TESTING B-BALL COSMOLOGY WITH THE CMB

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In D-term inflation models, the fluctuations of squark fields in the flat directions give rise to isocurvature density fluctuations stored in the Affleck-Dine condensate. After the condensate breaks up in B-balls, these can be perturbations in the baryon number, or, in the case where the present neutralino density comes directly from B-ball decay, perturbations in the number of dark matter neutralinos. The latter case results in a large enhancement of the isocurvature perturbation. In this case, the requirement that the deviation of the adiabatic perturbations from scale invariance due to the Affleck-Dine field is not too large imposes a lower bound on the magnitude of the isocurvature fluctuation of about $10^{-2}$ times the adiabatic perturbation. This should be observable by MAP and PLANCK.

1 AD condensate and B-ball decay

The quantum fluctuations of the inflaton field give rise to fluctuations of the energy density which are adiabatic. However, in the minimal supersymmetric standard model (MSSM), or its extensions, the inflaton is not the only fluctuating field. It is well known that the MSSM scalar field potential has many flat directions, along which a non-zero expectation value can form during inflation, leading to a condensate after inflation, the so-called Affleck-Dine (AD) condensate. When the Hubble rate becomes of the order of the curvature of the potential, given by the susy breaking mass $m_S$, the condensate starts to oscillate. At this stage B-violating terms are comparable to the mass term so that the condensate achieves a net baryonic charge. In the AD baryogenesis scenario the subsequent decay of the condensate will then generate the observable baryon number.

An important point is that the AD condensate is not stable but typically breaks up into non-topological solitons which carry baryon (and/or lepton) number and are therefore called B-balls (L-balls). For baryogenesis considerations, the most promising direction is the $d = 6$ ("$u^c u^c u^c$") direction, on

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which we shall focus on in the following. The formation of the B-balls takes place with an efficiency $f_B$, likely to be in the range 0.1 to 1. The properties of the B-balls depend on SUSY breaking and on the flat direction along which the AD condensate forms. We will consider SUSY breaking mediated to the observable sector by gravity. In this case the B-balls are unstable but long-lived, decaying well after the electroweak phase transition has taken place, with a natural order of magnitude for decay temperature $T_d \sim \mathcal{O}(1)$ GeV. This assumes a reheating temperature after inflation, $T_R$, is less than about $10^4$ GeV. Such a low value of $T_R$ is in fact necessary in D-term inflation models because the natural magnitude of the phase of the AD field, $\delta_{\text{CP}}$, is of the order of 1 in D-term inflation and along the $d=6$ direction AD baryogenesis implies that the baryon to entropy ratio is $\eta_B \sim \delta_{\text{CP}}(T_R/10^9 \text{ GeV})$, so that $T_R \simeq \mathcal{O}(1)$ GeV would be the most natural choice. It is significant that a low reheating temperature can naturally be achieved in D-term inflation models, as these have discrete symmetries in order to ensure the flatness of the inflaton potential which can simultaneously lead to a suppression of the reheating temperature.

2 Fluctuations of the AD field

The AD field $\Phi = \phi e^{i\theta}/\sqrt{2} \equiv (\phi_1 + i\phi_2)/\sqrt{2}$ is a complex field and, in the currently favoured D-term inflation models, is effectively massless during inflation. Therefore both its modulus and phase are subject to fluctuations with

$$\delta\phi_i(\vec{x}) = \sqrt{V} \int \frac{d^3k}{(2\pi)^3} e^{-i\vec{k} \cdot \vec{x}} \delta_{\vec{k}},$$

where $V$ is a normalizing volume and where the power spectrum is the same as for the inflaton field,

$$\frac{k^3|\delta_{\vec{k}}|^2}{2\pi^2} = \left(\frac{H_I}{2\pi}\right)^2,$$

where $H_I$ is the value of the Hubble parameter during inflation.

In D-term inflation models the phase of the AD field receives no order $H$ corrections after inflation and so its fluctuations are unsuppressed. The fluctuations of the phase correspond to fluctuations in the local baryon number density, or isocurvature fluctuations, while the fluctuations of the modulus give rise to adiabatic density fluctuations. For given background values $\bar{\theta}$ and $\bar{\phi}$, (with $\bar{\theta}$ naturally of the order of 1) one finds

$$\left(\frac{\delta\theta}{\tan(\theta)}\right)_k = \frac{H_I}{\tan(\theta)\bar{\phi}} = \frac{H_I k^{-3/2}}{\sqrt{2}\tan(\theta)\bar{\phi} I},$$
where $\phi_f$ is the value of $\phi$ when the perturbation leaves the horizon. The magnitude of the AD field $\Phi$ remains at the non-zero minimum of its potential until $H \simeq m_S$, after which the baryon asymmetry $n_B \propto \sin(\theta)$ forms. Thus the isocurvature fluctuation reads

$$\left( \frac{\delta n_B}{n_B} \right) \equiv \delta_B^{(i)} = \left( \frac{\delta\Phi}{\tan(\theta)} \right) \delta^i$$

The adiabatic fluctuations of the AD field may dominate over the inflaton fluctuations, with potentially adverse consequences for the scale invariance of the perturbation spectrum, thus imposing an upper bound on the amplitude of the AD field. In the simplest D-term inflation model, the inflaton is coupled to the matter fields $\psi_-$ and $\psi_+$ carrying opposite Fayet-Iliopoulos charges through a superpotential term $W = \kappa S \psi_- \psi_+$. At one loop level the inflaton potential reads

$$V(S) = V_0 + \frac{g^4 \xi^4}{32\pi^2} \ln \left( \frac{\kappa^2 S^2}{Q^2} \right); \quad V_0 = \frac{g^2 \xi^4}{2},$$

where $\xi$ is the Fayet-Iliopoulos term and $g$ the gauge coupling associated with it. COBE normalization fixes $\xi = 6.6 \times 10^{15}$ GeV. In addition, we must consider the contribution of the AD field to the adiabatic perturbation. During inflation, the potential of the $d = 6$ flat AD field is simply given by

$$V(\phi) = \frac{\lambda^2}{32M^6} \phi^{10}.$$ 

Taking both $S$ and $\phi$ to be slow rolling fields one finds that the adiabatic part of the invariant perturbation is given by

$$\zeta = \delta \rho / (\rho + p) = \frac{3}{4} \frac{\delta \rho_{S}^{(a)}}{\rho_{\gamma}} \propto \frac{V'(\phi) + V'(S)}{V'(\phi)^2 + V'(S)^2} \delta \phi .$$

Thus the field which dominates the spectral index of the perturbation will be that with the largest value of $V'$ and $V''$.

3 A lower bound on the isocurvature amplitude

The index of the power spectrum is given by $n = 1 + 2\eta - 6\epsilon$, where $\epsilon$ and $\eta$ are defined as

$$\epsilon = \frac{1}{2} M^2 \left( \frac{V'\phi}{V} \right)^2, \quad \eta = \frac{M^2 V''}{V}.$$
The present lower bounds imply that the condition that the spectral index is acceptably close to scale invariance essentially reduces to the condition that the spectral index is dominated by the inflaton, \( V'(\phi) < V'(S) \) and \( V''(\phi) < V''(S) \). The latter requirement turns out to be slightly more stringent and implies a lower bound on the AD condensate field \( \phi \) with \( \phi \lesssim \frac{0.48 (g/\lambda)^{1/4} (M \xi)^{1/2}}{\lambda^2} \).

As a consequence, there is a lower bound on the isocurvature fluctuation amplitude. Because the B-ball is essentially a squark condensate, in R-parity conserving models its decay produces both baryons and neutralinos (\( \chi \)), which we assume to be the lightest supersymmetric particles (LSPs), with \( n_\chi \simeq 3n_B \). This case is particularly interesting, as the simultaneous production of baryons and neutralinos may help to explain the remarkable similarity of the baryon and dark matter neutralino number densities. With B-ball decay temperatures \( T_d \sim \mathcal{O}(1) \) GeV, the decay products no longer thermalize completely and, so long as \( T_d \) is low enough that they do not annihilate after B-ball decay, retain the form of the original AD isocurvature fluctuation. Therefore in this scenario the cold dark matter particles can have both isocurvature and adiabatic density fluctuations, resulting in an enhancement of the isocurvature contribution relative to the baryonic case.

The total LSP number density is the sum of the thermal relic density \( n^{(th)}_\chi \) and the density \( n^{(B)}_\chi = 3f_B n_B \) originating from the B-ball decay. (Their relative importance depends on \( T_R \); for \( T_R \lesssim \mathcal{O}(1) \) GeV one would find \( n^{(th)}_\chi \simeq 0 \).) The isocurvature fluctuation imposed on the CMB photons is then found to be

\[
\frac{\delta \rho_\gamma^{(i)}}{\rho_\gamma} \simeq -\frac{4}{3} \left( 1 + \frac{m_B}{3f_B m_\chi} \right) \left( \frac{\Omega_\chi - \Omega^{(th)}_\chi}{\Omega_m} \right) \delta_B^{(i)} \equiv -\frac{4}{3} \omega \delta_B^{(i)},
\]

where \( \rho_\gamma^{(B)} \) is the LSP mass density from the B-ball, \( \Omega_m (\Omega_\chi) \) is total matter (LSP) density (in units of the critical density). Thus

\[
\beta \equiv \left( \frac{\delta \rho_\gamma^{(i)}}{\delta \rho_\gamma^{(a)}} \right)^2 = \frac{1}{9} \omega^2 \left( \frac{M^2 V'(S)}{V(S) \tan(\theta) \phi} \right)^2,
\]

It then follows that the lower limit on \( \beta \) is

\[
\beta \gtrsim 2.5 \times 10^{-2} g^{3/2} \lambda^{1/2} \omega^2 \tan(\theta)^{-2}.
\]

Thus significant isocurvature fluctuations are a definite prediction of the AD mechanism.
Isocurvature perturbations give rise to extra power at large angular scales but are damped at small angular scales. The amplitude of the rms mass fluctuations in an $8h^{-1}{\rm Mpc}^{-1}$ sphere, denoted as $\sigma_8$, is about an order of magnitude lower than in the adiabatic case. Hence COBE normalization alone is sufficient to set a tight limit on the relative strength of the isocurvature amplitude. Small isocurvature fluctuations are, however, beneficial, in the sense that they would improve the fit to the power spectrum in $\Omega_0 = 1$ CDM models with a cosmological constant (or $\Omega_0 = 1$, $\Lambda = 0$ CDM models with some hot dark matter).

Detecting isocurvature fluctuations at the level of $\beta \sim 10^{-4}$ should be quite realistic at MAP and Planck. Thus the forthcoming CMB experiments offer a test not only of the inflationary Universe but also of the B-ball variant of AD baryogenesis.

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References

1. E.W. Kolb and M.S. Turner, *The Early Universe* (Addison-Wesley, Reading MA, USA, 1990).
2. See e.g. M. Dine, L. Randall and S. Thomas, *Nucl. Phys.* B458 (1996) 291.
3. I. A. Affleck and M. Dine, *Nucl. Phys.* B249 (1985) 361.
4. For a recent review on baryogenesis, see e.g. A. Riotto and M. Trodden, [hep-ph/9901362](http://arxiv.org/abs/hep-ph/9901362).
5. A. Kusenko and M. Shaposhnikov, *Phys. Lett.* B418 (1998) 104.
6. K. Enqvist and J. McDonald, *Phys. Lett.* B425 (1998) 309.
7. A. Cohen, S. Coleman, H. Georgi and A. Manohar, *Nucl. Phys.* B272 (1986) 301.
8. A. Kusenko, *Phys. Lett.* B404 (1997) 285.
9. K. Enqvist and J. McDonald, *Nucl. Phys.* B538 (1999) 321.
10. K. Enqvist and J. McDonald, *Phys. Rev. Lett.* 81 (1998) 3071.
11. K. Enqvist and J. McDonald, *Phys. Lett.* B440 (1998) 59.
12. E. Halyo, *Phys. Lett.* B387 (1996) 43; P. Binetruy and G. Dvali, *Phys. Lett.* B388 (1996) 241.
13. C. Kolda and J. March-Russell, [hep-ph/9802358](http://arxiv.org/abs/hep-ph/9802358).
14. K. Enqvist and J. McDonald, [hep-ph/9811415](http://arxiv.org/abs/hep-ph/9811415).
15. R. Stomper, A. J. Banday and K. M. Gorski, *Ap. J.* **468** (1996) 8.
16. M. Kawasaki, N. Sugiyama and T. Yanagida, *Phys. Rev.* **D54** (1996) 2442; T. Kanazawa, M. Kawasaki, N. Sugiyama and T. Yanagida, astro-ph/9805102.
17. S. D. Burns, astro-ph/9711303.