Dark Radiation in LARGE Volume Models

Michele Cicoli∗∗, Joseph P. Conlon†, Fernando Quevedo∗∗∗†

∗ ICTP, Strada Costiera 11, Trieste 34014, Italy
∗∗ INFN, Sezione di Trieste, Italy
† Rudolf Peierls Centre for Theoretical Physics, 1 Keble Road, Oxford, OX1 3NP, UK
∗∗∗ DAMTP, Centre for Mathematical Sciences, Wilberforce Road, Cambridge, CB3 0WA

We consider reheating driven by volume modulus decays in the LARGE Volume Scenario. Such reheating always generates non-zero dark radiation through the decays to the axion partner, while the only competitive visible sector decays are Higgs pairs via the Giudice-Masiero term. In the framework of sequestered models where the cosmological moduli problem is absent, the simplest model with a shift-symmetric Higgs sector generates $1.56 \leq \Delta N_{eff} \leq 1.74$. For more general cases, the known experimental bounds on $\Delta N_{eff}$ strongly constrain the parameters and matter content of the models.

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I. INTRODUCTION

The cosmological Standard Model (SM) starts with a period of inflation. During this period, the energy of the universe is dominated by the vacuum energy of a slowly rolling scalar field. At some point inflation ends and, irrespective of the overall particle spectrum or number of hidden sectors, the energy has to be transferred predominantly into thermal relativistic SM degrees of freedom via a process of reheating.

Constraints on this are measured via $N_{eff}$, the effective number of neutrino species. $N_{eff}$ is measured both at BBN and CMB times, and in practice measures the fraction of the total energy density that lies in the thermal photon plasma. At CMB temperatures $N_{eff}$ is determined in terms of the total energy by

$$\rho_{total} = \rho_{\gamma} \left( 1 + \frac{7}{8} \left( \frac{4}{11} \right)^{4/3} N_{eff} \right).$$

In the SM $N_{eff,\text{BBN}} = 3$ and $N_{eff,\text{CMB}} = 3.04$, due to partial reheating of the neutrinos from $e^+e^-$ annihilation. The presence of additional dark radiation, decoupled from the SM and relativistic at both BBN and CMB temperatures, leads to $\Delta N_{eff} \equiv N_{eff} - N_{eff,\text{SM}} > 0$.

Observations show a mild but consistent preference for $\Delta N_{eff} > 0$. At CMB times WMAP, ACT and SPT report $N_{eff} = 4.34_{-0.86}^{+0.88}$. At BBN times an excess has also been reported but the evidence depends on the relic Helium abundance. Based on D/H, reports $N_{eff} = 3.9 \pm 0.44$. A recent general overview is [4].

As the inflationary universe is vacuum energy dominated, dark radiation must arise during or after reheating. In the context of string models of the early universe, there are two main challenges in understanding reheating. The first, the cosmological moduli problem (CMP) is to understand how re-heating can occur at all, and the second is to ensure that it is primarily the SM that is reheated.

We recall first the CMP [5]. String theory contains many moduli associated to the complicated Calabi-Yau geometry. Moduli are typically Planck-coupled scalars which are expected to obtain vevs during inflation, leading to post-inflationary production of moduli through the vacuum misalignment mechanism. Moduli oscillate coherently as matter, redshift slowly and come to dominate the energy density of the universe. As Planck-coupled fields, their characteristic decay rate is $\Gamma \sim \frac{m_3^3}{16\pi m_p^2}$. A reheating temperature $T \gtrsim \mathcal{O}(1)$ MeV, necessary for BBN, then requires $m_3 \gtrsim 30$ TeV. For ‘generic’ models, $m_3 \sim m_{3/2} \sim M_{soft}$, leading to a tension with supersymmetric solutions of the hierarchy problem.

There is a cognate problem with respect to decays to gravitini [6]. Even if $m_3 \gg M_{soft}$, provided $m_3 \gtrsim 2m_{3/2}$ the decay mode $\phi \to \psi_{3/2}\bar{\psi}_{3/2}$ is kinematically open. This decay mode is problematic as for $m_{3/2} \lesssim 30$ TeV the gravitino decays could affect the successful BBN predictions.

The second problem is to ensure that only the SM is reheated. String theory generally contains many extra sectors in addition to the SM. These include additional hidden gauge and matter sectors, as well as light axion-like particles. Excessive branching ratios to these hidden sectors would lead to an overproduction of dark matter or $\Delta N_{eff} \gg 1$ and a failure of the BBN predictions. There is also a practical difficulty. Calabi-Yaus have many - easily $\mathcal{O}(100)$ - moduli which in generic models of moduli stabilisation have parametrically similar masses and lifetimes. A study of reheating then requires a coupled analysis of all moduli and their decay modes. Such an analysis is not only impractical, it is also highly sensitive to the post-inflationary initial conditions as the relative energy densities in each modulus field depends on the magnitude of the initial modulus misalignment and its non-perturbative production rate at pre-heating.
II. DARK RADIATION IN LVS VACUA

We shall analyse reheating within the LARGE Volume Scenario (LVS) [7], in which the volume is stabilised at exponentially large values. For definiteness, we take a canonical strong Swiss-cheese form for the volume

\[ V = \frac{1}{\lambda} \left( r_b^{3/2} - \sum_i \lambda_i r_{s,i}^{3/2} \right), \tag{2} \]

but our conclusions would hold also for more general LVS scenarios. For recent explicit constructions of LVS models see \[8\].

LVS has some particular advantages rendering this analysis meaningful. First, the moduli masses take a distinctive hierarchy, with the volume modulus by far the lightest of the non-axionic fields (a are axionic partners)

\[ M_P \equiv 2.4 \times 10^{18} \text{GeV}, \]
\[ M_{\text{string}} \sim M_P/\sqrt{V}, \]
\[ m_{\tau,i} \sim m_{a,i} \sim M_P \ln V/V, \]
\[ m_{3/2} \sim m_U \sim m_S \sim M_P/\sqrt{V}, \]
\[ m_{\tau_b} \sim M_P/\sqrt{3}/2, \]
\[ m_{a_b} \lesssim M_P e^{-2\pi V^{2/3}} \sim 0. \]

Of particular importance for this paper is the volume axion \(a_b\). As it can only obtain mass by effects non-perturbative in the volume, it is effectively massless on all cosmological scales. Its presence and masslessness are universal implications of LVS [7, 12], and so it is important to extract physical implications of \(a_b\).

Second, as \(V \gg 1\) and \(\tau \propto m^{-3}\), the volume modulus \(\tau\) outlives all other moduli by factors polynomial in volume. With radiation redshifting as \(a^{-4}\) and matter as \(a^{-3}\), this implies any radiation produced by the decays of other moduli will be highly diluted by the time the volume modulus decays. In LVS it is therefore reasonable to expect that the universe goes through a period when its energy density is dominated by coherent oscillations of the volume modulus \(\tau\), with reheating driven by the decays of \(\tau\).

What is \(m_{\tau_b}\)? Naively, in gravity mediation \(M_{\text{soft}} \sim m_{3/2}\). However the no-scale susy breaking structure of LVS leads to large cancellations in the soft terms, both at tree level and loop level (anomaly mediation \[9\], see however \[11\] 22). In the case where the SM is realised by branes at a singularity \[10\], gaugino masses are suppressed to a scale \(M_{1/2} \sim M_P/\sqrt{V}\). For soft scalar masses, the level of calculation in \[10\] could only establish suppression of soft scalar masses to a scale \(m_\phi^2 \sim M_P^2/\sqrt{V}\).

Going beyond this requires \(\alpha'\) corrections to the Kähler potential that have not been computed (in particular, the \(\zeta(3)\mathcal{R}^4\) correction in the matter Kähler metric). We shall assume the level of volume-sequestering is the same for both gaugino and scalar masses, and so (for an analysis of the non-sequestered case see \[12\])

\[ M_{\text{soft}} \sim \frac{M_P}{\sqrt{V}} \ll m_{\tau_b} \sim \frac{M_P}{\sqrt{3}/2}. \]

Setting \(M_{\text{soft}} \sim 1\) TeV then leads to \(m_{\tau_b} \sim 3 \times 10^6\) GeV, comfortably heavy enough to avoid the CMP.

A. Reheating via volume modulus decays

To study reheating in LVS we therefore focus on perturbative decays of the volume modulus (ref. \[14\] focused instead on decays of the inflaton). We require the relative fractions of radiation entering hidden and SM degrees of freedom, \(f_{\text{hidden}}\) and \(f_{\text{SM}} \equiv 1 - f_{\text{hidden}}\). In turn this requires the couplings of the volume modulus - the overall breathing mode - to both visible and hidden degrees of freedom (some moduli-matter couplings have already been computed in \[15\]).

1. Decays to volume axions

The basic theory we start with is

\[ K = -3 \ln(T_b + \bar{T}_b). \tag{3} \]

Although LVS is a complicated many-modulus model, eq. \(\text{(3)}\) captures the couplings of the volume modulus: the small cycle moduli are much smaller and much heavier than the overall volume modulus and any mixing is volume-suppressed. This gives

\[ \mathcal{L} = \frac{3}{4T_b} \partial_\mu \tau_b \partial^\mu \tau_b + \frac{3}{4T_b} \partial_\mu a_b \partial^\mu a_b. \tag{4} \]

We canonically normalise the volume modulus by writing \(\Phi = \sqrt{\frac{2}{3}} \ln \tau_b\), giving

\[ \mathcal{L} = \frac{1}{2} \partial_\mu \Phi \partial^\mu \Phi + \frac{1}{2} \left( \frac{3}{2} \exp(-2\sqrt{\frac{2}{3}} \Phi) \right) \partial_\mu a_b \partial^\mu a_b. \tag{5} \]

Eq. \(\text{(5)}\) contains the self-coupling of the volume modulus to its corresponding axion. From \(\text{(5)}\) we can directly calculate the decay rate \(\Phi \rightarrow a_b a_b\), obtaining

\[ \Gamma_{\Phi \rightarrow a_b a_b} = \frac{1}{48\pi M_P^2} \frac{m_\Phi^4}{3}. \tag{6} \]

As the bulk axion only interacts via Planck-suppressed couplings \[18\], the decays to \(a_b\) correspond to dark radiation.

2. Decays to visible sector fields

Let us consider decays to MSSM particles. Here we are using the expression ‘MSSM’ but the exact model is not so crucial and the considerations apply more generally.
Gauge bosons: Couplings of the volume modulus to
gauge bosons arise through the dependence of the gauge
kinetic function on $T_b$. However as the MSSM is realised
locally, the tree-level gauge couplings $f_a = S + h_a i T_{a,i}$
are independent of $T_b$. Such a dependence is certainly
induced radiatively, as the volume determines the string
scale and hence the high scale from which couplings start
running. At loop level there is then a term $\frac{b_a n_{SM}}{4\pi}$ ln $\nu$ in
the gauge kinetic function, which induces a coupling
\[ \lambda_a^{O_{SM}} \Phi \bar{F}_{\mu \nu} F^{\mu \nu} + \ldots \] (7)

However as a radiative correction this only gives
$\Gamma_{\Phi \rightarrow A_{SM}^* A_{SM}^*} \sim (\alpha_{SM})^2 \frac{m^3}{M_p^2}$, which is highly subdomi-
nant compared to eq. (3).

Matter scalars: The couplings to MSSM matter scalars
$C$ are set by the Kähler potential, which extends (3) to
$K = -3 \ln(T_b + \bar{T}_b) + \frac{CC}{(T_b + \bar{T}_b)}$. (8)
The $(T_b + \bar{T}_b)^{-1}$ dependence is fixed by the local nature
of the MSSM. After canonical normalisation the relevant
modulus-matter interaction is
\[ \frac{1}{2} \sqrt{3} \Phi (\bar{C} \square C + C \square \bar{C}) . \] (9)

Note that to eliminate $\Phi \partial_\mu \bar{C} \partial^\mu C$ couplings it is crucial to
use the $K_{C \bar{F}}$ terms. This interaction also vanishes for on-
shell massless particles, giving $\Gamma_{\Phi \rightarrow C C} \sim \frac{m_{SM}^2}{M_p^2} \ll \frac{m_p^3}{M_p^2}$.

Matter fermions, gauginos and Higgsinos: The couplings
to fermions can also be derived from the Kähler
potential and turn out to be of the form
\[ \lambda_{\Phi \bar{\Phi} \bar{\chi} \sigma^m D_m \chi} . \] (10)

However such couplings are chirality suppressed and give
$\Gamma_{\Phi \rightarrow f f} \sim \frac{m_{SM}^2}{M_p^2} \ll \frac{m_p^3}{M_p^2}$. Actually, as for gauge bosons we
expect this decay mode to be generated radiatively, again
giving a subdominant contribution $\Gamma \sim (\frac{m_p}{M_p})^2 \frac{m_p^3}{M_p^2}$.

Higgs bosons: We finally consider decays to Higgs
bosons. These, uniquely, have the possibility of a
Giudice-Masiero coupling in the Kähler potential
\[ K = -3 \ln(T_b + \bar{T}_b) + \frac{H_u H_u + H_d H_d}{(T_b + \bar{T}_b)} + \left( \frac{ZH_u H_d}{(T_b + \bar{T}_b)} + \text{h.c.} \right) . \]

Here $Z$ is an undetermined constant (to leading order in
an inverse volume expansion). If the Higgs sector has a
shift symmetry (as considered in [14]), then $Z = 1$ since
\[ K = \ldots + \frac{(H_u + \bar{H}_d)(\bar{H}_u + H_d)}{(T_b + \bar{T}_b)} + \ldots . \]

At string scale energies the Giudice-Masiero term can be
forgotten by anomalous $U(1)s$, in which case this term
is generated after breaking to the SM. After canonical
normalisation the resulting Lagrangian is
\[ \mathcal{L} = \frac{1}{2} \partial_\mu \Phi \partial^\mu \Phi + \partial_\mu H_u \partial^\mu \bar{H}_u + \partial_\mu H_d \partial^\mu \bar{H}_d \]
\[ + \frac{1}{\sqrt{6}} \left[ \Phi (\bar{H}_u \square H_u + \bar{H}_d \square H_d) + Z H_u H_d \square \Phi + \text{h.c.} \right] . \] (11)

The last term in the second line of eq. (11) allows for the
direct decay $\Phi \rightarrow H_u H_d$, with
\[ \Gamma_{\Phi \rightarrow H_u H_d} = \frac{2 Z^2 m_T^3}{48 \pi M_p^2} . \] (12)
The decay $\Phi \rightarrow H_u H_d$ is the only unsuppressed, and
determine dominant, MSSM decay mode. There will also
be subleading contributions to eq. (12) from SM radia-
tive corrections and running of the interaction strength
from $M_{\text{string}}$ to $m_\Phi$.

3. Decays to other axions

Local closed string axions: Another source of light
fields are closed string axions, the imaginary parts of
Kähler moduli. The existence of such axions is model
dependent as they can be dynamically lifted by non-
perturbative effects. Take the Kähler potential
\[ K = -3 \ln(T_b + \bar{T}_b) + \frac{Z_{ax}}{2} \left( \frac{X + \bar{X}}{T_b + \bar{T}_b} \right)^\gamma \] (13)

Here the axion is $\text{Im}(X)$ and for a local closed string
axion $\lambda = 3/2$. Moreover, $\gamma = 2$ if the saxion $\text{Re}(X)$ is
shrunk at a singularity while $\gamma = 3/2$ if the saxion is in the
geometric regime. Writing $X = \frac{1}{\sqrt{2}} (A + i B)$, after
normalisation the relevant couplings are
\[ \mathcal{L} = \frac{1}{2} \partial_\mu \Phi \partial^\mu \Phi + \frac{1}{2} \partial_\mu A \partial^\mu A + \frac{1}{2} \partial_\mu B \partial^\mu B \]
\[ + \sqrt{\frac{2}{3}} \frac{\lambda}{4M_p} (A^2 - B^2) \square \Phi . \] (14)

Note $Z_{ax}$ has disappeared from the Lagrangian and for $\gamma = 3/2$ one would have additional cross couplings
$\Phi \partial_\mu a \partial^\mu B$ which are suppressed by $\mathcal{V}^{-7/6}$, and so can
be neglected. This gives rise to
\[ \Gamma_{\Phi \rightarrow B B} = \left( \frac{\lambda}{(3/2)} \right)^2 \frac{9}{16} \frac{m_T^3}{48 \pi M_p^2} , \] (15)

which if present is an extra and competitive source of hidden
radiation.

Open string axions: Note that the presence of a QCD
axion does not guarantee the existence of local closed
string axions, as the QCD axion could arise from the
angular component $\theta$ of open string matter fields
$C = \rho e^{i \theta}$ which obtain a vev $\langle \rho \rangle \neq 0$. However, decays to
open string axions $\theta$ can be shown to be suppressed as they would be induced by terms of the form

$$-\sqrt{\frac{2}{3}} \left( \frac{\langle \rho \rangle}{M_P} \right)^2 \Phi \Phi \theta + \frac{4}{3} \left( \frac{\langle \rho \rangle}{M_P} \right)^2 \Phi \partial_\mu \theta \partial^\mu \theta.$$  

(16)

The first term in (16) gives rise to $\Phi \rightarrow \theta \Phi$ decays which are mass suppressed while the second term yields $\Phi \rightarrow \theta \Phi$ decays which could compete with the $\Phi \rightarrow a_b a_b$ decays for $\langle \rho \rangle \sim M_P$. However in general $\rho$ is fixed by D-terms and its vev is set by the Fayet-Iliopoulos term $\xi$.

$$\langle \rho \rangle / M_P \sim \xi \sim 1/V \ll 1.$$  

(17)

Axions as dark radiation have been considered in a different context in [17].

4. Decays to hidden sector fields

**Gauge bosons on the large cycle:** Another source of competitive decays are from light gauge bosons from branes wrapping the bulk cycle. Such branes might be present if an O7-plane wraps this cycle. In this case the bulk modulus has the direct coupling

$$n_g F_{\mu \nu} F^{\mu \nu} \rightarrow \frac{1}{4} F_{\mu \nu} F^{\mu \nu} + \frac{1}{4} \sqrt{\frac{2}{3} M_P F_{\mu \nu} F^{\mu \nu}},$$  

(18)

which induces a decay rate

$$\Gamma_{\Phi \rightarrow A_{\text{bulk}}^c A_{\text{bulk}}^a} = \frac{n_g}{2} \frac{m_\phi^3}{48\pi M_P^2}.$$  

(19)

where $n_g$ is the number of gauge generators. Such a bulk gauge group would represent a new hyperweak interaction, with $\alpha_{\text{bulk}} \lesssim 10^{-4}$. Note that decays to Higgses and superpartners would be kinematically forbidden since these sectors would not be sequestered, and so $M_{\text{soft}} \sim m_{3/2} \gg m_\phi$.

**Gauge bosons on small cycles:** In this case the decays to gauge bosons and gauginos would be kinematically forbidden in the presence of strong dynamics at a high scale (gaugino condensation). Hence the only worrisome decays are to light gauge bosons but they would be loop-suppressed similar to eq. [7].

**Bulk closed string $U(1)$s:** One model-independent source of massless gauge bosons are the bulk closed string $U(1)$s that arise from reduction of $C_4$ along 3-cycles. In general there are $h^{1,2}_+ U(1)$s, where $h^{1,2}_+$ is the number of 3-cycle invariant under the orientifold projection. However the gauge kinetic function of these is set by the complex structure moduli, giving

$$\mathcal{L} = f(U) F_{\mu \nu} F^{\mu \nu} + f(U) \lambda \alpha^\mu D_\mu \lambda.$$  

(20)

As mixing between the Kähler and complex structure moduli is volume-suppressed at large volumes, $\tau_6$ only couples to such $U(1)$s via volume-suppressed mixing terms, and so such decay modes are negligible.

**Sequestered hidden sectors:** Other competing decay rates are to Higgs bosons belonging to sequestered hidden sectors at different singularities since they would behave as eq. [12].

5. Summary of leading decay channels

Let us summarise all the leading decay channels for the volume mode:

1. Volume axion
2. Visible sector Higgses
3. Local closed string axions (if present)
4. Light gauge bosons on the large cycle (if present)
5. Higgses in sequestered hidden sectors (if present)

B. Predictions for dark radiation

Given that the presence of the volume axion and visible Higgses is model-independent, let us start focusing only on the decays $\Phi \rightarrow a_b a_b$ and $\Phi \rightarrow H_u H_d$. The relative fraction of hidden radiation is

$$f_{\text{hidden}} = \frac{1}{1 + 2Z^2}.$$  

(21)

Let us now relate $f_{\text{hidden}}$ to $\Delta N_{\text{eff}}$. Conservation of comoving entropy $s = g(T) a^3 T^3$ implies the SM temperature $T_{\text{SM}}$ and energy density $g(T) T^4$ behave as

$$T_{\text{SM}} \sim \frac{1}{g^{1/3} a}, \quad \rho_{\text{SM}} = \rho_{\text{SM}}^{\text{init}} g^{1/3} a^4.$$  

The ratio of hidden to SM radiation at neutrino decoupling temperature $T_{\text{dec}}$ is then

$$\frac{\rho_{\text{hidden}}}{\rho_{\text{SM}}} = \frac{\rho_{\text{init}}^{\text{hidden}}}{\rho_{\text{SM}}^{\text{init}}} \left( \frac{g(T_{\text{dec}})}{g(T_{\text{reheat}})} \right)^{1/3}.$$  

The number of extra effective neutrino species is

$$\Delta N_{\text{eff}} = 3 \frac{\rho_{\text{hidden}}}{\rho_{\text{neutrinos}}} = 43 \frac{\rho_{\text{hidden}}}{7 \rho_{\text{SM}}} = \frac{43}{7} \frac{f_{\text{hidden}}}{1 - f_{\text{hidden}}} \left( \frac{g(T_{\text{dec}})}{g(T_{\text{reheat}})} \right)^{1/3}.$$  

(22)

The reheating temperature can be easily estimated as

$$T_{\text{reheat}} \sim \sqrt{\Gamma_{\Phi} M_P} \sim O(1) \text{ GeV} \left( \frac{m_\phi}{3 \times 10^3 \text{ GeV}} \right)^{3/2}.$$  

For $T_{\text{reheat}} < 1 \text{ GeV}$, $g(T_{\text{reheat}}) = 247/4$, whereas for $T_{\text{reheat}} > 5 \text{ GeV}$, $g(T_{\text{reheat}}) = 345/4$, generating a small
uncertainty in the final prediction for $\Delta N_{\text{eff}}$ which for $g(T_{\text{dec}}) = 10.75$ is

$$3.12 \kappa \leq \Delta N_{\text{eff}} \leq 3.48 \kappa$$

where $\kappa = f_{\text{hidden}}/(1 - f_{\text{hidden}}) = 1/(2Z^2)$. For $Z = 1$, $\kappa = 1/2$, and so we obtain $1.56 \leq \Delta N_{\text{eff}} \leq 1.74$ which is comparable in magnitude to the observational hints.

If we allow a generic number of Higgs doublets $n_H$ and local closed string axions $n_a$, the expression for $\kappa$ generalises to $\kappa = (1 + 9n_a/16)/n_HZ^2$. In figure 1 we present a plot of $\Delta N_{\text{eff}}$ versus $n_H$ and $Z$ for $n_a = 0$ (and so the QCD axion has to be an open string mode). If $n_a \neq 0$, and in particular $n_a \sim O(100)$ in the case of an ‘axiverse’ \cite{19}, $\Delta N_{\text{eff}}$ quickly grows above the allowed experimental bounds.

Let us briefly consider other possible sources of hidden radiation. If light gauge bosons on $\tau_b$ exist, there are various possibilities:

- If the bulk gauge group is not in thermal equilibrium with the SM and is not Higgsed, its gauge bosons constitute dark radiation.
- If the bulk gauge group is not in thermal equilibrium with the SM and is Higgsed, its massive gauge bosons can constitute dark matter if they are stable.

If the bulk gauge group is in thermal equilibrium with the SM (which might be obtained via kinetic mixing of bulk $U(1)$s with the ordinary photon), then the bounds on $\Delta N_{\text{eff}}$ require it to be Higgsed at a scale $\gtrsim 1$ MeV, and this mode counts as SM radiation.

Notice that also hidden photons produced by photon $\leftrightarrow$ hidden photon oscillations in the thermal bath could form dark radiation \cite{26}. Finally, in the presence of hidden sequestered scenarios at different singularities, it is very likely that $\Delta N_{\text{eff}}$ would become too large and dark matter would be overproduced.

### III. CONCLUSIONS

The main point to emphasise is that dark radiation is generic and unavoidable in the sequestered LARGE Volume Scenario. This relies only on reheating being driven by decays of the lightest modulus, which always has an open decay mode to its axion partner. The magnitude of this dark radiation depends on assumptions about the visible sector and the number of additional closed string axions, but can easily be at a level consistent with observational hints for $\Delta N_{\text{eff}}$. The bounds on $\Delta N_{\text{eff}}$ can be used to strongly constrain the couplings and matter content of the models, showing that it is very hard to achieve an axiverse in sequestered models. Future observations, such as those expected from PLANCK, will further constrain this scenario.

**Note added:** This paper is submitted simultaneously to the related work \cite{21}.

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[22] We shall not consider here field redefinitions such as in [13]. Independent of the subtle questions on the conditions required for their existence, if present and with the effects estimated in [13] they bring soft terms to the scale of the gravitino mass, restoring the full-blown CMP.