On the dark radiation problem in the axiverse

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Abstract. String scenarios generically predict that we live in a so called axiverse: the Universe with about a hundred of light axion species which are decoupled from the Standard Model particles. However, the axions can couple to the inflaton which leads to their production after inflation. Then, these axions remain in the expanding Universe contributing to the dark radiation component, which is severely bounded from present cosmological data. We place a general constraint on the axion production rate and apply it to several variants of reasonable inflaton-to-axion couplings. The limit merely constrains the number of ultralight axions and the relative strength of inflaton-to-axion coupling. It is valid in both large and small field inflationary models irrespectively of the axion energy scales and masses. Thus, the limit is complementary to those associated with the Universe overclosure and axion isocurvature fluctuations. In particular, a hundred of axions is forbidden if inflaton universally couples to all the fields at reheating. In the case of gravitational sector being responsible for the reheating of the Universe (which is a natural option in all inflationary models with modified gravity), the axion production can be efficient. We find that in the Starobinsky $R^2$-inflation even a single axion (e.g. the standard QCD-axion) is in tension with the Planck data, making the model inconsistent with the axiverse. The general conclusion is that an inflation with inefficient reheating mechanism and low reheating temperature may be in tension with the presence of light scalars.

Keywords: axions, inflation, string theory and cosmology

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1 Introduction

In order to solve several problems of the hot Big Bang cosmology, the inflationary stage of the Universe evolution has been proposed [1–3]. The simplest way to organize the close-to-exponential expansion of the Universe is to exploit the scalar field (inflaton) slowly rolling towards the minimum of the potential [4, 5]. Recent data from the Planck satellite put such strong constraints on the inflaton potential that the simplest quartic and quadratic cases turn out to be excluded [6]. The best agreement with the data is still exhibited by the large field inflation provided by the exponentially flat potential. Such a potential naturally arises in different models with the modified gravitational sector [1, 2, 7, 8] as well as in the supergravity framework [9–11]. All these models, being non-renormalizable, require an ultraviolet completion which is often thought to be a string theory.

However, the common prediction of many string scenarios is the existence of a plenty of light scalars (axions) arising as zero modes of antisymmetric gauge fields on the compactified dimensions [12]. They inherit a perturbative shift symmetry violated by instanton contributions [13, 14]. One of these scalars can play a role of QCD axion explaining the zero value of the QCD \( \theta \)-angle. The number of light axions is determined by the topology of extra dimensions and is likely to be about hundred [15].

Many potentially observable consequences of the string axiverse were discussed in literature [15]. Here we explore whether many light scalars are compatible with inflationary models. String axions can contribute to dark matter component, see e.g. [16, 17]. In order to avoid overproduction and suppress isocurvature fluctuations, one constrains the axion masses, couplings and initial conditions at the inflationary stage. At the same time, light scalars produced at reheating can contribute to the dark radiation, which implies an additional bound on the model parameters. Dark radiation amount at Big Bang Nucleosynthesis (BBN) and recombination is now strictly bounded by the Planck data [18]. In generic string scenarios, the number of effective relativistic degrees of freedom at these epochs is predicted to be much larger than the Planck results [18] allow, the discrepancy is mostly due to the contribution of moduli decays [19–21]. One can suppose, however, that the supersymmetry breaking scale, as well as the masses of all moduli, are larger than the Hubble parameter at inflation. In this case, the moduli are not produced and the axions are created only by the inflaton decay.\(^1\) We show that in this case, if the inflaton couples to SM particles only through the gravity (i.e. via suppressed by Planck mass operators, which is natural, e.g. for

\(^1\)Notice that, actually, inflation might be driven by one of the moduli fields.
$F(R)$ inflationary models), then even one light scalar is in tension with the recent Planck data. Thus, in the axiverse the inflaton must couple to matter much stronger. In that case, we put a lower bound on the reheating temperature which makes the existence of hundreds of light scalars consistent with the Planck data.

2 Reheating in the axiverse

In each model, the inflationary stage must be followed by the reheating process eventually populating the Universe with the hot plasma of SM particles. Therefore, the energy of the inflaton field must be somehow transformed into the usual matter. In general, during this process axion-like particles can well be produced. Let the production rate of each axion be $\Gamma_a$ while that of the SM particles be $\Gamma_{SM}$. Then, the number of additional degrees of freedom at BBN and recombination written traditionally as the number of additional neutrino components is

$$\Delta N_{\text{eff}} = N \frac{g_{\text{reh}}}{g_\nu} \left( \frac{g_{\text{BBN}}}{g_{\text{reh}}} \right)^{4/3} \frac{\Gamma_a}{\Gamma_{SM}}.$$  \hspace{1cm} (2.1)

Here $N$ is the number of axion species, $g_\nu = 2 \cdot 7/8$, $g_{\text{reh}}$ and $g_{\text{BBN}}$ are the effective number of relativistic degrees of freedom at reheating and nucleosynthesis stages, respectively. Their SM values are $g_{\text{reh}} = 106.75$ and $g_{\text{BBN}} = 10.75$. According to the latest Planck data, within the concordance $\Lambda$CDM cosmological model the effective number of relativistic species $N_{\text{eff}}$ is bounded as  

$$N_{\text{eff}} = 3.15 \pm 0.23.$$  \hspace{1cm} (2.2)

This parameter is expected to be measured in future CMB polarization experiments with much higher accuracy [23, 24].

One can observe from (2.1) that if the reheating is due to some universal mechanism (i.e. inflaton decays to all scalars including the SM Higgs with the comparable rates $\Gamma_a \sim \Gamma_{SM}/4$) then $\Delta N_{\text{eff}} \sim 0.7N$ which is clearly incompatible with bound (2.2) in the axiverse. This is the main finding of the present paper.

The bound (2.1), (2.2) constrains axion coupling to inflaton and the number of axions ultrarelativistic at BBN and recombination. It is applicable to both large- and small-field inflation with any mass pattern in the axion sector, as far as the axions remain ultrarelativistic. Thus, the obtained constraint is complementary to another bounds on the axion energy scales, masses and on the energy scale of inflation. These bounds are inferred from the Universe overclosure argument and absence of the axion isocurvature fluctuations [16, 22].

Formula (2.1) is exact if the axion branching ratio $\Gamma_a/\Gamma_{SM}$ is constant in time at the reheating epoch. Otherwise it is corrected by a numerical factor of order one, which is the case if the inflaton couplings to the axion and SM-particles are of the different nature, e.g. provided by the operators of different dimensions. Then $\Gamma_{SM}$ in eq. (2.1) defines the age of the Universe at reheating, $t_U \sim 1/\Gamma_{SM}$ and $\Gamma_a/\Gamma_{SM} \sim \Gamma_a t_U \ll 1$ refers to the fraction of inflaton energy transferred to axions.

The case when moduli decay is responsible for the reheating provides an example of the universal reheating mechanisms, which has been studied in refs. [17, 19–21]. Below we consider several other examples of models with universal reheating mechanisms relevant to the problem with axiverse.
2.1 The Starobinsky model

The Starobinsky model of inflation historically was the first successful model suggested for the exponential stage of the Universe expansion [1, 2]. However, it still provides the predictions which are in a perfect agreement with the present cosmological data [18]. In a Jordan frame, the model is described by the following action:

$$ S = -\frac{M_P^2}{2} \int \sqrt{-g} \, d^4x \left[ R - \frac{R^2}{6m^2} \right] + S_{\text{matter}}. $$

(2.3)

Here the reduced Planck mass is $M_P = M_{Pl}/\sqrt{8\pi} = 2.4 \times 10^{18}$ GeV and $S_{\text{matter}}$ denotes the action for all matter fields in the theory. The value of $m$, which is actually mass of an additional scalar degree of freedom (scalaron) responsible for inflation, is determined by the amplitude of scalar perturbations as $m = 1.3 \times 10^{-5} M_P$ [28].

After the Weyl transformation of the metric to the Einstein frame,

$$ g_{\mu\nu} \rightarrow e^{\sqrt{2/3}\phi/M_P} g_{\mu\nu}, $$

(2.4)

action (2.3) takes the form [25]

$$ S = \int \sqrt{-g} \, d^4x \left[ -\frac{M_P^2}{2} R + \frac{1}{2} \partial_\mu \phi \partial^\mu \phi - V(\phi) \right] + \tilde{S}_{\text{matter}}, $$

(2.5)

$$ V(\phi) = \frac{3m^2M_P^2}{4} \left( 1 - e^{-\sqrt{2/3}\phi/M_P} \right)^2. $$

(2.6)

Here $\tilde{S}_{\text{matter}}$ is the Weyl transformed action of the matter fields. Thus, conformal non-invariance in the matter sector naturally implies the interaction between scalaron $\phi$ and all other particles. Fermions and vector fields are Weyl invariant at the tree level in the high energy limit. Therefore, a key role in the process of reheating is played by the light scalars: SM Higgs boson and axions in the discussed framework (see ref. [31] for dilaton).

Axions $a_i$, $i = 1, \ldots, N$, unlike the Higgs, due to the perturbative shift symmetry [14] can not be non-minimally coupled to the gravity via terms $R\phi^2$. Thus, their kinetic terms are canonical in the Jordan frame (2.3), yielding the universal couplings to the inflaton. The decay rates of scalaron to Higgs pair and to axion pair are [26, 32]

$$ \Gamma_{\text{SM}} = \frac{m^3(1 + 6\xi_h)^2}{48\pi M_P^2}, \quad \Gamma_{\alpha} = \frac{m^3}{192\pi M_P^2}, $$

(2.7)

respectively, with the SM Higgs boson possibly non-minimally coupled to gravity via lagrangian term $L = \xi_h R h^2/2$. One observes that the decay rates to the axions and to the SM particles (Higgs bosons) are comparable.

The ultrarelativistic axions produced at reheating remain in the late Universe contributing to the energy density and pressure at BBN and recombination as additional

$$ \Delta N_{\text{eff}} = 0.7N, $$

(2.8)

neutrino flavors (we put (2.7) with $\xi_h = 0$ into (2.1)). Hence, even one additional light scalar (for example, standard QCD axion) is already in tension with the Planck bound (2.2) in the Starobinsky model. The similar result has been obtained in [27] for generic $F(R)$-gravity models.

To resolve the situation one can add some new $N_s$ scalars to the matter content of the SM. Then the r.h.s. of (2.8) gets suppressed by a factor $4/(4 + N_s)$, which makes the Starobinsky model in the axiverse populated by $N \sim 100$ axions consistent with the present cosmological bounds (2.2) if $N_s \sim 200$ scalars are added.
2.2 Inflation with non-minimal kinetic terms

Another way to obtain the favored by the Planck results exponentially flat potentials in a natural way is connected with the modification of the kinetic term of the inflaton. This idea is widely discussed in the context of supergravity [9, 10, 29, 37] where non-trivial kinetic terms come from the Kahler potential. In particular, it was realized as $\alpha$-attractors in refs. [33, 34], where the inflationary region corresponds to the pole in the kinetic term of the inflaton. In all these models the inflaton may be canonically normalized upon an appropriate field transformation. In general, one can expect that the kinetic terms of additional scalars (axions) are also non-minimal. Thus, the action consistent with the shift symmetry of axions $a_i$ reads

$$S = \int d^4x \sqrt{-g} \left( -\frac{M_P^2}{2} R + \frac{(\partial_{\mu}\phi)^2}{2} + \sum_{i=1}^{N} f_i(\phi) \left( \frac{\partial_{\mu}a_i}{2} \right)^2 - V(\phi) \right), \quad (2.9)$$

where $f_i(\phi)$ are some functions of the inflaton field.

Similarly one may expect non-renormalizable couplings to the SM fields yielding the most relevant for reheating two-body decays of inflaton,

$$S_{\text{int}} = \int d^4x \sqrt{-g} \left( y(\phi)|D\mu H|^2 - \frac{1}{4} g_{ij}(\phi) F_{\mu\nu,j} F^{\mu\nu} + z_i(\phi) \bar{\psi}_i \gamma^\mu D_\mu \psi_i \right), \quad (2.10)$$

Terms in the last set in (2.10) are proportional to the fermion masses through the equation of motion and hence their role in the reheating is negligible.

Near the minimum of $V(\phi)$ (we take it to be zero) one may anticipate the expansions

$$f_i(\phi) = 1 + \frac{\beta_i}{\Lambda} \phi + \gamma_i \frac{\phi^2}{\Lambda^2} + \ldots, \quad (2.11)$$

$$g_i(\phi) = 1 + \frac{\delta_i}{\Lambda} \phi + \ldots, \quad y(\phi) = 1 + \frac{\gamma}{\Lambda} \phi + \ldots.$$

From the point of view of the effective theory considered after inflation such terms suppressed by cutoff scale $\Lambda < M_P$ are naturally expected with $\beta_i, \gamma, \delta \sim 1$ from quantum corrections since the potential of the inflaton (as well as the gravity) is non-renormalizable. However, during inflation $f_i(\phi) \approx 1$ because of the approximate shift symmetry of the inflaton field with Planck-favored plateau-like potential providing no significant physical effects. But after inflation the expansions (2.11) start to work leading to the inflaton decay. For example, with $\gamma \neq 0$ the reheating process can go through the decay of the inflaton to the Higgs bosons. If the coupling to gauge bosons is suppressed for some reason we are left with the similar case as in the Starobinsky model where the dark radiation production is highly efficient (2.8). This model is cosmologically forbidden.

If the inflaton couples to the kinetic terms of all the matter fields in the model with $\gamma, \delta_i \sim 1$, then all the bosons will be produced roughly at equal amounts. In this case one can get an estimate for the axion contribution to the effective relativistic degrees $\Delta N_{\text{eff}}$ by putting $\Gamma_a/\Gamma_{\text{SM}} \sim 1/30$ in (2.1). In this way one obtains $\Delta N_{\text{eff}} \sim N/10$, which is certainly cosmologically forbidden for $N \sim 100$ given the constraint (2.2).

The problem of axion overproduction may be avoided if the inflaton couples to the SM particles stronger than to the axions. The former may be parametrized by means of the time of reheating (as we discuss at the beginning of section 2) or the reheating temperature $T_{\text{reh}}$, i.e. the temperature of the SM plasma at the moment when a half of the total energy density is already in the form of radiation.
Then, if the inflaton mass is $m$, one obtains from eqs. (2.9), (2.11),

$$\Gamma_{SM} \simeq 3 \frac{T_{reh}^2}{\sqrt{g_{reh} M_P}}, \quad \Gamma_a = \frac{\beta^2 m^3}{128 \pi \Lambda^2},$$  \hspace{1cm} (2.12)

Substituting (2.12) into the equation (2.1) we obtain for the amount of dark radiation:

$$\Delta N_{\text{eff}} = 0 \cdot 024 N^2 \frac{\beta^2 m^3 M_P}{\Lambda^2 T_{reh}^2}.$$  \hspace{1cm} (2.13)

For the reheating temperature high enough one can see that $N \sim 100$ may still be allowed by the Planck constraints. On the contrary, inefficient reheating with low $T_{reh}$ can easily throw the model out of the viable range (2.2).

Note in passing that the terms of the first order in the expansion (2.11) may be forbidden due to some symmetry (the simplest one is $Z_2, \phi \rightarrow -\phi$). In this case, the production of axions would be inefficient if the Universe is reheated due to the other inflaton couplings to matter provided by a lower order operators.

We discuss this case in more details in the next section using the inflaton non-minimally coupled to gravity as a realistic example.

### 2.3 Inflation driven by the scalar field non-minimally coupled to gravity

Although inflation models with reasonable (e.g. renormalizable without gravity) power-law potentials are disfavoured by the Planck data, switching on the non-minimal coupling of the inflaton to gravity can provide with the flat potential suppressing the tensor modes. Models of such type are widely discussed in the literature (see e.g., [33, 35, 36]) and include the Higgs inflation [8]. The production of light scalars determined by the linear coupling of the inflaton to the curvature was studied in [27]. Here we consider the class of models with quadratic coupling to the Ricci scalar. The action for the inflaton field $\phi$ reads (here we neglect the possible mass term for the inflaton at large field values),

$$S = \int d^4x \sqrt{-g} \left( -\frac{M_P^2}{2} \phi^2 R + \frac{\partial^2 \chi}{2} - \frac{\lambda}{4} \phi^4 \right).$$  \hspace{1cm} (2.14)

One can get rid of the non-minimal coupling by making use of the metric redefinition

$$\hat{g}_{\mu\nu} = \Omega^2 g_{\mu\nu}, \quad \Omega^2 = 1 + \frac{\xi \phi^2}{M_P^2}.$$  \hspace{1cm} (2.15)

After that, the action in the Einstein frame takes the form

$$S_{E} = \int d^4x \sqrt{-\hat{g}} \left\{ -\frac{M_P^2}{2} \hat{R} + \frac{(\partial^\mu \chi)^2}{2} - U(\chi) \right\},$$  \hspace{1cm} (2.16)

where canonically normalized field $\chi$ is defined by

$$\frac{d\chi}{d\phi} = \frac{\sqrt{\Omega^4 + 6 \xi^2 \phi^2 / M_P^2}}{\Omega^2}, \quad \text{and} \quad U(\chi) = \frac{1}{\Omega^4(\chi)} \frac{\lambda}{4} \phi^4(\chi).$$  \hspace{1cm} (2.17)

The kinetic term of axion gets coupled to the inflaton in the Einstein frame:

$$L_a = \frac{1}{2} \Omega^2 (\partial_{\mu} a)^2 = \frac{1}{2} \left( 1 + \frac{\xi \phi^2}{M_P^2} \right) (\partial_{\mu} a)^2.$$  \hspace{1cm} (2.18)
At first sight, this coupling is quadratic in inflaton and seems to be strongly suppressed by squared Planck scale. However, at and some time after inflationary epoch the inflaton field takes large values, \( \phi \sim M_P \). This makes the second term in parenthesis in (2.18) important to the extent which depends on the value of nonminimal coupling \( \xi \). Taking into account this fact we study two different cases for the value of \( \xi \) which finally yield different results.

**Large non-minimal coupling, \( \xi \gg 1 \).** This case includes the model of Higgs inflation [8]. In this limit for large field values of \( \phi>M_P/\sqrt{\xi} \) one obtains from (2.17)

\[
\phi \simeq \frac{M_P}{\sqrt{\xi}} \exp \left( \frac{\chi}{\sqrt{6}M_P} \right), \quad U(\chi) = \frac{\lambda M_P^4}{4\xi^2} \left( 1 + \exp \left( -\frac{2\chi}{\sqrt{6}M_P} \right) \right)^{-2}.
\] (2.19)

Inflation in models of this type is followed by harmonic oscillations with frequency \( \omega = \sqrt{\lambda/3M_P/\xi} \): the scalar potential is effectively quadratic while the amplitude of \( \chi \) is large enough,

\[
\chi \gg M_P/\xi.
\] (2.20)

The Universe is expanding as at the stage of matter domination: \( a \propto t^{2/3} \). In the original Higgs inflation [8] the reheating happens at this stage due to decays of Higgs to the SM particles [30]. The interaction lagrangian between the inflaton and any additional scalar \( a \) coming from the Weyl transformation (2.15) takes the same form as in the Starobinsky model of section 2.1:

\[
L_{\text{int}} = \frac{\chi}{\sqrt{6}M_P} \partial_\mu a \partial^\mu a.
\] (2.21)

If the reheating happens at this stage (eqs. (2.19), (2.21), (2.20)) then the decay rates of the inflaton are actually given by eq. (2.12). Here \( T_{\text{reh}} \) is the temperature of the SM plasma at the moment of equality between energy densities of radiation and inflaton excitations. For the Higgs inflation \( T_{\text{reh}} \simeq 6 \times 10^{13} \text{GeV} \) [30]. Let us evaluate the amount of dark radiation given the reference values of the Higgs inflation:

\[
\Delta N_{\text{eff}} \simeq 5.6 \times 10^{-8} \left( \frac{\omega}{1.3 \times 10^{-9} M_P} \right)^3 \left( \frac{6 \times 10^{13}}{T_{\text{reh}}} \right)^2.
\] (2.22)

One can observe that the axiverse with \( N \sim 10^2 \) is safe from the overproduction of dark radiation in models with high enough reheating temperature. Such models require non-gravitational couplings between the inflaton and SM particles, like gauge and Yukawa interactions between the inflaton (Higgs) and the SM fields in the example of Higgs inflation.

A side remark concerns models with the axions coupled to the SM particles in plasma that provide with more efficient mechanisms of the dark radiation production. The QCD axion with the decay constant in the range \( f_a \sim 10^{10} - 10^{12} \text{GeV} \) is a realistic example. Such axion couples to the SM particles via dimension-5 operators suppressed by \( \Lambda = f_a \) rather than \( \Lambda = M_P \). Therefore, it thermalizes in the SM plasma of temperature above \( 10^9 \text{GeV} \) which is the case for the Higgs inflation. Thus, the amount of dark radiation in this case is defined by the axion number density in thermal equilibrium. This leads to the value [38]

\[
\Delta N_{\text{eff}} = 0.026,
\] (2.23)

which can be measured in future CMB polarization experiments [23, 24]. At the same time, the thermal production of string axions with \( f_a \sim 10^{16} \text{GeV} \) is still inefficient.
It is worth noting that the reference value of the reheating temperature in (2.22) corresponds to the final amplitude of the inflaton oscillations of order $\chi \sim M_P/\xi$ [30]. In other words, for a quartic inflation with non-minimal coupling to gravity, the reheating temperature cannot be lower than the reference value in (2.22), if the system is still at the effective matter domination stage provided by (2.20). Hence eq. (2.22) imposes a kind of upper limit on the impact of axions in the model with efficient reheating.

If the reheating is less efficient than in the Higgs inflation, the evolution comes to the second stage. There the inflaton amplitude drops down to the value $\chi \sim M_P/\xi$ before the reheating started, and the potential and interactions of the canonically normalized inflaton field $\chi$ change the forms:

$$U(\chi) = \frac{\lambda}{4} \chi^4, \quad L_{\text{int}} \sim \frac{\chi^2}{M_P^2} \partial_\mu a \partial^\mu a.$$  \hfill (2.24)

After this moment the axion production becomes inefficient because, instead of the inflaton decay, we deal with the inflaton scattering suppressed by the squared Planck mass. Moreover, the scattering rate dilutes as $1/a^3$ due to the Universe expansion providing with negligible overall axion production.

Therefore at this stage, the axion production actually terminates providing the overall impact of axions to be of order (2.22) calculated for the reference value of the reheating temperature $T_{\text{reh}} \simeq 6 \times 10^{13}$ GeV. One can see that for all reasonable parameter choices we are left with a negligible amount of axion dark radiation. Note that the real reheating, that is a production of the SM plasma, may happen much later, than the reference temperature indicates, but it does not change this estimate in the slightest. Inflaton couplings to the SM particles must be of another form than the kinetic one of eq. (2.24), the latter is not sufficient for successful reheating.

**Small non-minimal coupling** $\xi \ll 1$. Here the change of variables (2.17) may be simplified:

$$\frac{d\chi}{d\phi} = \frac{1}{\sqrt{1 + \xi \phi^2/M_P^2}}, \quad \phi = \frac{M_P}{\sqrt{\xi}} \sinh \left(\frac{\sqrt{\xi} \chi}{M_P}\right).$$  \hfill (2.25)

The potential (2.14) transforms to

$$U(\chi) = \frac{\lambda M_P^4}{4\xi^2} \text{th}^4 \left(\frac{\sqrt{\xi} \chi}{M_P}\right).$$  \hfill (2.26)

Near the minimum, conformal factor (2.15) can be approximated as $\Omega^2 = 1 + \xi \chi^2/M_P^2$ providing the leading interaction term with scalars to be

$$L_{\text{int}} = \xi \frac{\chi^2}{M_P^2} \partial_\mu a \partial^\mu a.$$  \hfill (2.27)

Potential (2.26) is symmetric with respect to $\chi \rightarrow -\chi$ so no linear terms are expected. The production of axions in that case is inefficient due to the $M_P^2$ suppression of interaction (2.27) in accord with the expectations we discussed right below eq. (2.18).

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2We do not consider values of $\xi$ much larger than those of the Higgs inflation ($\xi \sim 10^5$) because it would lead to strong coupling for the inflaton self-coupling, $\lambda \gtrsim 1$. 

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3 Conclusions

In this paper, we investigate the validity of inflationary models in the string axiverse. Many light scalars can be produced at reheating and later contribute to the dark radiation component of the Universe which is strictly bounded by the recent Planck data. We found the general conditions for the efficient production of the light scalars at the Universe reheating. Namely, if the inflaton decays to two axions via the dimension-5 Planck-scale suppressed operators then the amount of the dark radiation is controlled by the reheating temperature. For example, inflationary models with reheating via Planck suppressed couplings of the inflaton to the SM particles (which seems to be common in the supergravity framework) predict too much dark radiation making them inconsistent with the cosmological observations. We should stress that our results are directly applicable not only in the string framework but for any light scalars (Nambu-Goldstone bosons, dilaton) which may appear in a concrete cosmological model.

However, there are two ways how to make the inflation consistent with the presence of extra light scalars. The first way is to provide the additional couplings between the inflaton and SM fields which are not suppressed by the Planck mass. It would raise the reheating temperature leaving no time for the axions production after inflation. Another way around is to deal with models possessing a symmetry which either forbids or strongly suppresses the inflaton decay to axions (the symmetry must not prevent the successful reheating, of course).

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