observed baron-to-photon ratio has an entirely new origin, i.e., it is not due to microscopic baryon violating, processes, which are ineffective at temperatures less than 10^{10} GeV or (ii) the GH model is a cosmological disaster.

In this letter we perform a similar analysis, using the simple potential introduced by Polonyi [7] for $N = 1$ supergravity. It has the property of spontaneously breaking supersymmetry (SUSY) while allowing for the possibility of fine tuning the cosmological constant to zero [8]. Moreover it has recently been used as a simple example of the so-called "hidden sector" whereby supergravity effects induce the weak breaking scale [9]. It has characteristic properties similar to those of GH i.e., μ^4 energy density at the origin, ϕ_0 is of order M_{Pl}; and ϕ decouples from light states. It differs from GH however in one way: There are no states which obtain mass proportional to ϕ. Thus the high temperature behavior of the Polonyi model is not the same as GH. Below, we study the high temperature behavior of the Polonyi model and the fate of the early universe.

A key ingredient in our analysis is the decoupling of the ϕ field. The ϕ field interacts with radiation and with itself only through gravitational strength interactions. As a result, the thermal averaged interaction rate of the ϕ field at temperature T must be on dimensional ground of order

$$\langle n\sigma v\rangle \sim T^5/M^4 \tag{1}$$

where $M = M_{\mathrm{Pl}}/\sqrt{8\pi} = 2.4\times 10^{18}$ GeV, and n is the number density ($n \sim T^3$).

Another manifestation of decoupling is the ϕ decay rate [1],

$$\Gamma \sim m_\phi^3/M^2 \sim \mu^6/M^5 . \tag{2}$$

These rates are to be compared to the expansion rate of the universe which is given by the expression

$$\dot{R}/R = (\rho/3M^2)^{1/2} . \tag{3}$$

For a radiation dominated universe, the energy density ρ is

$$\rho \cong \tfrac{1}{30} N\pi^2 T^4 \tag{4}$$

where $N = N_{\mathrm{b}} + \frac{7}{8}N_{\mathrm{f}}$, and $N_{\mathrm{b(f)}}$ are the number of light bosonic (fermionic) helicity states. In a minimal low energy supersymmetric model $N = \frac{15}{4}\times 61$ (61 corresponds to 45 quarks and leptons, 12 gauge fields and four Higgs fields) or $N \simeq 225$. Thus we have

$$\dot{R}/R \simeq T^2/M . \tag{5}$$

For the ϕ field to be in thermal equilibrium we must have

$$\langle n\sigma v\rangle \geqslant \dot{R}/R . \tag{6}$$

Clearly this will occur only for $T \geqslant M$.

We shall henceforth *assume* that ϕ is in thermal equilibrium at $T = M$. This assumption is implicit in our subsequent use of the free energy $V^{\mathrm{T}}(\phi)$ to obtain the initial value of ϕ.

For $T < M$ the ϕ field decouples and we shall treat the "low" momentum modes (corresponding to wavelengths of order the causal horizon) as a classical field satisfying the classical equations of motion.

It has been demonstrated by Cremmer et al. [8] that the general matter coupling to simple supergravity is determined by two arbitrary functions $(\phi_a, \phi_a^*, \phi, \phi^*)$ and $f_{\alpha\beta}(\phi_a, \phi)$ (where α, β are indices in the adjoint representation of any gauge group, ϕ is the Polonyi field and ϕ_a, $a = 1, \ldots, N-1$ are all other scalar fields in the theory.) For our present calculations we shall take $f_{\alpha\beta} = \delta_{\alpha\beta}$. We then take

$$\mathcal{G} = -|\phi|^2/M^2 - \ln(|g|^2/M^6) - |\phi_a|^2/M^2 - \ln|h|^2 , \tag{7}$$

where $g \equiv \mu^2(\beta_0 M + \phi)$ is the Polonyi potential with $\beta_0 = 2-\sqrt{3}$, $\mu \sim 10^{10}$ GeV, $M = M_{\mathrm{Pl}}/\sqrt{8\pi} = 2.4\times 10^{18}$ GeV, and $h \simeq h(\phi_a)$ is an arbitrary nonvanishing function of all other fields. In terms of $\mathcal{G}$ we obtain the classical potential

$$V_0 = \exp(-\mathcal{G})(|\mathcal{G}'|^2 + |\mathcal{G}_a|^2 - 3) , \tag{8}$$

where $\mathcal{G}' \equiv \partial\mathcal{G}/\partial\phi$ and $\mathcal{G}_a \equiv \partial\mathcal{G}/\partial\phi_a$.

We assume $\mathcal{G}_a = 0$ i.e., that all supersymmetry breaking is in the Polonyi sector. In order to treat the potential as a function of ϕ alone we let $\phi_a = 0$ (at high T), and consider the real dimensionless variable $\tilde{\phi} \equiv \mathrm{Re}\,\phi/M$. In this direction we find:

$$V_0(\tilde{\phi}) = \mu^4 \exp(\tilde{\phi}^2)[1 - 3\beta_0^2 + (\beta_0^2 - 1)\tilde{\phi}^2 + 2\beta_0(\tilde{\phi}^2 - 2)\tilde{\phi} + \tilde{\phi}^4] ,$$
$$\tilde{\phi}_0 = \sqrt{3} - 1 , \quad V(\tilde{\phi}_0) = \partial V/\partial\phi|_{\tilde{\phi}_0} = 0 . \tag{9}$$

[We have let $h(\phi_a = 0) \equiv 1$.]

The one-loop corrections to the finite temperatures free energy $V^{\mathrm{T}}(\phi)$ are given by the following expression [10]

$$V^{\mathrm{T}}(\phi) = V_0(\phi) + V_1(\phi) ,$$
$$V_1(\phi) = -\tfrac{1}{90}\pi^2 N(T) T^4 + \tfrac{1}{24}T^2(\mathrm{tr}\, M_{\mathrm{S}}^2 + \tfrac{1}{2}\,\mathrm{tr}\, M_{\mathrm{f}}^2 + 3\,\mathrm{tr}\, M_{\mathrm{V}}^2) + [\Lambda^2/(4\pi)^2](\mathrm{tr}\, M_{\mathrm{S}}^2 - \mathrm{tr}\, M_{\mathrm{f}}^2 + 3\,\mathrm{tr}\, M_{\mathrm{V}}^2) + O(M^4) , \tag{10}$$

where Λ is a momentum space cut-off of order M, $N(T) = N_{\mathrm{b}} + \frac{7}{8}N_{\mathrm{f}}$, and $M_{\mathrm{S,f,V}}^2$ are respectively scalar, fermion and vector mass matrices squared (the gravitino is included in M_{f}).

We are only interested in terms which depend on $\tilde{\phi}$. For the Polonyi model $N(T)$ is independent of $\tilde{\phi}$ for T larger than the weak scale $\sim\mu^2/M$. M_{V}^2 is also independent of $\tilde{\phi}$. We are therefore only interested in M_{S}^2 and M_{F}^2. We find

$$\mathrm{tr}\, M_{\mathrm{S}}^2 = 2(\partial^2 V/\partial\phi_a\partial\phi_a^* + \partial^2 V/\partial\phi\partial\phi^*)$$
$$= 2\exp(-\mathscr{G})[|\mathscr{G}_{ab}|^2 + |(\mathscr{G}')^2 - \mathscr{G}''|^2 - 2 + (N-1)(|\mathscr{G}'|^2 - 2)] ,$$
$$\mathrm{tr}\, M_{\mathrm{f}}^2 = 2\exp(-\mathscr{G})[|\mathscr{G}_{ab}|^2 + 2] , \tag{11}$$

where

$$\mathscr{G}' = -(\tilde{\phi} + 1/(\beta_0 + \tilde{\phi})) , \quad \mathscr{G}'' = 1/(\beta_0 + \tilde{\phi})^2 ,$$
$$\exp(-\mathscr{G}) = \mu^4 \exp(\tilde{\phi}^2)(\beta_0 + \tilde{\phi})^2 , \tag{12}$$

and $N \sim 50$ is the number of matter fields. Thus the leading contribution to V^{T} is given by

$$V^{\mathrm{T}}(\tilde{\phi}) \simeq V_0(\tilde{\phi}) + [\kappa + \kappa'(T^2)][V_0(\tilde{\phi}) + \mu^4 \exp(\tilde{\phi}^2)(\beta_0 + \tilde{\phi})^2] , \tag{13}$$

where

$$\kappa = [2(N-1)/(4\pi)^2]\Lambda^2/M^2 ,$$
$$\kappa' = \tfrac{2}{24}(N-1)T^2/M^2 .$$

The minimum of the zero temperature potential [10] is shifted to

$$\langle\tilde{\phi}\rangle_{T=0} = \sqrt{3}(\zeta - 1) + O(1-\zeta)^2 ,$$
$$\zeta = 1 - [(N-1)/3(4\pi)^2]\Lambda^2/M^2 , \tag{14}$$

and we must renormalize β_0 so that the energy at the minimum remains zero. We have

$$\beta = (2 - \sqrt{3})\zeta + O(1-\zeta)^2 . \tag{15}$$

In the same approximation we find the finite T minimum of $V^{\mathrm{T}}(\tilde{\phi})$ given by

$$\langle\tilde{\phi}\rangle_T = \langle\tilde{\phi}\rangle_{T=0} - \tfrac{1}{24}(N-1)T^2/M^2 , \tag{16}$$

and the energy in the potential is

$$V^{\mathrm{T}}(\langle\tilde{\phi}\rangle_T) \sim 23\mu^4 \exp(\tilde{\phi}_0^2)\tfrac{1}{24}(N-1)T^2/M^2 . \tag{17}$$

The above analysis suggests that a reasonable initial condition for the Polonyi field for $T \sim M$ (and momenta small compared to m_ϕ) is given by

$$\phi_i = \langle\tilde{\phi}\rangle_{T=M} M \lesssim \langle\tilde{\phi}\rangle_{T=0} M . \tag{18}$$

Note ϕ_i is not at the origin, but since the shift away from the minimum is of order 1 in M, we still have an intermediate scale energy density in the field ϕ.

The evolution of ϕ is determined by the classical equations of motion

$$\dot{\rho}_\phi + 3(\dot{R}/R)\dot{\phi}^2 + \Gamma_\phi\dot{\phi}^2 = 0 ,$$
$$\dot{\rho}_r + 4(\dot{R}/R)\rho_r - \Gamma_\phi\dot{\phi}^2 = 0 ,$$
$$(\dot{R}/R)^2 = \tfrac{1}{3}M^{-2}(\rho_\phi + \rho_r) , \tag{19}$$

with

$$\rho_\phi = \tfrac{1}{2}\dot{\phi}^2 + \tfrac{1}{2}(R/R_M)^{-2}(\partial_i\phi)^2 + \tilde{V}^{\mathrm{T}}(\phi) ,$$
$$\rho_r = \tfrac{1}{30}\pi^2 N(T) T^4 ,$$
$$\Gamma_\phi = \mu^6/M^5 . \tag{20}$$

In (20) $\tilde{V}^{\mathrm{T}}(\phi)$ is the internal energy given by

$$\tilde{V}^{\mathrm{T}}(\phi) = V(\phi) - T\partial V^{\mathrm{R}}(\phi)/\partial T - \rho_r ,$$

in the definition of π_ϕ we have kept the covariant spatial gradient term, $\nabla_i = (R/R_M)^{-1}\,\partial_i$, since we consider only ϕ with wavelength of order the causal horizon size l_m at $T = M$ (this is the longest wavelength mode which could be thermalized at $T = M$). We are free to choose the initial scale size $R_M = l_M$. In that case eq. (20) becomes

$$\rho_\phi = \tfrac{1}{2}\dot{\phi}^2 + \tfrac{1}{2}l_M^{-2}(R_M/R)^2\phi^2 + \bar{V}^T(\phi) .$$

We will consider two extreme choices for l_M:

(i) $l_M \simeq 10^{28}/M$.

(ii) $l_M \simeq 1/M$.

Choice (i) corresponds to initial conditions which are necessary in order to be consistent with the homogeneity, isotropy and flatness of the observed universe assuming there is *never* an inflationary epoch in the early universe. The more reasonable choice (ii) could be consistent with the observed universe only in an inflationary universe scenario.

We shall analyze these two cases separately.

(i) $l_M \simeq 10^{28}/M \simeq 10^{10}\,\mathrm{GeV}^{-1}$. In this case l_M is much larger than the ϕ Compton wavelength and the spatial gradient term can be neglected. The finite temperature corrections of the preceding subsection are thus relevant. There are then two relevant regimes for the evolution of ϕ. Consider first $M > T \geqslant \mu$.

In this regime the universe is radiation dominated, and

$$\dot{R}/R \simeq T^2/M \gg m_\phi . \tag{21}$$

In the equation for ϕ the damping term dominates and $\dot{\phi} \simeq 0$. The ϕ field thus slowly rolls down the potential and loses very little energy in the process. In fact in the time it takes to make a few oscillations about the minimum, the temperature drops from M to μ and only a small fraction ($f \sim \frac{1}{2}$) of the initial energy was redshifted away. Γ_ϕ is totally irrelevant on this timescale.

For $T < \mu$, $\dot{R}/R \leqslant m_\phi$, and the oscillation frequency of ϕ is large compared to the expansion rate. During this period, to a good approximation $\langle\dot{\phi}^2\rangle_{\text{time averaged}} \simeq \rho_\phi$. As a result the energy density ρ_ϕ redshifts as $1/R^3$ compared to the radiation energy density $\rho_r \sim 1/R^4$. The first indication of a possible problem is obtained by setting $\Gamma_\phi = 0$ and comparing ρ_ϕ to ρ_r at today's radiation temperature. We have $f \sim \frac{1}{2}$ where $f \equiv \rho_\phi/\rho_r$ at $T = \mu$. Therefore at $T \leqslant \mu$,

$$\rho_\phi/\rho_r = (\mu/T)f \tag{22}$$

which is 10^{23} today at $T = 2.7\,\mathrm{K} \sim 10^{-13}\,\mathrm{GeV}$. Thus, the energy in the ϕ field would dominate the energy density of the universe. This corresponds to a Hubble expansion rate which is ~ 20 orders of magnitude larger than the presently observed rate. Obviously this is a disaster.

Note however $\Gamma_\phi \neq 0$, so the energy in ϕ will eventually be converted into radiation. Let's estimate the temperature T_D at which this energy is released and the resultant entropy created. During the epoch of ρ_ϕ domination, the universe is "matter" dominated, and

$$(\dot{R}/R)^2 = \rho_\phi/3M^2 = [\rho_\phi(T_0)(T/T_0)^3]/3M^2 . \tag{23}$$

Using $\rho_\phi(\mu) = \mu^4$, we obtain the age of the universe, t, given in terms of the temperature as

$$t = (\mu/T)^{3/2}\sqrt{3}M/\mu^2 . \tag{24}$$

If we assume the energy in ρ_ϕ is instantly converted to radiation, when ϕ decays at $t = \Gamma_\phi^{-1}$, the initial temperature is

$$T_D = (3M^2\Gamma_\phi^2/\mu)^{1/3} \simeq 4\times 10^{-8}\,\mathrm{MeV} , \tag{25}$$

where we have used $M = 2.4\times 10^{18}\,\mathrm{GeV}$, $\mu = 3\times 10^{10}\,\mathrm{GeV}$, and Γ_ϕ given in (2). The energy density in the ϕ field at T_D is converted into radiation with a temperature $T_f^4 = \rho_\phi(T_D)$. The entropy increase is given by $\Delta \equiv T_f^3/T_D^3$:

$$\Delta = (\mu^4/3\Gamma_\phi^2M^2)^{1/4} \simeq 5\times 10^{15} , \tag{26}$$

for the same values of μ, M, and Γ_ϕ. In fig. 1,

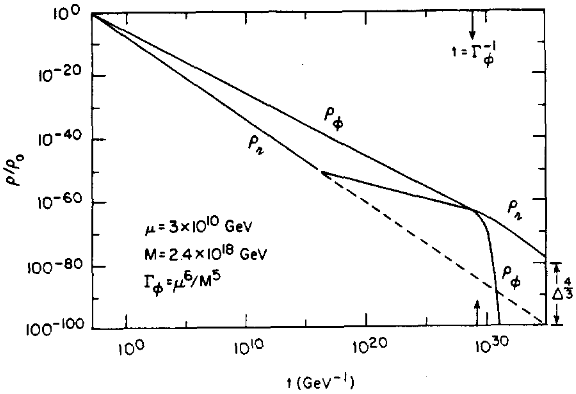

Fig. 1. This figure illustrates the transfer of the energy from the ϕ field to radiation as ϕ decays. The dashed line for ρ_ϕ would result if $\Gamma_\phi = 0$.

we give the results of a numerical evaluation of Δ, and the result is in excellent agreement with the above approximations. However, this result is unacceptable in the usual baryon number generation scenarios as the final temperature is too low to produce baryon number, and any pre-existing baryon number will be diluted by Δ.

(ii) $l_M = M^{-1}$: In this case the ϕ^2 term dominates in the effective potential for these modes. The initial condition for ϕ in this case is given by $\phi_i \simeq 0$, $\dot{\phi}_i \simeq 0$ [‡1]. The ϕ decay rate is also suppressed by a factor $m_\phi/(R^{-2}l_M^{-2} + m_\phi^2)^{1/2}$. The ϕ field remains at the origin until $R_M/l_M R = R^{-1}$ is of order m_ϕ. Recall that the entropy $\sim(RT)^3 \simeq 1$ remains constant until an inflationary epoch. We will assume that there is an inflationary epoch caused by some other sector of the theory. This is necessary to increase $(RT)^3$ to its present value of 10^{88}. We assume that the additional inflation commences at some temperature T^* ($\mu \leqslant T^* \leqslant M$) with vacuum energy T^{*4}. (Note we do not consider $T^* < \mu$ since in this case the baryon density of the universe will be diluted without the possibility for regeneration.)

The end state of this epoch is given by the final temperature T_f of order T^* and final scale size $R_f > 10^{28}R^*$. (This generates enough entropy for the inflationary scenario to work [5].) The timescale Δt for this inflationary epoch thus satisfies

$$R_f/R^* = \exp(H^*\Delta t) \geqslant 10^{28}\,, \tag{27}$$

or $\Delta t \geqslant 65/H^*$.

We see that during the inflationary epoch the momentum factor becomes irrelevant. The damping factor is given by

$$H^* = (\dot{R}/R)_{\text{inflation}} \simeq T^{*2}/M\,. \tag{28}$$

For T^* of order $\mu(H^* \sim m_\phi)$ the relevant timescale for damping τ_d is given by $\tau_d \sim 1/H^*$ and in a time Δt the energy has damped by a factor

$$f = \exp(-2\Delta t/\tau_d) \simeq 10^{-56}\,. \tag{29}$$

For $T^* \gg \mu(H^* \gg m_\phi)$, however, we have

$$\tau \sim H^*/m_\phi^2 \gg 1/H^*\,. \tag{30}$$

Since $1/H^*$ is the typical timescale for inflation it does not appear that the ϕ field will lose sufficient energy in this case; even for $T^* \sim 10\mu$.

Even though f can be 10^{-56}, the field has relaxed to a state which is *not* the zero-temperature ground state. Indeed, the potential contains terms of the order $(T^{*4}\phi^2/M^2)$ due to the cosmological constant T^{*4}. These terms, however, are ignorant about the location of the ground state.

In conclusion, we have discussed in this letter a cosmological problem for the Polonyi potential. Although we have focused on this model in particular, the problem seems generic to any weakly coupled scalar field with intermediate scale energy density in the potential. We have found a possible solution to the problem in the case where $l_M \simeq 10^{28}/M$ by having on O'Raifeartaigh-type potential in the hidden sector. In this case, thermodynamic considerations force the fields to be at a minimum near the origin. (This has also been studied by M. Dine and D. Nameshanski.) This solution, however, seems to fail when one considers the case $l_M \simeq M^{-1}$. We will report these results separately.

One of us (EWK) would like to thank Steve Barr and David Reiss for conversations about the Polonyi model.

[‡1] In this case, only one degree of freedom $[\phi_k(t) \equiv \int d^3x \exp(i\boldsymbol{k}\cdot\boldsymbol{x})\,\phi(x,t)$ with $k = l_M^{-1}]$ is relevant. Since $T \sim k$, the initial conditions are determined by quantum mechanical (rather than thermodynamic) considerations alone. We have the ground state expectations $\langle\phi_k\rangle \sim 0$, $\langle\dot{\phi}_k\rangle \sim 0$, with large quantum fluctuations. The ϕ field begins to roll down the potential when R has reached a size $R^{-1} \sim \sim \mu^2/M$.

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