Quintessential Baryogenesis

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Abstract

The simplest explanation for early time acceleration (inflation) and the late time acceleration indicated by recent data is that they have a common origin. We investigate another generic cosmological implication of this possibility, that the baryon asymmetry of the universe may be generated in such models. We identify several novel features of baryogenesis in such a universe, in which a rolling scalar field is always part of the cosmological energy budget. We also propose a concrete mechanism by which the baryon asymmetry of the universe may be generated in this context. We analyze the generic properties of and constraints on these cosmologies, and then demonstrate explicitly how a complete cosmology may develop in some specific classes of models.

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I. INTRODUCTION

Rolling scalar fields are a mainstay of modern cosmology. This is perhaps best-illustrated by the inflationary paradigm \[1,2,3\], in which most implementations involve a scalar field rolling towards the minimum of its potential in such a way that the potential energy of the field is the dominant component of the energy density of the universe. There are, however, many other cosmological instances in which scalar fields are invoked.

During the last few years a new consistent picture of the energy budget of the universe has emerged. Large scale structure studies show that matter (both luminous and dark) contributes a fraction of about 0.3 of the critical density, while the position of the first acoustic peak of the cosmic microwave background power spectrum indicates that the total energy density is consistent with criticality. The discrepancy between these two measurements may be reconciled by invoking a negative pressure component which is termed *dark energy*. While there are a number of different observational tools to study dark energy – number counts of galaxies \[4\] and galaxy clusters \[5\] for example – the most direct evidence to date comes from the light-curve measurements of intermediate redshift type IA supernovae \[6,7\]. Consistency between these observations and others such as weak gravitational lensing \[8\] and large scale structure surveys \[9\] implies that the dark energy \(X\) satisfy \(\Omega_X \sim 0.7\) and that the equation of state be \[10,11\]

\[
w_X \equiv \frac{p_X}{\rho_X} \leq -0.6 ,
\]

leading to the acceleration of the universe.

It is of course possible that this mystery component is a cosmological constant \(\Lambda\), for which \(w_\Lambda = -1\). However, understanding the nature of such an unnaturally small \(\Lambda\) is at least as difficult as understanding one that is zero. Alternatively, it has been suggested \[12-14\] that if the cosmological constant itself is zero, the dark energy component could be due to the dynamics of a rolling scalar field, in a form of late-universe inflation that has become known as *quintessence*. Although there are a number of fine-tuning problems associated with this idea, it does provide a way to ensure the late-time acceleration of the universe, albeit at the expense of introducing a second (after the inflaton) cosmologically relevant rolling scalar field. While not addressing the cosmological constant problem itself, and suffering from fine-tuning, quintessence itself has the advantage of avoiding a future horizon in space-time,
and hence makes consistency with what is known about perturbative string theory more likely.

It is natural to wonder whether the inflaton and the quintessence field might be one and the same [15]. In fact, specific models for this have been proposed [15]-[17]. Clearly such models are attractive because we need only postulate a single rolling scalar, but may be problematic either theoretically or phenomenologically.

In this paper we investigate how we may further limit the proliferation of rolling scalar fields required in modern cosmology by studying how the scalar field responsible for late-time acceleration of the universe might also solve another outstanding cosmological puzzle. Specifically we will be interested in the role that such a field may play in the generation of the baryon asymmetry of the universe. The spectacular success of primordial nucleosynthesis requires that there exist an asymmetry between baryons and antibaryons in the universe at temperatures lower than an MeV. This is quantified by the requirement

$$4 \times 10^{-10} \leq \eta \equiv \frac{n_B}{s} \leq 7 \times 10^{-10},$$

where $n_B \equiv n_b - n_{\bar{b}}$, with $n_{b(\bar{b})}$ the number density of (anti)baryons and $s$ is the entropy density. To generate such an asymmetry, the underlying particle physics theory must satisfy three necessary conditions – the Sakharov conditions [18]. These are baryon number $B$ violation, the violation of the discrete symmetries $C$ and $CP$ and a departure from thermal equilibrium, this last condition resulting from an application of the $CPT$ theorem. In this paper we are interested in how these conditions may be met within the context of dark energy models.

The relationship between early-time acceleration – inflation – and baryogenesis has been explored in some detail (for example see [19]-[33]). Here we investigate the opposite regime, that the quintessence field may be associated with the generation of the baryon asymmetry. Naturally, it would be particularly efficient if a single scalar field could be responsible for three fundamental phenomena in cosmology – inflation, baryogenesis and dark energy, and indeed we will show that baryogenesis occurs quite generically in models in which a single scalar is responsible for the two periods of cosmic acceleration.

The outline of this paper is as follows. In section II we will review some details about quintessence and explain how inflation and quintessence may be unified by generalizing the quintessential inflation model of Peebles and Vilenkin [16]. In section III we will describe
how quintessence and quintessential inflation may naturally yield a baryon asymmetry without the introduction of any new fields into the theory. We term this model *quintessential baryogenesis*, borrowing the phrasing from Peebles and Vilenkin. This turns out to depend to some extent on the details of quintessential inflation. In section IV we will discuss experimental and astrophysical constraints on our models and comment on how we may test the physics involved. We offer our comments and conclusions in the final section of the paper.

II. QUINTESSENCE AND GENERALIZED QUINTESSENTIAL INFLATION

As we have mentioned already, one approach to the dark energy problem is to assume that there is some as yet unknown process that sets the vacuum cosmological constant of the universe to zero, but that there exists a cosmologically-relevant scalar field in the universe, that has yet to reach its global minimum and therefore contributes an effective vacuum energy to the total. This idea has been termed quintessence and we shall briefly review it here.

The Einstein equations in cosmology may be written as

\[
\left(\frac{\dot{a}}{a}\right)^2 = \frac{8\pi}{3M_p^2} \rho, \tag{3}
\]

\[
\frac{\ddot{a}}{a} = -\frac{4\pi}{3M_p^2} (\rho + 3p), \tag{4}
\]

where we are using the Friedmann, Robertson-Walker (FRW) ansatz for the metric

\[
d s^2 = -d t^2 + a(t)^2 \left[ d r^2 + r^2 d\Omega_2^2 \right]. \tag{5}
\]

Here the energy density \( \rho \) and pressure \( p \) for a real homogeneous scalar field \( \phi \) are given by

\[
\rho_\phi = \frac{1}{2} \dot{\phi}^2 + V(\phi), \tag{6}
\]

\[
p_\phi = \frac{1}{2} \dot{\phi}^2 - V(\phi), \tag{7}
\]

respectively, with \( V(\phi) \) the potential, and where we have defined the Planck mass by \( G \equiv M_p^{-2} \). The scalar field itself obeys

\[
\ddot{\phi} + 3 \left(\frac{\dot{a}}{a}\right) \dot{\phi} + \frac{dV(\phi)}{d\phi} = 0, \tag{8}
\]

with a prime denoting a derivative with respect to \( \phi \).
Now, to explain the current data indicating an accelerating universe, it is necessary to have the dominant type of matter at late times be such that $\dot{a} > 0$. If this matter is to be $\phi$, then (4) implies $\rho_\phi + 3p_\phi < 0$ which, since we conventionally write $p_\phi \equiv w_\phi \rho_\phi$ translates into an equation of state parameter that obeys

$$w_\phi < -\frac{1}{3}. \quad (9)$$

With an appropriate choice of potential the resulting cosmic acceleration can be arranged to occur at late times, and provides an explanation for the supernova data [3, 4].

Let us now consider extending these ideas to incorporate inflation [15]-[17]. This further constrains our potential, and we will present a general analysis, indicating where appropriate how the particular results of [16] are recovered.

Consider a generic potential which, for convenience and for comparison with other work, we will express in the form

$$V(\phi) = \begin{cases} V_1(\phi), & \phi \in (-\infty, 0] \\ V_2(\phi), & \phi \in (0, \infty) \end{cases} \quad (10)$$

such that $V_1(\phi) \geq 0$, $V_2(\phi) \geq 0 \ \forall \phi$, $V_1(0) = V_2(0)$ and $V'_1(0) = V'_2(0)$. We shall in general be interested in cases where $(V(\phi)$ is monotonically decreasing and concave ($V''(\phi) < 0$ and $V'''(\phi) > 0 \ \forall \phi$). This is a generalized form of the potential used by Peebles and Vilenkin [16] in which

$$V_{PV}^1(\phi) = \lambda(\phi^4 + M^4), \quad (11)$$
$$V_{PV}^2(\phi) = \frac{\lambda M^8}{\phi^4 + M^4}. \quad (12)$$

Since we require inflation to occur at early times, when the expectation value of $\phi$ is large and negative, we would like $V_1(\phi)$ to satisfy the slow-roll conditions

$$\epsilon \equiv \left[ \frac{V'_1(\phi)}{V_1(\phi)} \right]^2 \frac{M_p^2}{16\pi} \ll 1,$$
$$\tilde{\eta} \equiv \left| \frac{V''_1(\phi)}{V_1(\phi)} \right| \frac{M_p^2}{8\pi} \ll 1, \quad (13)$$

for sufficiently large and negative $\phi$.

Inflation ends when the slow-roll conditions are violated and the potential and kinetic energies of the inflaton are comparable with each other. We denote this epoch by the
subscript $x$, following Peebles and Vilenkin [16], so that the above statement reads

$$V_1(\phi_x) \simeq \frac{1}{2} \phi_x^2.$$  \hspace{1cm} (14)

In the models in this paper we will always have $\dot{\phi}_x \simeq -M_p$ and (8) then implies that the Hubble parameter at this epoch is given by

$$H_x = \sqrt{\frac{8\pi V_1(-M_p)}{3M_p^2}}.$$  \hspace{1cm} (15)

In traditional inflationary models the inflaton then rapidly transfers its energy to other fields either through perturbative effects (reheating) or parametric resonance (preheating). Here, however, there is no such effect, and it is the kinetic energy of the field $\phi$ that is the dominant component of the energy density of the universe immediately after the end of inflation. Following Joyce [34] we term this behavior *kination*. Since time derivatives of scalar fields scale as $\dot{\phi}(t) \propto a(t)^{-3}$ we may use (14) to show that the evolution of the energy density in $\phi$ during the kination era obeys

$$\rho_\phi(a) \simeq V_1(-M_p) \left(\frac{a_x}{a}\right)^6.$$  \hspace{1cm} (16)

Now, assuming spatial flatness (which the epoch of inflation will ensure in general), we may solve the Friedmann equation (3) for the cosmic scale factor to obtain

$$\left[\frac{a(t)}{a_x}\right]^3 = 3\sqrt{\frac{8\pi V_1(-M_p)}{3M_p^2}} t.$$  \hspace{1cm} (17)

Since the universe is kinetic energy dominated during this epoch, we may also obtain

$$\phi(t) = \frac{M_p}{\sqrt{12\pi}} \ln \left(\frac{t}{t_x}\right) - M_p = \frac{3M_p}{\sqrt{12\pi}} \ln \left(\frac{a}{a_x}\right) - M_p,$$  \hspace{1cm} (18)

where we have imposed $\phi(t_x) = -M_p$ and

$$t_x = \frac{1}{3} \sqrt{\frac{3M_p^2}{8\pi V_1(-M_p)}}$$  \hspace{1cm} (19)

is the cosmic time at which inflation ends.

A successful cosmology requires the universe be radiation-dominated at the time of nucleosynthesis, since otherwise the precision predictions of that theory are no longer in agreement with observations. The lack of conventional reheating, the conversion of the potential energy
of the inflaton to particle production, in quintessential inflation means that the requisite radiation must be produced another way. In fact, the radiation era in these models is due to the subtle behavior of quantum fields in changing geometries.

At the end of inflation, the FRW line element undergoes an abrupt change from that associated with cosmic expansion (exponential or power-law) to that associated with kination. Massless quantum fields in their vacua in the inflation era are no longer in vacuum in the kination era, corresponding to gravitational particle production. This effect is analogous to Hawking radiation, and has been explored in detail \[33\]-\[37\] in the cosmological context of interest here.

The radiation density produced in this way is

\[ \rho_r = RH_x^4 \left( \frac{a_x}{a} \right)^4, \]  

(20)

where \( R \sim 10^{-2} \), and the number density of massless particles produced is given by

\[ n \sim RH_x^3 \left( \frac{a_x}{a} \right)^3. \]  

(21)

At such early times in the universe, thermal equilibrium is not yet established due to the rapid pace of cosmic expansion. The massless particles produced by the effects of quantum fields in our changing space-time only establish thermal equilibrium when the Hubble parameter has dropped to a value \( n\sigma \sim H \), where \( \sigma \) is the particle-antiparticle annihilation cross-section.

Since

\[ \sigma \sim \frac{\alpha^2 a^2}{H_x^2 a_x^2}, \]  

(22)

where \( \alpha \sim 0.01 - 0.1 \) is a coupling constant, we obtain

\[ \frac{a_{th}}{a_x} \sim \frac{1}{\alpha \sqrt{R}}. \]  

(23)

Therefore, thermalization takes place at a temperature

\[ T_{th} \sim \alpha R^{3/4} \sqrt{\frac{8\pi V_1(-M_p)}{3M_p^2}}. \]  

(24)

This is the highest temperature at which there is thermal equilibrium in the universe.

Now that we know how both the scalar field and the radiation evolve during the kinetic energy dominated era, we may easily calculate the scale factor at which radiation-domination
occurs. Demanding that $\rho_\phi(a_r) = \rho_r(a_r)$ we obtain
\[ \frac{a_r}{a_x} \sim \frac{3M_p^2}{8\pi \sqrt{RV_1(-M_p)}}, \tag{25} \]
and the temperature $T_r$ at which radiation domination begins is then simply calculated to be
\[ T_r \sim \left(\frac{8\pi}{3}\right)^{3/2} R^{3/4} V_1(-M_p) M_p^3, \tag{26} \]
so that we have
\[ \frac{T_{\text{th}}}{T_r} \sim \left(\frac{3}{8\pi}\right) \alpha \frac{M_p^2}{V_1(-M_p)}. \tag{27} \]

Thus, for $\alpha \sim 0.1$, since the scale of the potential is no greater than $M_p^4$, we see that $T_{\text{th}} > T_r$ and so when radiation-domination begins the universe is immediately in thermal equilibrium.

Clearly, to obtain our standard cosmology it is necessary to have a period of radiation domination (followed by matter domination) before dark energy domination begins. This means requiring that $a_r < a_*$, so that at $T_r$ the universe becomes dominated by radiation, with the scalar field evolving in the background, its potential and kinetic energies subdominant to the radiation. In particular, nucleosynthesis takes place at $T_{\text{nuc}} \sim 1\text{MeV}$. To ensure that the universe is radiation dominated at this epoch we should conservatively require $T_r \geq 10\text{MeV}$. Using (26), this allows us to bound the energy scale associated with quintessential inflation by
\[ V_1^{1/4}(-M_p) \geq 10^{-5}M_p. \tag{28} \]

Such a bound does not exist for the standard inflationary paradigm because reheating effects ensure that the universe is radiation-dominated immediately following inflation. In quintessential inflation however, there is a comparatively small amount of radiation produced through gravitational particle production, so that radiation-domination occurs much later.

For a significant time subsequent to this, cosmic evolution is much the same as in the standard cosmology, with a matter dominated epoch eventually succeeding the radiation era. Although the scalar field is not important during these times, the density fluctuations seeded by quantum fluctuations in $\phi$ during inflation lead to structure formation and temperature fluctuations in the cosmic microwave background radiation. In order to obtain agreement
with the COBE anisotropy measurements, we must require

\[
\frac{V^{3/2}(\phi_i)}{M_p^3|V'(\phi_i)|} \sim 5.2 \times 10^{-5},
\]

where \(\phi_i\) denotes the value of \(\phi\) 60 efolds before the end of inflation. Note that in the Peebles-Vilenkin model this constraint translates to \(\lambda \sim 10^{-14}\), similar to that tuning required by standard chaotic inflationary potentials.

Finally, let us turn to this extreme future of the universe, in which the scalar field is rolling in the potential \(V_2(\phi)\) and becomes responsible for quintessence at an epoch denoted by a subscript \(\ast\). It is clear that when quintessence begins, the energy density of the field \(\phi\) once again becomes dominated by its potential energy density. This implies

\[
V_2(\phi_\ast) = V_1(-M_p) \left(\frac{a_x}{a_\ast}\right)^6,
\]

where

\[
\phi_\ast = \frac{3M_p}{\sqrt{12\pi}} \ln \left(\frac{a_\ast}{a_x}\right) - M_p.
\]

It is a challenge similar to that for conventional quintessence to ensure that this epoch occurs at the present time and yields the correct ratio of matter to dark energy. However, such considerations apply far after baryogenesis and we shall refer the reader to other treatments for the details of how this occurs [15]-[17].

III. QUINTESSENTIAL BARYOGENESIS

In order for the quintessence field \(\phi\) to play a role in baryogenesis, we must consider how \(\phi\) couples to other fields. In principle, the inflaton and quintessence field may lie in any sector of the theory, the phenomenologically safest of which would be one in which there are only gravitational strength couplings to other particles. This is presumably the best we can do, since attempts to protect gravitational-strength couplings through global symmetries can be thwarted by wormholes and quantum gravitational effects [38]-[41]. We will adopt a conservative approach and assume that \(\phi\) couples to standard model fields with couplings specified by a dimensionless constant and an energy scale which we shall leave as a free parameter for the moment and later constrain by observations and the condition that our model produce a sufficient baryon asymmetry.
We consider terms in the effective Lagrangian density of the form

\[ L_{\text{eff}} = \frac{\lambda'}{M} \partial_\mu \phi J^\mu, \]  

where \( \lambda' \) is a coupling constant, \( M < M_p \) is the scale of the cutoff in the effective theory and \( J^\mu \) is the current corresponding to some continuous global symmetry such as baryon number or baryon number minus lepton number. Further, let us assume that \( \phi \) is homogeneous. We then obtain

\[ L_{\text{eff}} = \frac{\lambda'}{M} \dot{\phi} \Delta n \equiv \mu(t) \Delta n, \]  

where \( n = J^0 \) is the number density corresponding to the global symmetry and we have defined an effective time-dependent “chemical potential” \( \mu(t) \equiv \lambda' \dot{\phi}/M \).

Recall that we need to satisfy the Sakharov criteria in order to generate a baryon asymmetry (for reviews see [42]-[45]). The first of these requires baryon number \( B \) to be violated. At this stage, to maintain generality, we shall leave the mechanism of baryon number violation unspecified, and will address particular cases later. Possible sources are the decay of superheavy grand-unified gauge bosons or anomalous electroweak processes at finite temperature. Further, the standard model is maximally C-violating due to its chiral structure, and the coupling (32) is \( CP \)-odd. In this sense, no explicit \( CP \)-violation is required in this model. The third Sakharov criterion requires a departure from thermal equilibrium if \( CPT \) is a manifest symmetry. However, the crucial point about baryogenesis in the presence of the rolling scalar field \( \phi \) is that \( CPT \) is broken spontaneously by the explicit value taken by \( \langle \dot{\phi} \rangle \neq 0 \). Thus, the particular model of baryogenesis that is important here is spontaneous baryogenesis [46], which is effective even in thermal equilibrium. We will refer to this model, in which the rolling scalar responsible for both inflation and dark energy also provides a source for spontaneous baryogenesis as quintessential baryogenesis. In our model it is the field \( \phi \) that plays the role of Cohen and Kaplan’s thermion [46]. The idea that the inflaton could drive spontaneous baryogenesis was discussed briefly in [46], where it was correctly noted that accelerated expansion would reduce the baryon number generated during inflation to a negligible magnitude, and therefore that barogenesis during the reheating phase was the only possibility. However, as we shall see, in the context of quintessential inflation there exists a significant range of postinflationary cosmic history during which spontaneous baryogenesis may occur.
To understand how spontaneous (and hence quintessential) baryogenesis works, note that in thermal equilibrium we have

\[
\Delta n(T; \xi) = \int \frac{d^3 \mathbf{p}}{(2\pi)^3} \left[ f(E, \mu) - f(E, -\mu) \right],
\]

(34)

where \( \xi \equiv \mu/T \) is a parameter and \( f(E, \mu) \) is the phase-space distribution of the particles of the current \( J^\mu \), which may be Fermi-Dirac or Bose-Einstein. Thus, for \( \xi < 1 \)

\[
\Delta n(T; \mu) \simeq gT^3 \frac{\xi}{6} + \mathcal{O}(\xi^2),
\]

(35)

where \( g \) is the number of degrees of freedom of the field corresponding to \( n \). Therefore,

\[
\Delta n(T; \mu) \simeq \frac{N'g}{6M} T^2 \dot{\phi}.
\]

(36)

Now recall that the entropy density is given by

\[
s = \frac{2\pi}{45} g_* T^3,
\]

(37)

where \( g_* \) is the effective number of relativistic degrees of freedom in thermal equilibrium at temperature \( T \). Whatever the mechanism of baryon number violation, there will exist a temperature \( T_F \) below which baryon number violating processes due to this mechanism become sufficiently rare that they freeze out. For \( T < T_F \) these processes can no longer appreciably change the baryon number of the universe. Computing the freeze-out value of the baryon to entropy ratio we then obtain

\[
\eta_F \equiv \eta(T_F) \equiv \frac{\Delta n}{s}(T_F) \simeq 0.38 \sqrt{\frac{\lambda' g}{g_*}} \frac{\dot{\phi}(T_F)}{MT_F}.
\]

(38)

Quintessential baryogenesis is effective at temperatures \( T_{th} > T > T_F \), with corresponding scale factors \( a_{th} \equiv a(t_{th}) < a(t) < a_F \equiv a(t_F) \). We have seen that \( T_{th} < T_r \), so that baryon number violating interactions are first in equilibrium during the kination epoch in which the evolution of the scalar field can be written as

\[
\dot{\phi} \simeq \sqrt{2V_1(-M_p)} \left( \frac{a_x}{a} \right)^3.
\]

(39)

If we assume that \( T_F > T_r \) then (38) yields

\[
\eta_F \sim 3.8 \times 10^{-3} \sqrt{\frac{2V_1(-M_p)}{M_p^4}} \frac{M_p}{T_F} \left( \frac{M_p}{M} \right) \left( \frac{a_x}{a_F} \right)^3 \lambda',
\]

(40)
where we have used \((g/g_*) \sim 10^{-2}\). Writing, for convenience,

\[ a_F = \gamma a_{\text{th}} , \]  

(41)

with \(1 \leq \gamma \leq 10^6\) (so that \(T_F = T_{\text{th}}/\gamma\)) and using (23) we then obtain

\[ \eta_F \sim 5.4 \times 10^{-3} \times \sqrt{\frac{8\pi}{3}} \lambda' \gamma^{-2} R^{3/4} \alpha^2 \left( \frac{M_p}{M} \right) . \]  

(42)

Inserting the values for \(R\) and \(\alpha\) yields

\[ \eta_F \simeq 2 \times 10^{-7} \times \lambda' \gamma^{-2} \left( \frac{M_p}{M} \right) . \]  

(43)

Notice that, as expected, this final result is linear in the effective chemical potential, and contains a power of \(\gamma\) reflecting the appropriate amount of redshifting occurring during the kination epoch between the temperatures \(T_{\text{th}}\) and \(T_F\).

To make further progress we must calculate \(T_F\) and \(\phi(T_F)\) in order to find \(\gamma\), and this requires knowledge of the dynamics of \(\phi\). To do this correctly it is necessary to consider the possible effects of back-reaction of the coupling to \(J^\mu\) on the dynamics of \(\phi\). Taking account of the effective Lagrangian into account, the equation of motion of \(\phi\) becomes

\[ \ddot{\phi} + 3H \dot{\phi} + V'(\phi) + \frac{\lambda'}{M} \Delta n + 3 \frac{\lambda'}{M} H \Delta n = 0 . \]  

(44)

Using (35) we obtain

\[ \left[ 1 + \frac{\lambda' g_6}{6} \left( \frac{T}{M} \right)^2 \right] \left( \ddot{\phi} + 3H \dot{\phi} \right) + V'(\phi) = 0 . \]  

(45)

Therefore, for \(T < M\), we are justified in neglecting the extra term and we may safely neglect the back-reaction on \(\phi\). This approximation will typically easily be satisfied, and in particular it is well-justified in the Peebles-Vilenkin model \([16]\) in which \(T \ll M\).

How the baryon excess evolves after this point depends on the value of \(T_F\) and on the relevant current in equation (32). If \(T_F \leq T_{c}^{\text{EW}} \sim 100\ \text{GeV}\), the critical temperature of the electroweak phase transition, then all baryon number violation ceases at \(T_F\) and \(\eta(T < T_F) = \eta_F\). However, if \(T_F > T_{c}^{\text{EW}}\), then we must take into account the effects of anomalous electroweak processes at finite temperature. These can be involved in directly generating the baryon asymmetry (electroweak baryogenesis), in reprocessing an asymmetry in other quantum numbers into one in baryon number (for example in leptogenesis) or in diluting the asymmetry created by any baryogenesis mechanism which is effective above
the electroweak scale and does not produce a $B - L$ asymmetry. It is important to realize that, in the context of quintessential inflation, the quantitative effects of these electroweak processes may differ substantially from those in the standard cosmology \cite{34, 47} since in our case, the electroweak phase transition may occur during kinflation rather than radiation domination.

In the electroweak theory baryon number violating processes at zero temperature are mediated by a saddle-point field configuration known as the sphaleron \cite{48}. We shall therefore refer to the finite temperature configurations relevant here as thermal sphalerons.

However, given the constraints we shall present in section \ref{sec:constraints}, we shall see that it is necessary to have $\gamma \geq 10^2$ and no electroweak dilution in order to generate a sufficient baryon asymmetry.

We have now provided quite a general description of quintessential baryogenesis. While this has allowed us to demonstrate the generic features of our model, we cannot calculate the magnitude of the actual baryon asymmetry generated without first specifying a mechanism of baryon number violation (and hence a value for $T_F$) and a value for the dimensionless combination $\lambda' M_p / M$. Let us now turn to some concrete examples.

\subsection*{A. Baryon Number Violation Through Non-renormalizable Operators}

If there exists baryon number violating physics above the standard model, then this physics will manifest itself in non-renormalizable operators in the standard model. For the purposes of this section we will actually be interested in operators that violate the anomaly-free combination $B - L$. In that case the value $\eta_F$ calculated via \eqref{eq:etaF} will be the final baryon to entropy ratio $\eta$, since anomalous electroweak processes preserve this combination of quantum numbers. Consider the effective 4-fermion operator

$$\mathcal{L}_{B-L} = \frac{\tilde{g}}{M_X^2} \psi_1 \psi_2 \bar{\psi}_3 \bar{\psi}_4, \quad (46)$$

where $\psi_i$ denote standard model fermions. Here $\tilde{g}$ is a dimensionless coupling, obtained after integrating out the $B - L$ violating effects of a particle of mass $M_X$. The rate of baryon number violating processes due to this operator is, as usual, defined by $\Gamma_{B-L}(T) = \langle \sigma(T)n(T)v \rangle$, where $\sigma(T)$ is the cross-section for $\psi_1 + \psi_2 \to \psi_3 + \psi_4$, $n(T)$ is the number density of $\psi$ particles, $v$ is the relative velocity and $\langle \cdots \rangle$ denotes a thermal average. For
temperatures $T < M_X$ we have $n(T) \sim T^3$, $\sigma(T) \sim \tilde{g}^2 T^2 / M_X^4$, and $v \sim 1$ which yields

$$\Gamma_{B-L}(T) \sim \frac{\tilde{g}^2}{M_X^4} T^5.$$  \hspace{1cm} (47)

The high power of the temperature dependence in this rate results from the fact that (46) is an irrelevant operator in the electroweak theory and, as we shall see, is crucial for the success of our mechanism. These interactions are in thermal equilibrium in the early universe, but because their rate drops off so quickly with the cosmic expansion they will drop out of equilibrium at the temperature $T_F$ defined through

$$\Gamma_{B-L}(T_F) = H(T_F).$$  \hspace{1cm} (48)

Thus,

$$T_F \simeq \left( \frac{3}{8\pi} \right)^{3/4} \frac{1}{\tilde{g} R^{3/8}} \sqrt{\frac{M_p^2}{V_1(-M_p)} M_X^2}.$$  \hspace{1cm} (49)

As a definite example, let us take $\lambda M_p / M \sim 8$. As we shall see in IV, this is as large as is allowed by current constraints. In this case (43) implies that we need $\gamma \sim 10^2$ in order to obtain the correct BAU. This then implies that we must have

$$M_X \sim \left( \frac{8\pi}{3} \right)^{5/8} R^{9/16} \sqrt{\frac{\alpha \tilde{g} V_1(-M_p)}{M_p^2}}.$$  \hspace{1cm} (50)

For reference, note that in the Peebles-Vilenkin model [16] this becomes $M_X^{PV} \sim 10^{11}$ GeV, the intermediate scale that appears in some supersymmetric models.

## B. Grand Unified Baryon Number Violation

A natural source of baryon number violation to consider is that arising from gauge-mediated interactions in grand unified theories (GUTs), in which quarks and leptons lie in the same representation of the GUT gauge group. However, as we shall demonstrate briefly here, this is not a viable source for quintessential baryogenesis.

Consider a GUT gauge boson $X$ with mass $M_X$. This particle decays through a renormalizable coupling with decay rate

$$\Gamma_{\text{GUT}} \simeq \frac{\tilde{g}_X^2}{16\pi} \begin{cases} M_X & \text{for } T \leq M_X, \\ T & \text{for } T \geq M_X. \end{cases}$$  \hspace{1cm} (51)
It is then relatively straightforward to show that $X$ particles decay in equilibrium (as is required by our mechanism) only if

$$M_X \leq 10^9 \left( \frac{T_r}{\text{GeV}} \right)^{1/2} \text{GeV}.$$  \hspace{1cm} (52)

Since GUT gauge bosons are much heavier than this value this source of baryon number violation is not useful for our model.

One might be tempted to consider a particle with the appropriate mass produced by some physics beyond the standard model. However, it quickly becomes apparent that decays through renormalizable interactions can never work, since these particles then decay and the interactions are never again out of equilibrium (there is no freeze-out). Thus, as we commented earlier, nonrenormalizable operators seem essential for decays of massive particles to work.

C. Electroweak Baryon Number Violation

Another important source of baryon number violation in baryogenesis models comes from anomalous electroweak processes at finite temperature. Thermal sphalerons are in equilibrium at temperatures above the critical temperature of the electroweak phase transition $T_{c}^{\text{EW}} \sim 10^2 \text{GeV}$ and thus, the final baryon to entropy ratio generated purely through electroweak processes is given by (43) with $\gamma_{\text{EW}} \sim 10^7$. This yields

$$\eta_{\text{EW}} \sim 2 \times 10^{-21} \left( \frac{\lambda' M_p}{M} \right),$$

which, since $\lambda' M_p/M < 8$, is far too low to explain the observed baryon asymmetry (2).

IV. CONSTRAINTS AND TESTS

The presence of an extremely light scalar field in the universe has the potential to lead to a number of observable consequences in the laboratory and in cosmology. In the case of quintessence these effects have been analyzed in some detail by Carroll [49]. Particularly strong constraints arise due to couplings of the form

$$\mathcal{L}^{(1)}_{\text{eff}} \equiv \beta_{FF} \frac{\phi}{M} F_{\mu\nu} \tilde{F}^{\mu\nu},$$

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where $F_{\mu\nu}$ and $\tilde{F}^{\mu\nu} \equiv \frac{1}{2} \epsilon^{\mu\rho\sigma} F_{\rho\sigma}$ are the electromagnetic field strength tensor and its dual respectively. If, as in quintessence models, the field $\phi$ is homogeneous and time varying, then it affects the dispersion relation for electromagnetic waves and leads to a rotation in the direction of polarized light from radio sources \cite{50, 51}. The bound obtained is

$$|\beta_{F\tilde{F}}| \leq 3 \times 10^{-2} \left( \frac{M}{M_p} \right),$$

(55)

where we have assumed that $\phi$ rolls over about a Planck mass during the last half a redshift of the universe.

To avoid this bound, quintessence models usually struggle to have such couplings be as small as possible. In our model however, we are making important use of the coupling \cite{32}. If the relevant current is that for baryon number $J_B^\mu$, then using the anomaly equation we may rewrite our term \cite{32} as

$$\mathcal{L}_{\text{eff}} = -\frac{\lambda'}{M} \phi n_f \left( \frac{g^2}{32\pi^2} W^a_{\mu\nu} \tilde{W}^{a\mu\nu} - \frac{g'^2}{32\pi^2} B_{\mu\nu} \tilde{B}^{\mu\nu} \right),$$

(56)

where $g$ and $g'$ are the gauge couplings of $SU(2)_L$ and $U(1)_Y$, respectively, $n_f$ is the number of families and $\tilde{W}^{\mu\nu}$ and $\tilde{B}^{\mu\nu}$ are the duals of the $SU(2)_L$ and $U(1)_Y$ field strength tensors respectively.

Thus, the constraint \cite{53} should apply to give

$$\left( \frac{n_f g^2}{32\pi^2} \right) \lambda' \leq 3 \times 10^{-2} \left( \frac{M}{M_p} \right),$$

(57)

Putting in the appropriate standard model numbers we obtain

$$\lambda' \frac{M_p}{M} < 8.$$ 

(58)

If the relevant current is $J_{B-L}^\mu$ then the above argument does not apply formally. However, we may still generally expect an analogous coupling to the term $F^{\mu\nu} \tilde{F}_{\mu\nu}$ of a similar order.

We have already demonstrated that, for an appropriate scale $M_X$, successful quintessential baryogenesis takes place for $\lambda' < 8$, and so it is possible to generate the observed BAU and to evade existing constraints. If quintessential baryogenesis is correct, then it may be that the relevant coupling $\lambda'$ lies just below the existing observational bounds and that future studies of the rotation of polarized light from distant galaxies will reveal the presence of such a term.
One might also worry about the laboratory constraints on the operator corresponding to a coupling to $F^{\mu\nu} \tilde{F}_{\mu\nu}$ resulting from (56) after electroweak symmetry breaking. For example, in the standard model these can lead to contributions to the electric dipole moments of the electron and the neutron. These effects have been considered \cite{52, 53} in the context of electroweak baryogenesis. Applying the experimental bounds \cite{54, 55} the relevant bound is significantly weaker than the cosmological one quoted above.

We should also comment briefly on the naturalness of the potentials required. We have taken the quintessence and quintessential inflation paradigms and explored another generic effect in these models. It should be pointed out however, that the potentials require a great deal of fine-tuning. In particular, the extremely small value of the self-coupling $\lambda$ to explain the fluctuations in the CMB is unexplained. If we consider embedding the model in a supersymmetric theory we will come up against the problems that standard quintessence faces, that the flatness of the potential required at late times in the universe may be ruined by the soft supersymmetry breaking terms required to obtain a realistic phenomenology today.

\section{Comments and Conclusions}

Modern particle cosmology concerns the search for dynamical explanations for the initial conditions required by the standard FRW cosmology. Particular attention has been paid recently to one of these initial condition problems, that of the size of the vacuum energy contribution to the total energy density of the universe. The root of this issue is that vacuum energy, or something approximating it, can lead to the acceleration of the universe.

The best fit cosmology to all current observational data is one in which the universe undergoes two separate epochs of acceleration. The first of these, inflation, is the only clear way to seed adiabatic, scale-free perturbations in the cosmic microwave background radiation. The second epoch, that of dark-energy domination, is required to simultaneously understand the power spectrum of the CMB and the expansion of the universe at intermediate redshifts as revealed by type IA supernovae. These requirements have led cosmologists to introduce a new scalar field to account for the newly required late-time acceleration.

In this paper we have extended the approach of Spokoiny \cite{15} and of Peebles and Vilenkin \cite{16} in exploring the extent to which the dynamics of a single scalar field can be
FIG. 1: The evolution of the universe in this model. As the scalar field rolls down its potential the universe goes through a succession of phases, beginning with inflation, generating the baryon asymmetry along the way, and ending with dark energy domination.

It seems a particularly powerful idea to us that a single rolling scalar field might be responsible for a number of the fundamental initial conditions required to make the standard cosmology work. In the case of the baryon asymmetry, this allows us to associate the existence of an asymmetry with the spontaneous breaking of CPT and the direction of the rolling of the scalar field.

The evolution of the universe we envisage may be summarized as follows (see figure 1). At the earliest times in the universe, inflation occurs due to the potential energy dominance of the field $\phi$ which begins rolling at very large and negative values. Inflation ends when the kinetic energy of the scalar field becomes important and the slow-roll conditions are violated. Since our potential does not have a minimum at finite $\phi$, unlike typical inflationary models,
conventional reheating does not occur. Instead, matter is created gravitationally due to the mismatch of vacuum states between the approximately de-Sitter state of inflation and that of the kination era. Since kinetic energy density redshifts more rapidly than radiation energy density, the universe eventually becomes radiation-dominated. At this stage the rolling scalar has negligible effect on the expansion rate of the universe. However, the direction of rolling spontaneously violates $CPT$. If $\phi$ couples to other fields, as we expect it to generically, then the expectation value of the baryon number operator in this background in thermal equilibrium is nonzero. Thus a baryon excess is generated. After the electroweak phase transition baryon number violation is no longer effective in the universe and the baryon number existing at that time is frozen in. The scalar field continues to evolve and in the late universe, after matter-domination has begun, its potential energy can once again become dominant leading to a new period of dark-energy domination.

The couplings required to make quintessential baryogenesis effective may be generated in a number of different ways, for example by gravitational effects coupling the inflaton/dark energy sector to visible sector fields. We have considered the current experimental constraints on the necessary operator and have found that there exist considerable regions of parameter space in which our mechanism is consistent. Further, it is possible that a restricted region of this parameter space may be accessible to future experiments.

We have left a number of questions unanswered and will return to them in future work. Perhaps the most pressing issue is one that plagues rolling scalar models of dark energy in general, namely the question of technical naturalness of the potentials involved, and their stability to quantum corrections. We have omitted any discussion of this here, while laying out the general features of the model, but these issues must be addressed to put our mechanism on firmer ground. For example, it may be most natural to identify the field $\phi$ with a pseudo-Goldstone boson $[27, 56]$, since its coupling to the current $J^\mu$ is derivative. However, this is a general issue for quintessence models, and is not specific to our baryogenesis mechanism. We have therefore chosen to concern ourselves with this issue separately.

Taken at face value, current observations imply that our universe is entering an accelerating phase that may be governed by the rolling of a scalar field. If we are to understand the physics of such a field then it is important that we investigate other ways in which it may impact cosmology and particle physics. In particular, if, as we have suggested here,
the field is responsible for the generation of the baryon asymmetry, then the result will be a more economical and attractive cosmology.

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