Can the baryon asymmetry arise from initial conditions?

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Can the Baryon Asymmetry Arise From Initial Conditions?

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In this letter, we quantify the challenge of explaining the baryon asymmetry using initial conditions in a universe that undergoes inflation. Contrary to lore, we find that such an explanation is possible if net $B - L$ number is stored in a light bosonic field with hyper-Planckian initial displacement and a delicately chosen field velocity prior to inflation. However, such a construction may require extremely tuned coupling constants to ensure that this asymmetry is viably communicated to the Standard Model after reheating; the large field displacement required to overcome inflationary dilution must not induce masses for Standard Model particles or generate dangerous washout processes. While these features are inelegant, this counterexample nonetheless shows that there is no theorem against such an explanation. We also comment on potential observables in the double β-decay spectrum and on model variations that may allow for more natural realizations.

INTRODUCTION

The observed flatness, isotropy, and homogeneity of the universe on large scales strongly suggest that the universe underwent a period of inflation at early times (for a review see [1]). An important consequence of this epoch is that pre-inflationary relics are exponentially diluted away, thereby eliminating monopoles, cosmic strings, and other topological defects from the observed universe. Such a dilution also depletes any primordial $B - L$ number, so it is generally accepted that a dynamical mechanism satisfying the Sakharov conditions [2] is required to generate the matter asymmetry after inflation.

To see the difficulty explicitly, consider the most favorable initial condition for preserving an asymmetry stored in Standard Model (SM) particles: a universe populated with a degenerate gas of massless, fully asymmetric neutrinos ($n_\nu \neq n_{\bar\nu} = 0$) at zero temperature. Denoting pre-inflationary quantities with an i subscript, the energy density of this ensemble is

$$\rho_{\nu,i} = \frac{1}{4\pi^2}(3\pi^2 n_{\nu,i})^{4/3} \,, \tag{1}$$

which is completely determined by the number density $n_{\nu,i}$, so after N inflationary e-folds, the diluted asymmetry is

$$Y_\nu \equiv \frac{n_\nu}{s} = \frac{45(4\pi^2\rho_{\nu,i})^{3/4}}{6\pi^4 g_* T_R^3}\, e^{-3N} \,, \tag{2}$$

where s is the entropy density, T_R is the reheating temperature, and g_* is the number of relativistic degrees of freedom. Even with a Planckian energy density and the lowest viable reheating temperature the asymmetry is at most

$$Y_\nu^{\rm max} \sim 4\times 10^{-15}\left(\frac{10\,{\rm MeV}}{T_R}\right)^3 \left(\frac{\rho_{\nu,i}}{m_{Pl}^4}\right)^3, \tag{3}$$

so even if $Y_\nu^{\rm max}$ is all converted into a comparable baryon asymmetry, it falls many orders of magnitude below the observed value $Y_B^{\rm obs} \equiv (n_B - n_{\bar B})/s = (8.6 \pm 0.1)\times 10^{-11}$ [3]. It is also clear that, for an initial asymmetry of order $\sim Y_\nu^{\rm max}$, the corresponding energy density $\rho_{\nu,i}$ easily dominates the pre-inflationary universe, thereby preventing inflation from starting in the first place.

However, this observation assumes that the number density is stored in fermions for which the energy and number densities are related by Fermi statistics as in Eq. (1). If, instead, the asymmetry is stored in a bosonic field, these quantities may be decoupled from each other and it is possible, in principle, to engineer a relatively low energy configuration with a large number density.

In this letter, we study the conditions under which a large asymmetry can exist prior to inflation, survive ~ 60 e-folds of dilution, and viably yield the observed baryon asymmetry after reheating. We find that such a maneuver is possible, but requires super-Planckian field values and significant tuning to stabilize the electroweak sector and prevent the asymmetry from being washed out. Nonetheless, the simple realization presented here serves as a proof of principle to demonstrate that a sizable asymmetry can, indeed, survive inflation and be transferred to SM particles. Although [4] initially suggested the possibility of storing an asymmetry in a light boson and transferring it to the SM, to our knowledge this work presents the first explicit model to realize such a mechanism in a universe with inflation.

MODEL DESCRIPTION AND INITIAL SETUP

Inspired by the Affleck-Dine mechanism [5], we consider a complex scalar that carries a global conserved ϕ number, which we will later identify with $B - L$.[1] The Lagrangian for this setup is

$$\mathscr{L} = |\partial_\mu\phi|^2 - m_\phi^2|\phi|^2 - V(|\phi|) \,, \tag{4}$$

which preserves ϕ number in its potential. Defining $\phi = \frac{1}{\sqrt{2}} r e^{i\theta}$, this can be written in polar coordinates as

$$\mathscr{L} = \frac{1}{2}\left(\dot r^2 + r^2\dot\theta^2 - m_\phi^2 r^2\right) - V(r) \,. \tag{5}$$

The ϕ number density is

$$n_\phi = i(\phi^*\partial_t\phi - \phi\,\partial_t\phi^*) = r^2\dot\theta \,, \tag{6}$$

so the angular equation of motion can be written as

$$a^{-3}\,\partial_t(a^3 n_\phi) = -\partial_\theta V = 0 \,, \tag{7}$$

where a is the FRW scale factor. Like in Affleck-Dine models, we assume that ϕ is initially displaced far from the origin in field space. Unlike in Affleck-Dine, this displacement is established *before* inflation and we allow no explicit ϕ number violation in the Lagrangian (*i.e.* the potential is independent of θ) so Eq. (7) represents the conservation of comoving ϕ number.

To generate a sizable asymmetry from initial conditions, we require that

- The asymmetry in ϕ-number before inflation must be exponentially large to survive 60 e-folds of inflation, which implies $n_{\phi,i} = r_i^2\dot{\theta}_i \gg m_{Pl}^3$.
- This large number density must not dominate the the energy budget of the pre-inflationary universe, $\rho_{\phi,i} < \rho_{\rm inf}$, so that the inflaton's potential can successfully drive inflation in the presence of such a large number density.
- ϕ number must consistently be identified with $B-L$ and that the corresponding asymmetry must be transferred to the SM through $B-L$ preserving interactions.

Naively, these are difficult tasks since values of $n_{\phi,i}$ that survive 60 e-folds require very large field values, which generically favor a comparably large energy density

$$\rho_{\phi,i} \;=\; \frac{1}{2}\left(\dot{r}_i^2 + r_i^2\dot{\theta}_i^2 + m_\phi^2 r_i^2\right) + V(r_i)\,. \tag{8}$$

However, consider initial conditions in which

$$r_i \gg m_{Pl} \quad , \quad \dot{r}_i \ll \sqrt{\rho_{\rm inf}} \quad , \quad m_\phi \ll \dot{\theta}_i \ll \frac{\sqrt{\rho_{\rm inf}}}{r_i} \quad , \tag{9}$$

where $\rho_{\rm inf}$ is the energy density during inflation. In this limit, ϕ is relativistic prior to inflation and the physical energy density is dominated by the angular momentum term, so we have

$$\rho_{\phi,i} \;\simeq\; \frac{r_i^2\dot{\theta}_i^2}{2} = \frac{n_{\phi,i}\,\dot{\theta}_i}{2} \ll \rho_{\rm inf}\,, \tag{10}$$

where the contributions from $V(r_i)$ and $m_\phi^2 r_i^2$ have been tuned away by the bare cosmological constant. Note that the number density $n_{\phi,i} = r_i^2\dot{\theta}_i$ can be very large while the average energy per particle $\rho_{\phi,i}/n_{\phi,i} \sim \dot{\theta}_i$ can be small for suitable choices of r_i and $\dot{\theta}_i$. Note also that m_ϕ must be negligible compared to all other scales in the problem because, for a non-relativistic ensemble, the energy density ($\rho_\phi \geq m_\phi n_\phi$) cannot be decoupled from the number density. However, if ϕ is the Nambu-Goldstone boson (NGB) of a broken global symmetry, its mass is protected by a shift symmetry and can naturally be light; we will return to this possibility below.

Assuming instantaneous reheating to temperature T_R immediately following inflation, the post-inflationary ϕ asymmetry is

$$Y_\phi \;=\; \frac{n_\phi}{s} = \frac{45\,e^{-3N}\,n_{\phi,i}}{2\pi^2 g_* T_R^3}\,, \tag{11}$$

where g_* is the number of relativistic degrees of freedom. Requiring $N=60$ e-folds during inflation, we find

$$Y_\phi \simeq 10^{-10}\left(\frac{53\,\mathrm{TeV}}{T_R}\right)^3\left(\frac{r_i}{10^{67}\,\mathrm{GeV}}\right)^2\left(\frac{\dot{\theta}_i}{10^{-50}\,\mathrm{GeV}}\right), \tag{12}$$

where we have taken $g_* = 100$. Although the initial field value r_i and angular velocity $\dot{\theta}_i$ have taken on grotesquely hierarchical values, the asymmetry in ϕ is in the right ballpark and the energy density in Eq. (10) remains sub-Planckian

$$\rho_{\phi,i} \simeq 2\times 10^{-43}\, m_{Pl}^4\left(\frac{r_i}{10^{67}\,\mathrm{GeV}}\right)^2\left(\frac{\dot{\theta}_i}{10^{-50}\,\mathrm{GeV}}\right)^2, \tag{13}$$

which, therefore, allows the inflaton potential to dominate the total energy density at early times.

FIELD EVOLUTION POST-INFLATION

To understand how the field evolves during and after inflation, we solve the classical equations of motion

$$a^{-3}\partial_t(a^3\dot{r}) = m_\phi^2 r - r\dot{\theta}^2 - \partial_r V \simeq 0 \tag{14}$$

$$\partial_t(a^3 r^2\dot{\theta}) = -\partial_\theta V = 0\,, \tag{15}$$

where, based on Eq. (9), the $m_\phi^2 r$ and $r\dot{\theta}^2$ terms are negligible in Eq. (14). We also assume that $|\partial_r V|^{1/3} \ll H, \rho_{\rm inf}^{1/4}$ during inflation so that the radial dynamics are dominated by Hubble expansion; since r_i is hyper-Planckian and $\rho_\phi \ll \rho_{\rm inf}$, the $\partial_r V$ term cannot significantly affect the radial displacement. Finally, we demand that all interactions preserve ϕ number, so $\partial_\theta V = 0$ in Eq. (15). Aside from these requirements, the details of this potential are not important for our purposes. Note that the initial values r_i and $\dot{\theta}_i$ in Eq. (9) need not necessarily correspond to a minimum of the potential; they can be imposed by fiat when inflation begins.

Introducing c_r and c_θ as constants of integration with mass-dimension 1 and using the scale factor during inflation $a(t) \propto e^{Ht}$, the system in Eqs. (14) and (15) becomes

$$\dot{r} \simeq c_r^2 e^{-3Ht} \;\;,\;\; \dot{\theta} = \frac{c_\theta^3}{r_i^2}e^{-3Ht}\,, \tag{16}$$

where, for $\dot{\theta}$ we have used the fact that $\dot{r} \sim e^{-3Ht}$, so the initial value r_i is approximately preserved with Hubble expansion. In our regime of interest, Eq. (14) can be approximated as $\partial_t(a^3\dot{r}) = 0$, which yields

$$r(t) \;=\; r_i + \frac{c_r^2}{3H}\left(1 - e^{-3Ht}\right) \;\to\; r_i\;, \tag{17}$$

where the limiting value holds after imposing initial conditions $\dot{r}(0) = 0$ (equivalently $c_r = 0$) and $r(0) = r_i$ from Eq. (9) . There is a corresponding expression for $\theta(t)$ with the replacement $c_r^2 \to c_\theta^3/r_i^2$ and $r_i \to \theta_i$, so the initial radial and angular positions are quickly locked in when inflation begins.

We have also verified numerically that including small corrections from $m_\phi^2 r$, $r\dot{\theta}^2$, and $\partial_r V$ terms in Eq (14) does not affect the solution in Eq. (17) as long as our parametric assumptions hold.

For this setup, we can safely assume that the quantum fluctuations during inflation are negligible as $\langle|\delta\phi|^2\rangle \sim H^2 \ll r_i^2$ for our chosen initial configuration. Also since Eq. (13) implies $\rho_\phi \ll \rho_{\rm DM}$, isocurvature fluctuations are highly suppressed and do not affect CMB observables; below we will also require that ϕ to thermalize with the SM at later times, so any such fluctuations would be erased.

TRANSFERRING THE ASYMMETRY

To communicate the conserved ϕ number to SM fields, we consistently assign ϕ a nonzero $B-L$ number Q_ϕ, we include the $B-L$ preserving interaction

$$\mathscr{L} \supset \frac{1}{\Lambda^{D-2}}(\partial_\mu\phi)\,\hat{\mathcal{O}}^\mu_{B-L} + h.c.\,, \tag{18}$$

where $\hat{\mathcal{O}}_{B-L}$ is some SM operator with mass dimension D and nonzero $B-L$ number such that the combined $\partial_\mu\phi\,\hat{\mathcal{O}}^\mu_{B-L}$ interaction is a $B-L$ singlet. Note that non-derivative operators of the form $\phi\,\hat{\mathcal{O}}_{SM}$ communicate $B-L$ breaking to the SM through r_i insertions, which must be sequestered from the SM to avoid washing out the primordial asymmetry and maintaining sub-Planckian SM energy densities. Achieving this separation may require tuning since there is no shift symmetry to forbid such interactions (see below for a discussion of this issue). After reheating populates the SM, the interaction in Eq. (19) thermalizes ϕ with SM particles, thereby transferring the net $B-L$ number stored in ϕ to the rest of the thermal bath.

One possible realization involves interpreting ϕ number as lepton number so that its asymmetry can be transferred via

$$\mathscr{L} \supset \frac{1}{\Lambda}(\partial_\mu\phi)N^{c\dagger}\bar{\sigma}^\mu N + MNN^c + y_\nu HLN + h.c., \tag{19}$$

where flavor indices have been suppressed and the SM has been extended to include gauge singlet N, its Dirac partner N^c, and the scalar ϕ whose $B-L$ numbers are $+1$, -1 and -2, respectively so that all terms in Eq. (19) preserve $B-L$. Note that the current-like operator $N^{c\dagger}\bar{\sigma}^\mu N$ is not the usual RH neutrino vector current and that MNN^c is a *Dirac* mass and we assume $\Lambda \gg M$. Although "Majoron" interactions of the form ϕNN are allowed by all the symmetries of the setup, the coefficients of such operators must be tuned to ensure that dangerous Majorana masses $\sim r_i NN$ do not arise to washout the asymmetry. Here the $y_\nu HLN$ term is merely an interaction which allows N to decay above the weak scale; not as a source of all active neutrino masses.[2]

Assuming N and N^c are produced thermally during reheating, their dynamics give rise to an additional symmetric population of ϕ which is distinct from the cold pre-inflationary condensate. So long as the $\phi N \leftrightarrow \phi N$ scattering rate exceeds the Hubble rate, the symmetric ϕ population will thermalize with the radiation bath. However, this leading order process only exchanges energy, but not $B-L$ number. To communicate the asymmetry in the ϕ condensate, we demand that the rate for the leading $B-L$ exchanging process $\phi N \to \phi\phi^* N^c$ (depicted in Fig. 1) exceeds Hubble before the N decay to SM states. In the $\Lambda \gg T \gg M$ limit, the cross section for this interaction is approximately $\sigma \sim T^4/8(4\pi)^3\Lambda^6$ so the typical ϕ from the nonthermal pre-inflationary population will exchange $B-L$ number with an N in the thermal bath so long as $n_N\sigma \gtrsim H$, which requires

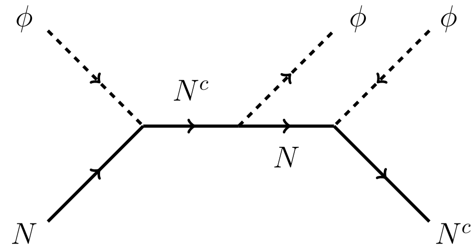

FIG. 1: Scattering process that exchanges lepton number between the pre-inflationary ϕ population to the $N\ N^c$ system after reheating when the condensate evaporates away. Lower order interactions (e.g. $\phi N \to \phi N$) can exchange energy, but do not transfer $B-L$ number.

$$\Lambda \lesssim 0.14\left(\frac{m_{Pl}T_R^5}{\sqrt{g_*}}\right)^{1/6} = 1.2\times 10^3\,\text{TeV}\left(\frac{T_R}{50\,\text{TeV}}\right)^{5/6}, \tag{20}$$

assuming $g_* = 100$. Ssatisfying the condition in Eq. (20) also guarantees thermalization between the N, N^c, and the full ϕ population, including both symmetric-thermal and asymmetric condensate components.[3] After the asymmetry is transferred, $N \to h\nu_L$ decays transmit lepton number to the active neutrinos.

For this particular example, we demand that this transfer occur before the electroweak phase transition so that sphalerons can distribute the asymmetry to the quarks; however this requirement is relaxed if ϕ also carries baryon number. Following the procedure in [6, 7], assuming the interaction in Eq. (18) achieves thermal equilibrium before the electroweak phase transition, the baryon asymmetry is

$$Y_B = \frac{28\,Q_\phi Y_\phi}{22\,Q_\phi^2 + 79}\,, \tag{21}$$

where Q_ϕ is ϕ's $B-L$ number. It is clear that the asymmetry in ϕ from Eq. (12) is sufficient to generate a large baryon asymmetry. For a discussion of the experimental constraints on this scenario, see the supplementary material [8].

TUNING CONSIDERATIONS

Since the ϕ radial displacement is super-Planckian and remains comparably large well after inflation, this value must be sequestered from SM fields. For any scalar field, there is an irreducible tuning in the Higgs portal coupling $H^\dagger H|\phi|^2$ which must be chosen to prevent destabilizing $r_i^2 H^\dagger H$ corrections; indeed, this is an extreme realization of the usual electroweak hierarchy problem. Since no symmetry forbids this operator, its coefficient can only be naturally small if H and ϕ are localized away from each other in a higher dimensional model. In such a scenario, the Higgs portal coefficient can be exponentially suppressed by wave function overlap, but additional model building may be required for ϕ to thermalize with the SM and transfer its asymmetry after reheating. At the non-renormalizable level, there is no symmetry forbidding $B-L$ violating "Planck slop" with potentially large insertions of r_i; if these arise in a full theory of quantum gravity, they must be similarly tuned to prevent washout after reheating.

The scalar ϕ must also be extremely light to decouple its number and energy densities, so its mass and interactions must also be exponentially tuned in this setup. For the particular example in Eq. (19), the coefficients of $B-L$ preserving "Majoron" interactions ϕNN and $\phi^* N^c N^c$ must also be highly tuned to ensure that r_i does not give rise to a large Majorana masses; such terms introduces dangerous $B-L$ breaking washout processes that can erase the asymmetry after reheating.

If, instead, ϕ is interpreted as an NGB, there is partial protection from some of the tuning presented in the example given above. Most favorably, its couplings automatically involve derivatives, so there is no tuning required to protect SM operators from destabilizing r_i insertions. Furthermore, the model in [9] demonstrates that, for an ensemble of p symmetry breaking scalar fields with suitable quartic interactions, the compact field range of an NGB can be exponentially enhanced by factors as large as $\sim 3^p$ relative to the usual maximum excursion $\sim 2\pi f$, where f is the symmetry breaking scale. For $f \sim m_{Pl}$ and $p = 100$, this enhancement allows an NGB to have a field range as large as $\sim 10^{67}$ GeV, which can generate an asymmetry sufficiently large to survive 60 e-folds of dilution and yield the observed asymmetry after thermalization with the SM– see Eq. (12).

However, the NGB realization does not automatically eliminate the Higgs portal tuning encountered in the fundamental scalar example above. Since exponentially enlarging the compact NGB field range requires a large population of natural, Planckian scalars, each of these ~ 100 fields must be prohibited from coupling to the $H^\dagger H$ bilinear and destabilizing the electroweak scale with Planck scale mass parameters. However, as discussed above, higher dimensional locality can exponentially reduce the Higgs portal coupling, so combined with the NGB realization, which protects the mass and ensures derivative interactions, there is a potential roadmap for a more natural theory.

DISCUSSION

In this letter we have presented a simple model to demonstrate that the baryon asymmetry can, in fact, arise from initial conditions in a universe that undergoes inflation. This is accomplished using a light, bosonic field with a large number density prior to inflation. At face value, the benchmark model requires severe tuning to prevent super-Planckian field values from destabilizing the electroweak sector; however, we regard this setup merely as a proof of principle rather than as a compelling model of nature. If, contrary to the example model considered here, the asymmetry carrying field is an NGB in an extra-dimensional setup, it may be possible to build a natural theory that preserves many of the features discussed in this paper. However, doing so requires additional model building to engineer a super-Planckian field range for the asymmetry carrying field.

To be concrete in our discussion, we have chosen a specific operator ansatz to communicate the ϕ asymmetry to the SM through the right-handed neutrino portal, but the key features of our discussion ff do not depend on this choice. For instance, instead of using $\hat{\mathcal{O}}^\mu_{B-L} = N^{c\dagger}\bar{\sigma}^\mu N$ in Eq. (19), we could have assigned ϕ a net baryon number and transferred $B-L$ to quarks using singlet squarks from the Minimal Supersymmetric Standard Model via $\hat{\mathcal{O}}^\mu_{B-L} = (\partial^\mu \tilde{u}^c)\tilde{d}^c\tilde{d}^c$, with qualitatively similar results, up to different constraints introduced by this operator. However, regardless of the particular choice of $\hat{\mathcal{O}}^\mu_{B-L}$, it is essential that the combination $\partial_\mu\phi\,\hat{\mathcal{O}}_{B-L}$ preserve $B-L$ number and that this symmetry is not broken by the interactions that resolve this operator. Indeed, aside from the initial ϕ displacement before inflation, which spontaneously breaks $B-L$, all interactions in the example model from Eq. (19) are $B-L$ singlets.

We note in passing that the effective interaction between ϕ and SM neutrinos (see supplementary material [8]) induces a new double β-decay signature, which contributes a *background* to neutrinoless double β decay searches [10]. The final state for this $2n \to 2p2e\phi$ process has different kinematics relative to conventional double β decay and may be distinguishable using electron momentum distributions. However, since the model in Eq. (19) does not fully account for neutrino masses, the overall rate depends on neutrino yukawa couplings, which are not fixed in the setup described here.

Finally, although this model does not address the electroweak hierarchy problem, it contains many of the ingredients found in the relaxion solution [11] (*e.g.* super Planckian field values for a light, slowly evolving field during inflation) so it would be interesting to see if the setup presented here can accommodate the cosmological relaxation of the electroweak scale as well.

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Notes

[1]Although quantum gravity may spoil this global symmetry, we assume that the relevant dynamics occur at scales for which these effects can be neglected.

[2]After EWSB, for each generation, one linear combination of ν_L and N^c spinors acquires a Dirac mass with N, while the other remains massless. Thus, this simple model does not fully account for active neutrino masses, which are assumed to arise from additional interactions.

[3]For a more detailed discussion of (pseudo)scalar condensate thermalization in the context of axions, see [12–14] which perform more careful calculations. Although a similarly detailed treatment is necessary to determine the exact thermalization requirement here, Eq. (20) suffices as a rough estimate to show that this can easily take place. Note that, unlike axions produced via misalignment, here the asymmetric ϕ population is always present. For a discussion of this issue see [12], which addresses axion thermalization *except* in the case of misalignment production, for which the population is not present before the QCD phase transition.