Cosmology of Radiatively Generated Axion Scale

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Abstract

We discuss some cosmological aspects of supersymmetric axion models in which the axion scale is radiatively generated in terms of the weak scale and the Planck scale. They include thermal inflation, axions produced by the decay of oscillating Peccei-Quinn flatons, late time baryogenesis, and finally the possibility to raise up the cosmological upper bound on the axion scale in thermal inflation scenario.
1 Introduction

The axion solution to the strong CP problem [1] involves an intermediate scale for the spontaneous breaking of $U(1)_{PQ}$ which is far above the electroweak scale but still far below the Planck scale $M_P$. It would be interesting that this intermediate axion scale appears as a dynamical consequence of the electroweak scale and the Planck scale. This indeed happens [2] in some class of spontaneously broken supergravity (SUGRA) models. In this scheme, the early universe experiences the so-called thermal inflation and subsequently a period dominated by coherently oscillating flaton fields [3] which break $U(1)_{PQ}$ spontaneously. Here we wish to discuss some cosmological implications of such flatons [4]. They include axions produced by oscillating flatons, late time baryogenesis, and finally the possibility of raising up the cosmological upper bound on the axion scale in thermal inflation scenario.

2 Radiatively Generated Axion Scale

As an example, let us consider a variant of the model of Ref. [2] with superpotential

$$W = k\frac{\phi_1^n + 2\phi_2}{M_P^2} + h_N N\phi_1 + \cdots$$  \hspace{1cm} (1)$$

where $N$ is the right-handed neutrino superfield and the ellipsis denotes the part depending upon the fields in the supersymmetric standard model (SSM) sector. Here $\phi_1$ and $\phi_2$ correspond to flat directions when nonrenormalizable interactions and supersymmetry breaking effects are ignored. Including the radiative effects of the strong Yukawa coupling $h_N N\phi_1$, the soft mass-squared of $\phi_1$ becomes negative at scales around $F_a \simeq \langle \phi_1 \rangle$, and thereby driving $\phi_1$ to develop vacuum expectation value (VEV) at an intermediate scale. Neglecting the field $\phi_2$, the renormalization group improved scalar potential for $\phi_1$ is given by

$$V = V_0 - m_1^2|\phi_1|^2 + k^2|\phi_1|^{2n+4} M_P^{2n},$$  \hspace{1cm} (2)$$

where $m_1^2$ is positive and of order $m_{3/2}^2$, and $V_0$ is a constant of order $m_{3/2}^2 F_a^2$ which is introduced to make $V(\langle \phi_1 \rangle) = 0$. (Here we assume that SUSY is broken by hidden sector dynamics yielding the electroweak scale value of $m_{3/2} \simeq 10^2 \sim 10^3$ GeV.) Clearly the minimum of this scalar potential breaks $U(1)_{PQ}$ by

$$\langle \phi_1 \rangle \simeq F_a \simeq (m_{3/2} M_P^n)^{1/n+1},$$  \hspace{1cm} (3)$$
where we have ignored the coefficients of order unity. Note that the integer $n$ which determines the size of $F_a$ is determined by the Peccei-Quinn (PQ) charge assignment of the model.

### 3 Thermal Inflation

The above radiative mechanism generating the axion scale has substantial influence on the history of the universe [3]. At high temperature, $\phi_1$ receives a thermal mass $\delta m_1^2 \simeq |h_N|^2 T^2 \gg m_1^2$ leading to $\langle \phi_1 \rangle = 0$. This thermal mass is generated by right-handed neutrinos in the thermal bath. During this period, $\langle \phi_2 \rangle = 0$ also. When the temperature falls below $T \simeq V_0^{1/4}$, which is about $\sqrt{m_3/2} F_a$, the universe is dominated by the vacuum energy density $V_0$ and thus there appears a short period of thermal inflation. Below $T < m_1 \simeq m_3/2$, the effective mass of $\phi_1$ becomes negative and then $\phi_1$ develops an intermediate scale VEV given by Eq. (3). With $\langle \phi_1 \rangle \simeq F_a$, the other flaton field $\phi_2$ develops also a VEV of order $F_a$ through the $A$-type soft SUSY breaking term, $k A \phi_1^{n+2} \phi_2/M_p^n$, in the scalar potential. This procedure makes the thermal inflation end and subsequently the early universe experiences a period dominated by coherently oscillating PQ flaton field which corresponds to a linear combination of $\phi_1$ and $\phi_2$ orthogonal to the axion field.

After the period of coherent oscillation, the universe would be reheated by the decay products of the oscillating PQ flaton $\varphi$. The reheat temperature $T_{RH}$ is given by

$$T_{RH} \simeq g_{RH}^{-1/4} \sqrt{M_P \Gamma_\varphi} \simeq \left( \frac{N_{\text{eff}}}{10} \right)^{1/2} \left( \frac{10^{12}\text{GeV}}{F_a} \right) \left( \frac{300\text{GeV}}{M_\varphi} \right)^{3/2} \text{GeV},$$

where $g_{RH} \equiv g_*(T_{RH})$ counts the effective number of relativistic degree of freedom at $T_{RH}$, $M_\varphi$ denotes the flaton mass, and we parameterize the width of the flaton decay into thermalizable particles as $\Gamma_\varphi = N_{\text{eff}} M_\varphi^3 / 64\pi F_a^2$ with $N_{\text{eff}}$ presumed to be of order $10 \sim 10^2$.

### 4 Axion Energy Density at Nucleosynthesis

A feature peculiar to the PQ flaton is that its decay products include axions. Requiring that these axions do not spoil the big-bang nucleosynthesis (NS), we found [4]

$$\frac{B_a}{1 - B_a} \leq 0.24 \left( \frac{\delta N_\nu}{1.5} \right) \left( \frac{g_{RH}}{43/4} \right)^{1/3},$$

where $B_a$ denotes the effective branching ratio measuring how large fraction of flatons are converted into axions during the reheating, and $\delta N_\nu$ is the number of allowed extra neutrino
species, which is presumed to be in the range \(0.1 \sim 1.5\). This indicates that we need to tune the effective branching ratio \(B_a\) to be less than \(0.02 \sim 1/3\).

As is well known, generic axion models can be classified by two classes: hadronic axion models and Dine-Fischler-Srednicki-Zhitnitskii (DFSZ) axion models [1]. In hadronic axion models, all fields in the supersymmetric standard model (SSM) sector carry vanishing PQ charge. As a result, flaton couplings to SSM fields do vanish at tree level but they appear to be nonzero by radiative effects. Since flaton couplings to SSM fields are loop-suppressed, in hadronic type models, most of the oscillating flatons decay first into either axion pairs, or lighter flaton pairs, or flatino pairs, as long as the decays are kinematically allowed. Lighter flatons would experience similar decay modes, while flatinos decay into axion plus a lighter flatino. Then in the first round of reheating, most of flatons are converted into either axions or the lightest flatinos. The lightest flatinos will eventually decay into SSM particles. Then more than half of the original flatons are expected to be converted into axions, i.e. the effective branching ratio \(B_a \geq 1/2\), unless the flaton coupling to the lightest flatino is unusually large. This is in conflict with the NS limit (5) even for the most conservative choice \(\delta N_\nu = 1.5\), implying that hadronic axion models with radiatively generated axion scale have a difficulty with the big-bang NS unless the models are tuned to have an unusually large flaton coupling to the lightest flatino.

In DFSZ type models, flatons have tree level couplings to SSM fields which are of order \(M_{\text{SSM}}/F_a\) or \(M_{\text{SSM}}^2/F_a\) where \(M_{\text{SSM}}\) collectively denotes the mass parameters in the SSM, e.g. \(M_t, M_W, \mu, A\), and so on. If \(M_\varphi \gg M_{\text{SSM}}\), the reheating procedure would be similar to that of hadronic axion models, which is problematic. However if \(M_\varphi\) is comparable to \(M_{\text{SSM}}\), the NS limit (5) does not provide any meaningful restriction on DFSZ type models.

5 Late Time Baryogenesis

Thermal inflation driven by PQ flatons is expected to dilute away any pre-existing baryon asymmetry. However, PQ flatons themselves can produce baryon asymmetry after thermal inflation. A careful examination of the flaton couplings in DFSZ type models suggests that, among the decays into SSM particles, the decay channels to the top \((t)\) and/or stop \((\tilde{t})\) pairs are most important. Stops produced by the oscillating flatons would be in out-of-equilibrium and subsequently experience a \(B\) and \(CP\) violating decay to generate a baryon asymmetry provided
that the $B$-violating operator, e.g., $\chi''_{332} U_3^c D_3^c \tilde{D}_3^c$ and the corresponding complex trilinear soft-term are present [5]. Note that the PQ symmetry can be arranged so that dangerous lepton-number violating operators $LQD^c, LLE^c$ are forbidden for the proton stability.

In order for the baryon asymmetry (generated as above) not to be erased, the reheat temperature (4) has to be less than few GeVs [5]. This means that the above mechanism for baryogenesis can work only for $n = 2$ or $3$ [see Eqs. (3) and (4)]. The produced baryon asymmetry is [4]

$$\frac{\eta}{3 \times 10^{-10}} \simeq |\chi''_{332}|^2 \left( \frac{\text{arg}(Am_{1/2}^s)^2}{10^{-2}} \right) \left( \frac{10^{14}\text{GeV}}{F_a} \right) \left( \frac{M_\varphi}{300\text{GeV}} \right)^{1/2},$$

(6)

where $\text{arg}(Am_{1/2}^s)$ denotes the CP violating relative phase which is constrained to be less than about $10^{-2}$ for superparticle masses of order 100 GeV.

6 Raising up the Upper Bound on $F_a$

For a successful nucleosynthesis, we need $T_{RH} > 6$ MeV [6], implying that only $n = 1, 2,$ and 3 are allowed. As is well known, the axion scale is constrained by the consideration of the coherent axion energy density produced by an initial misalignment [1]. If there is no entropy production after the axion start to oscillate at around $T \simeq 1$ GeV, this lead to the usual bound: $F_a \leq 10^{12}$ GeV. When $n = 2$ or $3$, the corresponding axion scale $F_a \simeq (M_\varphi M_P^n)^{1/n+1}$ would exceed this bound. However in this case, the reheat temperature (4) goes below 1 GeV. Then the coherent axions may be significantly diluted by the entropy dumped from flaton decays, thereby allowing $F_a$ much bigger than $10^{12}$ GeV [7]. Taking into account of this dilution, we find [4]

$$\Omega_ah_{50}^2 \simeq \left( \frac{N_{\text{eff}}}{10} \right) \left( \frac{10^{12}\text{GeV}}{F_a} \right)^{0.44} \left( \frac{M_\varphi}{300\text{GeV}} \right)^{2.9} \left( \frac{\Lambda_{QCD}}{200\text{MeV}} \right)^{-1.9},$$

(7)

where we have used $\Gamma_\varphi \simeq N_{\text{eff}} M_\varphi^3/64\pi F_a^2$. The above result is valid only for $n \geq 2$. As we have anticipated, it shows that the case of $n = 2$ or $3$ with $F_a \simeq (M_\varphi M_P^n)^{1/n+1}$ yields a coherent axion energy density not exceeding the critical density although the corresponding $F_a$ exceeds $10^{12}$ GeV. Furthermore, in this case of $n = 2$ or $3$, axions can be a good dark matter candidate for an appropriate value of $M_\varphi$, which was not possible for $n = 1$.

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