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# Quintessence, inflation and baryogenesis from a single pseudo-Nambu-Goldstone boson

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**ABSTRACT:** We exhibit a model in which a single pseudo-Nambu-Goldstone boson explains dark energy, inflation and baryogenesis. The model predicts correlated signals in future collider experiments, WIMP searches, proton decay experiments, dark energy probes, and the PLANCK satellite CMB measurements.

**KEYWORDS:** cosmology, baryogenesis, inflation, quintessence, dark energy.

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## 1. Introduction

The most plausible candidate for a quintessence explanation of dark energy is a pseudo-Nambu-Goldstone boson (PNGB) of a spontaneously and explicitly broken global  $U(1)$  symmetry [1][2]. Denoting the scale of the explicit breaking by  $M$  and the scale of the spontaneous breaking by  $f$ , the effective mass of the PNGB is  $m_{\text{eff}} \sim M^2/f$ . The required value of this mass for a successful quintessence model is of order the Hubble parameter  $H_0 \sim 10^{-33}$  eV; this can be obtained naturally if the spontaneous breaking scale  $f$  is very large, roughly comparable to the Planck scale  $M_{\text{Planck}} \simeq 10^{19}$  GeV. Because of the symmetry, the PNGB has only derivative couplings to matter and radiation, plus couplings whose dimensionless strength is suppressed by powers of  $M/f \sim 10^{-22}$ . This evades a number of strong experimental and observational constraints on weakly-coupled ultralight scalars [3].

In a more general class of PNGB quintessence models, the effective mass of the PNGB will vary over time, *i.e.* it will be a function of the scale factor  $a$  obtained by solving the coupled cosmological equations of motion. Since  $f$  is comparable to the Planck scale, it is possible that  $m_{\text{eff}}$  was much larger at early times, without losing the key property that  $m_{\text{eff}}/f \ll 1$ . This raises the possibility that the quintessence PNGB field may also have been responsible for primordial inflation, albeit in some modification of the usual scenario.

A simple avenue towards quintessential inflation is then to assume that  $m_{\text{eff}}$  and the PNGB scalar potential  $V$  scale (at least roughly) like some power of the Hubble rate  $H$ . Any model with  $m_{\text{eff}} \sim H$  has the additional virtue of removing the coincidence problem, *i.e.* the fact that the ratio of scales  $M/f$  is of order  $H_0$  is no longer a coincidence, but rather has some dynamical origin.

Any model of quintessential inflation must explain why the energy density of the universe was dominated by radiation at the time of Big Bang Nucleosynthesis (BBN), even though the inflaton dominates the energy density now and dominated it in a primordial epoch as well. For a PNGB, the simplest explanation is that the PNGB decays to matter and radiation via derivative couplings of the form

$$\lambda_{ij} \frac{f}{M_{\text{Planck}}} g^{\mu\nu} \partial_\mu \theta \bar{\psi}_i \gamma_\nu \psi_j, \quad (1.1)$$

where  $\theta$  denotes the PNGB field rescaled by  $f$  to make it dimensionless, and the  $\lambda_{ij}$  are moderately small dimensionless couplings. For  $f \sim M_{\text{Planck}}$ , this will allow large entropy production from PNGB decays. Since the PNGB potential varies over time, the equation of state of the PNGB also varies. Thus it is natural to have periods of inflation interspersed with periods of radiation dominance.

The matter coupling (1.1) was introduced by Cohen and Kaplan in their thermodynamic model for baryogenesis [4]. In their scenario  $B$  or  $L$  violating processes occur via dimension six four-fermion operators suppressed by a relatively low scale  $\Lambda \sim 10^8$  GeV. Combined with a PNGB possessing dimension five couplings like (1.1), they generate a baryon asymmetry in thermal equilibrium (at temperatures  $\gtrsim 10^8$  GeV), and a further asymmetry at lower temperatures from PNGB decays.

In this paper we construct a natural model of PNGB quintessential inflation that also implements the Cohen-Kaplan mechanism for baryogenesis. At the same time that we provide quintessence and inflation, our model gives a simpler explanation of the baryon asymmetry than the original scenario of [4]. We are able to assume that the scale for the dimension six  $B$ ,  $L$  violating operators is of order  $10^{15}$  GeV, or of order  $10^{11}$  GeV for purely  $L$  violating operators. Thus we are slightly above or saturating the current experimental bounds from proton decay [5]. Our cosmological evolution begins with generic initial conditions, unlike the original models of [1][2], at an initial temperature  $T_0 \sim 10^{17}$  GeV, which is comfortably less than  $f \sim M_{\text{Planck}}$ . Our model requires no unnatural tunings other than that of the cosmological constant, a tuning that is unavoidable since quintessence does not solve the cosmological constant problem.

## 2. FRW cosmology driven by a Nambu-Goldstone boson

Consider a theory with a  $U(1)$  global symmetry under which some complex scalar field  $\Phi(x, t)$  transforms as

$$\Phi \rightarrow e^{i\alpha} \Phi, \quad (2.1)$$

where  $\alpha$  is a constant. Other fields, including fermions, may also transform nontrivially; in particular the symmetry may be a chiral symmetry. We imagine that this symmetry is spontaneously broken at some high scale, near  $M_{\text{Planck}}$ , determined by the vev  $f$  of the scalar. Expanding around this vev gives

$$\Phi = (f + \sigma(x, t)) e^{i\theta(x, t)}, \quad (2.2)$$

where  $\sigma$  is a real scalar and  $\theta$  is another real scalar which has been rescaled by  $1/f$  to be dimensionless. The effective theory at lower energies has the original symmetry nonlinearly realized, with the Nambu-Goldstone boson (NGB)  $\theta$  undergoing a shift:

$$\theta(x, t) \rightarrow \theta(x, t) + \alpha . \quad (2.3)$$

Obviously the NGB only has derivative couplings in this effective theory. Higher dimension operators are suppressed by powers of a high scale  $f$  which we are roughly equating to  $M_{\text{Planck}}$ . To leading order in  $f$  the NGB action is just

$$\int d^4x \sqrt{-g} \frac{1}{2} f^2 g^{\mu\nu} \partial_\mu \theta \partial_\nu \theta . \quad (2.4)$$

Let us further suppose that the energy density of the universe during some epoch is dominated by the NGB. Then the Friedmann equation is approximately

$$H^2 = \frac{f^2}{3k^2} \dot{\theta}^2 , \quad (2.5)$$

where  $k^2 = M_{\text{Planck}}^2/4\pi$ . The equation of motion for  $\theta$  (again to leading order in  $f$ ) is just

$$\ddot{\theta} + 3H\dot{\theta} = 0 , \quad (2.6)$$

while the continuity equation is given by

$$\dot{\rho}_\theta = -6H\rho_\theta . \quad (2.7)$$

The cosmological solutions for the spatially averaged vacuum expectation value of  $\theta(t)$  depend upon the initial conditions. If we assume that  $\dot{\theta}$  is initially zero, with  $\theta$  taking an arbitrary initial value, then the solution to (2.5)-(2.7) is a static universe. If instead we assume that both  $\dot{\theta}$  and  $\theta$  have arbitrary initial values, we get an expansion dominated (initially) by kination:

$$\theta(t) = -\ln a(t) , \quad H = \frac{1}{3t} , \quad a(t) = t^{1/3} , \quad (2.8)$$

where we have taken  $f = \sqrt{3}k$  to simplify notation. Of course in this solution  $\rho_\theta$  dilutes like  $1/a^6$ , as appropriate for kination, *i.e.*, an equation of state dominated by kinetic energy. Derivative couplings of  $\theta$  to ordinary matter will allow the  $\theta$  vacuum energy to be converted to a thermal radiation bath via decays. Since ordinary matter and radiation dilute like  $1/a^3$  and  $1/a^4$ , they will eventually dominate the expansion.

Let us interpret this solution in terms of the original global symmetry. The equation of motion for  $\theta$  is just the statement that the global current is covariantly conserved; the corresponding conserved global charge of the vacuum  $Q$  is proportional to  $a^3\dot{\theta}$ , which from (2.8) is indeed seen to be a constant. Once we include matter couplings the vacuum charge  $Q$  is no longer constant, but the total global charge including matter contributions is.

### 3. FRW cosmology driven by a pseudo-Nambu-Goldstone boson

We can modify the discussion in the previous section by introducing a nonvanishing potential for the  $\theta$  field and a noncanonical kinetic function:

$$\int d^4x \sqrt{-g} \left[ \frac{1}{2} F(\cos \theta) g^{\mu\nu} \partial_\mu \theta \partial_\nu \theta - V(\cos \theta) \right]. \quad (3.1)$$

$F$  and  $V$  explicitly break the global  $U(1)$  symmetry down to a discrete periodic remnant:

$$\theta \rightarrow \theta + 2\pi N, \quad (3.2)$$

where  $N$  is any integer. This kind of explicit breaking would arise if terms proportional to powers of  $\Phi + \Phi^*$  were present in the original action. Alternatively, there could be Yukawa couplings or derivative couplings of  $\Phi$  to fermions:

$$\lambda_{ij} \Phi \bar{\psi}_i \psi_j, \quad \lambda'_{ij} g^{\mu\nu} \partial_\mu \Phi \bar{\psi}_i \gamma_\nu \psi_j \quad (3.3)$$

In this case chiral symmetry breaking will induce, at the loop level, an effective action of the form (3.1). The simplest possibility, considered in [1, 2], gives

$$\int d^4x \sqrt{-g} \left[ \frac{1}{2} f^2 g^{\mu\nu} \partial_\mu \theta \partial_\nu \theta - M^4 (1 - \cos \theta) \right], \quad (3.4)$$

where  $M$  is some chiral symmetry breaking scale much smaller than  $f$ .

Because of this explicit breaking,  $\theta$  is now a pseudo-Nambu-Goldstone boson, with a mass of order  $M^2/f$ . It is technically natural for this mass to be small, since it vanishes in a symmetry limit.

Now consider FRW cosmology driven by such a PNGB. The PNGB equation of motion is:

$$\ddot{\theta} + 3H\dot{\theta} + \frac{M^4}{f^2} \sin \theta = 0. \quad (3.5)$$

The cosmological solutions again depend upon the initial conditions. One possibility, considered in [1, 2], is that  $\dot{\theta}$  is initially zero, with  $\theta$  taking an arbitrary initial value. Then the equation of motion (3.5) is heavily overdamped, and  $\theta$  remains approximately constant until  $H(t)$  decreases to the point where  $H \sim M^2/f$ . Thus the PNGB behaves approximately like dark energy. The coincidence problem is not solved unless one finds a rationale for why  $M^2/f$ , a ratio of two seemingly independent scales, is roughly equal to the Hubble rate today. Even then one also needs to explain why the initial value of  $\dot{\theta}$  is much less than  $H$ .

Another possibility is that both  $\dot{\theta}$  and  $\theta$  have generic initial values during some epoch where the PNGB dominates the expansion. Then for  $M/f \ll 1$  the cosmological solution is an oscillatory perturbation of the NGB solution (2.8). The global charge of the vacuum  $Q \propto a^3 \dot{\theta}$  is no longer conserved; in fact it increases with time like

$$\frac{\dot{Q}}{Q} \sim t. \quad (3.6)$$

This case is more generic than the first one, but does not provide an explanation of dark energy.

#### 4. From generic PNGBs to Slinky

The most general effective action for the PNGB invariant under the shift (3.2) is

$$\int d^4x \sqrt{-g} \left[ \frac{1}{2} F(\cos \theta) P(X) - V(\cos \theta) + \mathcal{L}_m \right], \quad (4.1)$$

where  $X = g^{\mu\nu} \partial_\mu \theta \partial_\nu \theta$ ,  $F$  and  $V$  are arbitrary functions, and  $P$  is an arbitrary polynomial. The matter lagrangian  $\mathcal{L}$  could contain both derivative couplings to  $\theta$  and couplings to functions of  $\cos \theta$ . If there are further explicit breakings of the global symmetry, the action may also contain a nonperiodic dependence on  $\theta$ .

The Slinky model of quintessential inflation introduced in [6]-[8] is a PNGB quintessence model of the general class just described, but with some special features. It does not require tuned initial conditions; in particular the initial value of  $\dot{\theta}$  is of order  $H$ . It is the simplest model with a periodic equation of state parameter for quintessence. The quintessence energy density dominates the Friedmann equation during an earlier epoch as well as during the present epoch, causing primordial inflation as well as present-day acceleration. In order for this to happen, both the potential and the kinetic energy must be proportional to the square of the Hubble parameter times periodic functions:

$$V \propto H^2, \quad \dot{\theta}^2 \propto H^2. \quad (4.2)$$

The first relation implies that these models will have a PNGB mass that decreases with time.

The explicit form of the Slinky is easily derived. We begin with the general action (4.1) and make the simplest nontrivial choice  $P(X) = X$ . Without loss of generality, we will take  $\theta(t) = 1$  today. For quintessence the equation of state parameter  $w(t)$  should be close to  $-1$  today; for simplicity we will take it to be exactly  $-1$ . Since

$$1 + w(t) = F(\cos \theta) \dot{\theta}^2. \quad (4.3)$$

This means that  $F(\cos \theta) \dot{\theta}^2 \rightarrow 0$  now. The simplest nonvanishing function that does this is

$$F = \frac{3k^2}{b^2} (1 - \cos \theta), \quad (4.4)$$

where  $b$  is a dimensionless constant given by

$$b = \sqrt{\frac{3}{4\pi}} \frac{M_{\text{Planck}}}{f}. \quad (4.5)$$

The explicit breaking represented by the cosine term in (4.4) does not have any tunable small parameter associated with it. Thus one might worry that this breaking is not small. However since this breaking occurs in the kinetic function, it generates global charge nonconserving processes that are suppressed by powers of momenta divided by  $k$ . Thus for dynamics well below the Planck scale this breaking is indeed small.

The PNGB potential  $V(\theta)$  must vanish when we turn off the explicit breaking. Of course this requirement is a tuning, reflecting the fact that quintessence models do not explain why the  $\theta$ -independent cosmological constant vanishes. In this limit we are kination dominated, with  $w(t) = -\cos\theta = 1$ . Thus the simplest form for the potential is

$$V \propto (1 + \cos\theta)H^2(\theta) . \quad (4.6)$$

We need to solve three equations of motion, beginning with the Friedmann equation:

$$H^2 = \frac{2}{3k^2}\rho_\theta . \quad (4.7)$$

The second equation is the scalar EOM:

$$0 = \ddot{\theta} + 3H\dot{\theta} + \frac{1}{2}\frac{F'}{F}\dot{\theta}^2 + \frac{V'}{F} , \quad (4.8)$$

where a prime denotes a derivative with respect to  $\theta$ . The third equation is the continuity equation:

$$\dot{\rho}_\theta = -3(1+w)H\rho_\theta , \quad (4.9)$$

where  $w(\theta)$  is the equation of state parameter for quintessence.

Plugging in (4.4) and (4.6) we get a solution with

$$H = H_0 e^{\frac{3}{2b}(\theta - \sin\theta)} \quad (4.10)$$

where  $H_0$  is the Hubble parameter today.

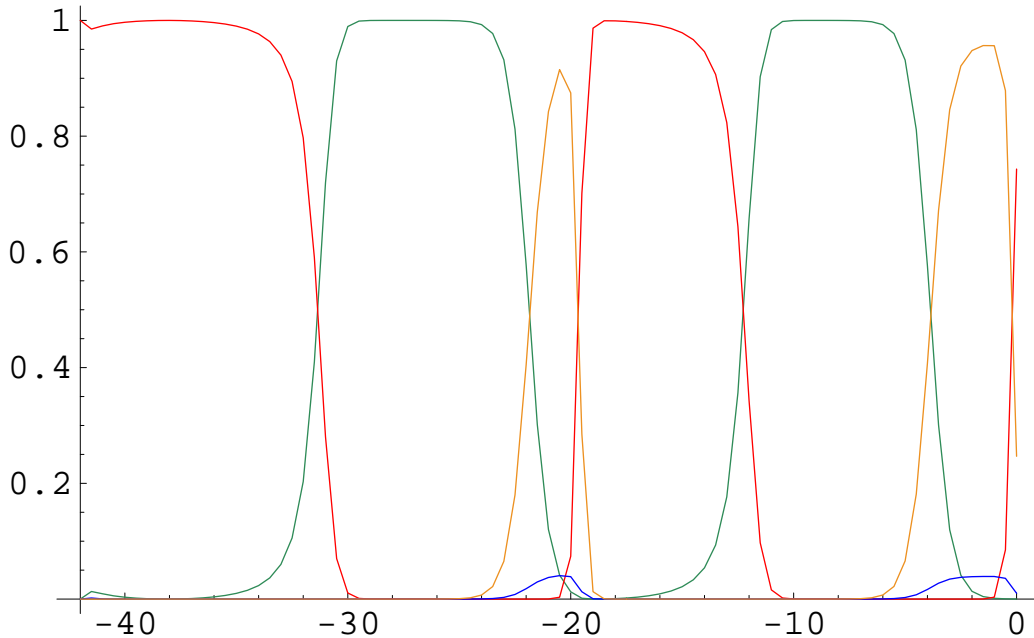
Now we have enough information to specify  $F$  and  $V$  completely:

$$\begin{aligned} F(\theta) &= \frac{6k^2}{b^2}(1 - \cos\theta) , \\ V(\theta) &= \frac{1}{2}\rho_0(1 + \cos\theta) \exp\left[\frac{3}{b}(\theta - \sin\theta)\right] , \end{aligned} \quad (4.11)$$

where  $\rho_0 = 3k^2H_0^2/2$  is the analog of  $M^4$  in the standard PNGB quintessence model discussed earlier. Note that the scalar potential has an additional explicit breaking that does not respect the periodicity (3.2). This breaking is small during any epoch such that  $H/k \ll 1$ .

Slinky is a PNGB quintessence model characterized by a periodic equation of state  $w = -\cos\theta$ . The effective mass of the PNGB is tied to the expansion rate, and therefore decreases over time. These two features combined allow the same PNGB to serve the role both of the inflaton in the early universe and the quintessence field that drives the accelerated expansion in the current epoch.

The number of inflationary epochs that have occurred is controlled by the dimensionless parameter  $b$ ; larger values of  $b$  mean more inflationary periods. Most values of  $b$  are ruled out by the requirement that the universe had at most a tiny inflationary component during BBN [9]. Note from (4.4) that  $b \lesssim 0.49$  implies that  $f > M_{\text{Planck}}$ . The Slinky models



**Figure 1:** The cosmological history of the simplest Slinky model. The x-axis is the logarithm base 10 of the scale factor. Shown are energy densities of radiation (green), dark matter (orange), baryons (blue) and dark energy (red), as a fraction of the critical density.

considered in [7] have values in the range  $0.09 < b < 0.4$ . This does not necessarily mean that these models require trans-Planckian vevs; for example, if we replace the single PNCB by  $N$  identical ones, then  $b$  is effectively rescaled as  $b \rightarrow b/\sqrt{N}$ . Thus small values for  $b$  do not indicate a scenario that is technically out of control.

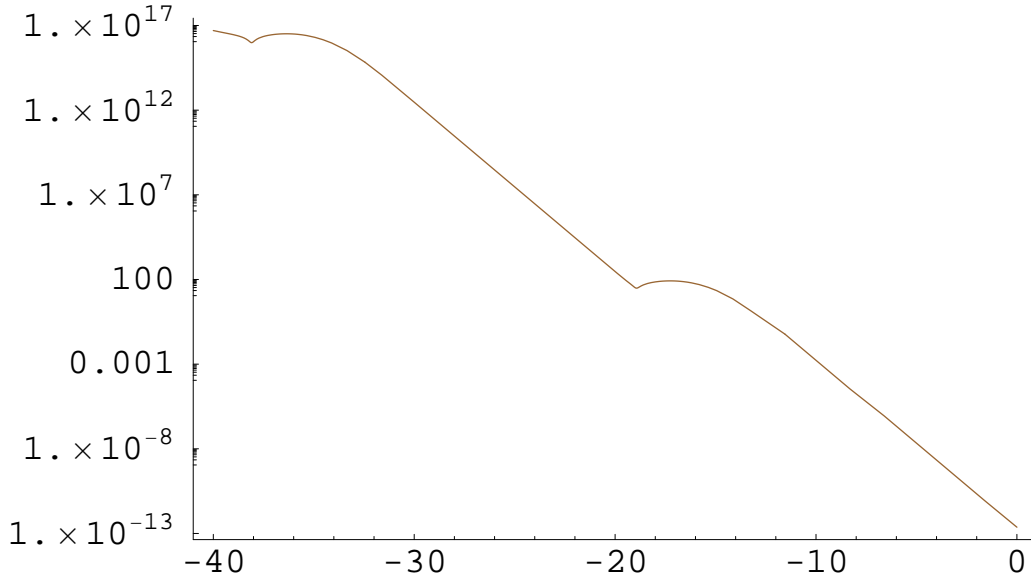
The nonstandard cosmological history of this model is depicted in Figure 1. The energy densities of radiation, dark matter, baryons and dark energy, computed as a fraction of the critical density, are plotted as function of the logarithm of the scale factor. We have chosen to start the cosmological evolution at a scale factor of  $10^{-42}$ , but this choice is not essential. We have assumed that the Friedmann equation is initially dominated by the PNCB. For simplicity we set the initial radiation density to zero, so all radiation arises from PNCB decays, as explained in the next section.

The radiation temperature as a function of the scale factor is shown in Figure 2. During the initial inflation the temperature is roughly constant at around  $10^{16}$  GeV, increasing mildly towards the end of this first inflationary epoch. This behavior is in contrast to standard primordial inflation, where the temperature first decreases very rapidly, then increases very rapidly (reheating). The new behavior is due to the time-varying equation of state.

## 5. Couplings to matter

The allowed couplings of  $\theta$  to matter can be divided into two classes:





**Figure 2:** The cosmological history of the simplest Slinky model. The x-axis is the logarithm base 10 of the scale factor. Shown are energy densities of radiation (green), dark matter (orange), baryons (blue) and dark energy (red), as a fraction of the critical density.

1. Couplings of functions of  $\cos \theta$  to matter. These couplings explicitly break the global  $U(1)$  symmetry. It is technically natural to take the dimensionless coupling constants for all such couplings to be small.
2. Derivative couplings of  $\theta$  to matter. Some of these couplings respect the full global  $U(1)$  symmetry, so there is no symmetry argument for tuning their dimensionless coupling constants to be small; indeed we will assume that they are of order one. However almost all such couplings are higher dimension operators suppressed by powers of momenta over powers of  $M_{\text{Planck}}$ .

Thus for Slinky, as well as for more generic PNGB quintessence models, it is a good approximation to only include couplings to matter of the form

$$\lambda_{ij} g^{\mu\nu} \partial_\mu \theta \bar{\psi}_i \gamma_\nu \psi_j, \quad (5.1)$$

where  $\lambda_{ij}$  is dimensionless. A caveat is that the Slinky potential (4.6) has an additional explicit breaking of the global  $U(1)$  symmetry, that we are assuming has no analog in the matter couplings.

Couplings of the form (5.1) allow the  $\theta$  field to decay into ordinary matter. This process can be modeled, albeit roughly [10], as an additional friction term in the  $\theta$  equation of motion:

$$0 = \ddot{\theta} + 3H\dot{\theta} + \Gamma\dot{\theta} + \frac{1}{2} \frac{F'}{F} \dot{\theta}^2 + \frac{V'}{F}. \quad (5.2)$$

Here  $\Gamma$  is the decay width, which for decays into pairs of fermions can be generically written

$$\Gamma = k_m m_{\text{eff}}(t) , \quad (5.3)$$

where  $k_m \sim \lambda_{ij}$  are dimensionless couplings which we will take to be approximately constant with magnitudes in the range .1 to .01;  $m_{\text{eff}} \propto H$  is the time varying mass of the PNGB, obtained by expanding the potential (4.6).

This process converts quintessence vacuum energy into matter and radiation. Thus Slinky produces the following nonstandard cosmological history:

- During a primordial epoch,  $w$  is close to -1 and inflation occurs.
- Eventually  $w$  increases towards +1, and the radiation produced from  $\theta$  decays, which is now diluting less rapidly than the quintessence energy, comes to dominate the expansion.
- The process repeats. In the original Slinky model [6], which we will use from now on, the second radiation dominated epoch overlaps with the time of Big Bang Nucleosynthesis.
- Well before BBN time,  $m_{\text{eff}}$  becomes so small that the PNGB decays to matter effectively turn themselves off, due to kinematic suppression. PNGB decays to photons continue, but these are loop-suppressed and so do not give large entropy production.
- Right now we are entering the third inflationary phase.

From the arguments given above, the only couplings in this model whose values affect the cosmological evolutions are  $b$  and the  $\lambda_{ij}$ . The value of  $b$  is coarsely adjusted such that we ensure that BBN time is radiation dominated and that  $w$  is close to  $-1$  today. The dominant  $\lambda_{ij}$  couplings are adjusted such that the radiation and matter fractions  $\Omega_r/\Omega_\Lambda$  and  $\Omega_{DM}/\Omega_\Lambda$  come out to their measured values today. For simplicity we assume that all of the dark matter is produced thermally, as *e.g.* in standard WIMP scenarios. The remaining matter couplings (for the moment) remain free.

We now see more clearly why the the temperature history represented in Figures 1 and 2 is so nonstandard. The coupling between matter and the PNGB field force them track each other, giving significant entropy production even after the first period of inflation . This entropy production turns off well before BBN time due to the kinematic suppression of PNGB decays.

## 6. Baryogenesis

If baryon number  $B$  and lepton number  $L$  are exactly conserved, then  $B$  or  $L$  violating decays of the PNGB simply do not occur. This can be seen from the coupling (5.1), which with integration by parts vanishes for any conserved current  $j_\mu = \bar{\psi}\gamma_\mu\psi$ .

However we certainly do not expect the global symmetries  $B$  and  $L$  to be respected by unification scale physics. We therefore introduce all possible  $B$  and  $L$  violating dimension six operators:

$$\frac{\psi_i \psi_j \psi_k \psi_l}{\Lambda_{ijkl}^2} \quad (6.1)$$

suppressed by some superheavy scales  $\Lambda_{ijkl}$ . Suppressing flavor labels, this implies that the  $B$  or  $L$  violating decay width of the PNGB can be written

$$\Gamma = \lambda^2 \frac{m_{\text{eff}}^5}{\Lambda^4} . \quad (6.2)$$

Thus PNGB evolution and decay during the first inflationary epoch will produce a net  $B$  and  $L$  asymmetry, including a  $B - L$  asymmetry that will survive the  $B + L$  washout by sphalerons during the electroweak phase transition.

At sufficiently high temperatures the  $B$  or  $L$  violating processes will be in thermal equilibrium. The equilibrium condition is [4]

$$\Gamma \gg \frac{\ddot{\theta}}{\dot{\theta}} = \frac{3}{2}(1 - \cos \theta) H_{PNGB} , \quad (6.3)$$

where  $H_{PNGB}$  denotes the Hubble rate that would result from the PNGB alone, given in equation (4.10). Figure (3) shows an estimate of the value of the scale factor at which  $B$  or  $L$  violating processes decouple, assuming  $\Lambda = 10^{15}$  GeV. The corresponding decoupling temperature is about  $8 \times 10^{15}$  GeV.

Once we are out of equilibrium the baryon asymmetry continues to grow via PNGB decays. The net baryon number density as a function of time can be written as [4]

$$\dot{n}_B = \lambda \frac{m_{\text{eff}}^5}{\Lambda^4} \frac{12k^2}{b^2} \dot{\theta} \sin \theta . \quad (6.4)$$

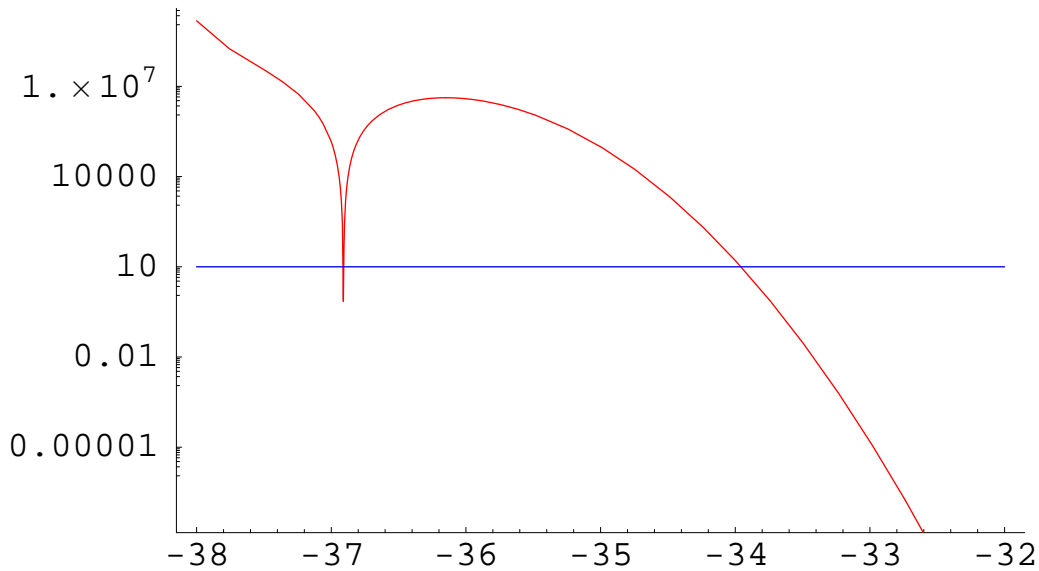
It is useful to express the baryon number density as a function of  $\theta$ , *i.e.*

$$n'_B = \lambda \frac{m_{\text{eff}}^5}{\Lambda^4} \frac{12k^2}{b^2} \sin \theta . \quad (6.5)$$

To calculate the baryon asymmetry produced, we have to include the expansion of the universe. This can be easily done by replacing all the volume factors in the equations above by the corresponding comoving volumes, proportional to  $a^3$ ,

$$\frac{d}{d\theta} (a^3 n_B) = \lambda a^3 \frac{m_{\text{eff}}^5}{\Lambda^4} \frac{12k^2}{b^2} \sin \theta . \quad (6.6)$$

To obtain the net baryon number density we have to numerically integrate this expression from the decoupling temperature to  $\theta_f$ , the temperature where baryon production from PNGB decays is cut off kinematically by the decrease in  $m_{\text{eff}}$ . For decays to baryons  $\theta_f$  corresponds to a scale factor of approximately  $10^{-27}$ , when  $m_{\text{eff}} \sim H(\theta) \simeq 1$  GeV.



**Figure 3:** The ratio of the relative rates of  $B$  or  $L$  violating processes versus the cosmological evolution of the PNGB field, as a function of the logarithm base 10 of the scale factor. We have taken  $\Lambda = 10^{15}$  GeV. The ratio is shown in red;  $B$  or  $L$  violation goes out of equilibrium below the blue line.

From the baryon density we want to extract the baryon to photon ratio,  $\eta$ . As usual the photon number density is given by

$$n_\gamma = \frac{2\zeta(3)}{\pi^2} T^3 . \quad (6.7)$$

Contrary to the standard scenarios [11, 12], where the baryon (more precisely,  $B-L$ ) to photon ratio does not change after baryogenesis, in our model there is significant entropy production due to the coupling of the PNGB field to radiation. The net effect of such a production will be to dilute any baryon to photon ratio produced before BBN, where entropy production stops. Therefore  $\eta$  can still be calculated at  $\theta_f$ , but we have to include the extra dilution factor  $\gamma$ , given by

$$\gamma = \frac{S_{\theta_f}}{S_{\theta_{BBN}}} . \quad (6.8)$$

Here  $S = g_* a^3 T^3$ , with  $g_*$  the effective number of relativistic degrees of freedom. Alternatively, we can calculate  $\eta$  directly at BBN time.

The only input parameters left to adjust in our simple model are the scale  $\Lambda$  and the dimensionless coupling  $\lambda$ . We consider two well-motivated scenarios. In the first scenario,  $\Lambda \simeq 10^{15}$  GeV, just about saturating the generic lower bounds from the nonobservation of proton decay [5]. We estimate the baryon to photon ratio  $\eta$  in this case to be

$$\eta \simeq 4\lambda \times 10^{-10} . \quad (6.9)$$

This gives the observed baryon asymmetry for  $\lambda \simeq 1$ .

The second scenario has  $\Lambda \simeq 10^{11}$  GeV, as might be appropriate in a leptogenesis model. Here we estimate

$$\eta \simeq 5\lambda \times 10^{-8}, \quad (6.10)$$

which gives the observed baryon asymmetry for  $\lambda \simeq 0.01$ .

The cosmological history of  $\eta$  is unlike that of previously discussed models of quintessential baryogenesis [13]-[20]. At early times the PNGB effective mass is large, and the baryon number violating processes are least suppressed. A baryon or lepton asymmetry is produced in thermal equilibrium, but this contribution is negligible compared to the asymmetry produced later by PNGB decays. This is basically because the PNGB decay process is enhanced by  $f^2$ , the square of the Planckian scalar vev, as seen in (6.4). The net baryon production turns itself off kinematically as  $m_{eff}$  decreases. The baryon to photon ratio drops dramatically during the subsequent second inflation era, where there is further large entropy production. The fact that the baryon to photon number density observed today is very small is thus due entirely to the existence of a second inflation era;  $B$  or  $L$  violation was only mildly suppressed at the time that most of the net baryon or lepton excess was created.

## 7. Conclusions

We have shown that a single noncanonical PNGB could be responsible for both primordial inflation and the present day accelerated expansion, while simultaneously generating the observed baryon excess. The baryon asymmetry is generated after the first inflationary period, from PNGB decays via dimension six operators that violate  $B$  and/or  $L$ . The ratio of baryon to photon number densities is greatly diluted later on, via entropy production from  $B$  and  $L$  conserving PNGB decays. The baryon energy density is completely negligible except during two eras: the present day and an era around the time of the electroweak phase transition.

While economical, technically natural, and consistent with current data, our model has some theoretical shortcomings. It explains dark energy but does not solve the cosmological constant problem. There is also no first principles explanation for why the PNGB potential scales like the square of the Hubble rate. Resolving these shortcomings presumably involves an ultraviolet completion of the PNGB effective theory into a more fundamental framework.

One of the distinguishing features of this model, is that it predicts substantial entropy production in the era between the electroweak phase transition and BBN. If dark matter is predominately composed of thermally produced WIMPs, this prediction can be tested by combining collider data with signals from direct and indirect WIMP searches [8], [21]-[25].

Because of the entropy production after the second inflation era (but before BBN), any pre-existing baryon asymmetry is much diluted. This requires that the operators responsible for  $B$  violation be not too much suppressed, forcing us close to the current experimental bounds for proton decay. Thus another prediction is that proton decay will be observed in one of the future proposed experiments [26]-[29].

Future dark energy probes will pin down the equation of state of dark energy with much greater precision. In our simple model we artificially set  $w = 1$  exactly at redshift  $z = 0$ ; more relevant is that  $w$  varies by about 2% as redshift is varied between 0 and 2. This variation is comparable to the one-sigma projected combined errors after the Stage IV dark energy probes [30].

Models of the type discussed here predict running of the spectral indices of the Cosmic Microwave Background (CMB) [31]. These effects may be large enough to extract from observations of the PLANCK satellite [32]. The smoking gun of a predictive model of quintessential inflation is that these CMB effects, a result of inflation, are directly related to the detailed equation of state of dark energy.

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