PHYSICAL REVIEW D **91,** 043503 (2015)

Inflationary constraints on modulus dominated cosmology

Koushik Dutta[1] and Anshuman Maharana[2]

[1]*Saha Institute of Nuclear Physics, 1/AF Salt Lake, Kolkata 700064, West Bengal, India*

[2]*Harish Chandra Research Intitute, Chattnag Road, Jhunsi, Allahabad 211019, Uttar Pradesh, India*

(Received 16 October 2014; published 4 February 2015)

We study moduli fields that arise in string/supergravity models in the context of the inflationary scenario. The early time cosmological dynamics involves generation of density perturbations by quantum fluctuations of the inflaton; the late time dynamics involves a modulus field dominating the energy density of the Universe and then its decay. We derive a relation which relates the modulus mass, inflationary observables and broad features of the reheating epoch. When viewed along with generic expectations regarding reheating and the initial field displacement of the modulus after inflation, this gives a bound on the modulus mass. For a large class of models, the bound is much stronger than the cosmological moduli problem bound.

DOI: 10.1103/PhysRevD.91.043503 PACS numbers: 98.80.Cq, 11.25.Wx

I. INTRODUCTION

The hot big bang model together with the inflationary paradigm provides a highly attractive framework for cosmology. Typically, it is assumed that after inflation the visible sector degrees of freedom reheat and have evolved adiabatically since then. In spite of the impressive successes of the models based on this assumption, it is important to keep in mind that from the point of view of supergravity and string models late time entropy production is very well motivated.

At tree level, string vacua usually contain massless scalar fields (the moduli) which interact only via Planck suppressed interactions. Moduli fields acquire masses from subleading effects in the effective action; their masses are expected to be well below the string scale. If the postinflationary mass of a modulus is below the Hubble scale during inflation (we will refer to such moduli as light moduli), it can dominate the energy density of the Universe at late times and then decay producing significant entropy (as we will briefly review in Sec. II). The modulus has to decay prior to nucleosynthesis in order to account for the success of big bang theory in predicting the abundances of light elements. This gives a bound on the mass of the modulus [1]—the cosmological moduli problem (CMP) bound; $m_\varphi \gtrsim 30$ TeV. At present, this is the strongest available bound on moduli masses and serves as an extremely useful guide for string model building. In this paper, we provide a new bound on moduli masses in string/supergravity models based on precision cosmic microwave background (CMB) data (as in the CMP bound our focus shall be on moduli which decay via Planck suppressed interactions).

If quantum fluctuations of the inflaton are responsible for the observed density perturbations, then the energy density at the time of horizon exit of the pivot mode can be obtained from CMB data via determination of the primordial scalar amplitude (A_s) and the tensor to scalar ratio (r). Any history we ascribe to the Universe between the time of horizon exit of the pivot mode and today should be consistent with the fact that the energy density at horizon exit has evolved to the observed energy density today. We find that in the context of cosmologies in which a modulus dominates the energy density at late times, this consistency condition implies that an increase in the number of e-foldings of modulus domination must be compensated by a decrease in the number of e-foldings during inflation. Thus there is a tension between having a very light modulus (which lives very long and leads to a large number of e-foldings of modulus domination) and having the number of e-foldings during inflation as determined from precision CMB data for a given inflationary model. Our bound is a manifestation of this tension.

For a large class of inflationary models the bound can be much stronger than the CMP bound. It is likely to have broad implications for moduli stabilization, supersymmetry breaking and inflationary model building in string theory.

II. LATE TIME MODULUS DOMINATED COSMOLOGY

At the end of inflation, the expectation value of a light modulus differs from the value of the field at the minimum of its postinflationary potential. This "initial displacement" $\varphi_{\rm in}$ can take place due to quantum/thermal effects [2] and/or explicit dependence of the modulus potential on the inflaton expectation value [3], and is expected to be of the order of $M_{\rm pl}$. At this stage, the modulus vacuum expectation value remains frozen at $\varphi_{\rm in}$ due to the large friction induced by the large value of the Hubble constant. With reheating, the energy associated with the inflaton gets converted to radiation; the Hubble constant decreases as the Universe expands. When the Hubble constant drops below the mass of the modulus, the friction term becomes irrelevant. The modulus begins to oscillate about the minimum of its potential. Subsequently, the energy density associated with the field begins to redshift like matter, at a rate much slower than that of radiation—the energy associated with the modulus can quickly dominate the energy density of the Universe. Eventually the modulus decays reheating the Universe. The last modulus to decay provides a set of "initial conditions" for cosmological evolution.

Properties such as dark matter density and baryon asymmetry are determined by the branching ratio of the various decay products. The phenomenology of such scenarios is an active area of study; see [4] and references therein. Late time modulus dominated cosmology has emerged as the preferred cosmological scenario in various string constructions; M-theory models [5,6] and large volume compactifications in type II B [7,8]. It has also been suggested to be a generic prediction of string theory [9].

III. INFLATIONARY CONSTRAINTS ON THE MASS OF LIGHT MODULI

Our working assumption will be that quantum fluctuations of the inflaton are solely responsible for the observed density perturbations. The history of the Universe will be taken to be as described in Sec. II. We will be explicitly including only one modulus in our analysis. Comments on the multiple modulus case (which can be considered generic in string compactifications) will be made in the text.

There are two phases of reheating (production of relativistic particles)—the one after inflation and one after the decay of the modulus. Each of these epochs can be characterized by the number of e-foldings in the epoch and the effective equation of state. We will denote the number of e-foldings during the reheating epoch (from decay of inflaton) after inflation by $N_{\rm re}$ and the effective equation of state during this epoch by $w_{\rm re}$. For the epoch of reheating after the decay of the modulus, the number of e-foldings will be denoted by $N_{\rm re2}$, and the effective equation of state by $w_{\rm re2}$. To streamline our discussion, we present the derivation for an instantaneous second reheating phase ($N_{\rm re2} = 0$) first. Later, we will discuss the effects of the second reheating phase and will see that the associated effects do not alter our conclusions at all; see Eq. (17). For the reheating phase after inflation we will be sensitive to only the broad characteristics $N_{\rm re}$ and $w_{\rm re}$; given this we hope to capture not only the transfer of energy from the inflaton to radiation but also of the decay of very heavy moduli (which decay much earlier) by this reheating phase.

We begin the derivation of our relation by writing the condition which determines the exit of a mode of comoving wave number k from the horizon $k = a_{\rm k} H_{\rm k}$ as

$$k = \frac{a_{\rm k}}{a_{\rm end}} \cdot \frac{a_{\rm end}}{a_{\rm re}} \cdot \frac{a_{\rm re}}{a_{\rm eq}} \cdot \frac{a_{\rm eq}}{a_{\rm decay}} . a_{\rm decay} H_k, \tag{1}$$

where the subscripts end, re, eq and decay indicate the end of inflation, end of reheating after inflation, equality of energy density between matter (in this case modulus energy density) and radiation, and decay of the modulus. Taking the logarithm of (1), one obtains

$$N_{\rm matdom} = -N_{\rm k} - N_{\rm re} - N_{\rm rad} - \ln k + \ln(a_{\rm decay}) + \ln H_k, \tag{2}$$

where $N_{\rm matdom}$ is the number of e-foldings in the matter (modulus) dominated era, $N_{\rm k}$ the number of e-foldings between the horizon exit of the pivot mode and end of inflation, $N_{\rm re}$ the number of e-foldings during the period of reheating after inflation and $N_{\rm rad}$ the number of e-foldings in the radiation dominated era.

Next, we obtain another expression for $N_{\rm matdom}$ based on the evolution of energy density. We begin by writing[1]

$$-N_{\rm matdom} = \frac{1}{3} \ln(\rho^{\rm matter}_{\rm decay} / \rho^{\rm matter}_{\rm eq}). \tag{3}$$

The energy density at the time of decay of the modulus can be expressed in terms of the reheat temperature $T_{\rm re2}$ (after the decay of the modulus) and the effective number of light species $g_{\rm re2}$ at the time of reheating

$$\rho^{\rm matter}_{\rm decay} \approx \rho_{\rm decay} = (\pi^2/30) g_{\rm re2} T^4_{\rm re2}, \tag{4}$$

and the reheat temperature can be related to the observed temperature of the CMB today (assuming no entropy production at a later stage) by

$$T_{\rm re2} = (43/11 g_{\rm s,re2})^{1/3} (a_0/a_{\rm decay}) T_0, \tag{5}$$

where $g_{\rm s,re2}$ is the effective number of light species for entropy. Combining Eqs. (4) and (5) and expressing $\ln \rho^{\rm matter}_{\rm eq}$ as

$$\ln(\rho^{\rm matter}_{\rm eq}) = \ln(\rho^{\rm radiation}_{\rm eq} / \rho_{\rm re}) + \ln(\rho_{\rm re}/\rho_{\rm end}) + \ln(\rho_{\rm end}), \tag{6}$$

Eq. (3) yields (we take $\rho^{\rm radiation}_{\rm re} \simeq \rho_{\rm re}$)

$$-\frac{3}{4} N_{\rm matdom} = \frac{1}{4} \ln(\pi^2 g_{\rm re2}/30) + \frac{1}{3} \ln(43/11 g_{\rm s,re2}) + \ln(a_0 T_0 / a_{\rm decay}) + N_{\rm rad} - \frac{1}{4} \ln(\rho_{\rm end}) + \frac{3}{4}(1 + w_{\rm re}) N_{\rm re}. \tag{7}$$

Adding (2) and (7) the dependence on both $N_{\rm rad}$ and $a_{\rm decay}$ drops out; we obtain

[1]An energy density ρ with a superscript will denote the energy density in a given form; $\rho^{\rm matter}_{\rm decay}$ is the energy density in the form of matter at the time of modulus decay.

$$\frac{1}{4}N_{\rm matdom} = \frac{1}{4}\ln(\pi^2 g_{\rm re2}/30) + \frac{1}{3}\ln(43/11 g_{\rm s,re2}) - N_k - \ln(k/a_0T_0) + \ln H_k - \frac{1}{4}\ln(\rho_{\rm end}) - \frac{1}{4}(1-3w_{\rm re})N_{\rm re}. \quad (8)$$

We note that the above equation can be thought of as a generalization of the formula which gives the total number of e-foldings in inflationary models (see e.g. [10]).

Now, $N_{\rm matdom}$ can be expressed in terms of the modulus mass and lifetime by using the explicit form of the scale factor as a function of time. Recall that if the equation of state is w the evolution of the scale factor between times t_1 and t_2 is given by

$$(a(t_2)/a(t_1))^{\frac{3}{2}(1+w)} = 1 + \frac{3}{2}(1+w)H(t_1)(t_2-t_1). \quad (9)$$

By demanding that the time elapsed between the end of inflation and the decay of the modulus is the lifetime of the modulus (and assuming $N_{\rm matdom} \gg 1; N_{\rm re}, N_{\rm rad} > 1$), we obtain

$$N_{\rm matdom} \approx \frac{2}{3}\log\left(\frac{3}{2}H_{\rm eq}\tau_{\rm mod}\right), \quad (10)$$

where $\tau_{\rm mod}$ is the lifetime of the modulus. It can easily be checked that the above (approximate) expression is also correct in the regime $N_{\rm matdom} \gg 1$ and $N_{\rm re}, N_{\rm rad} \ll 1$. $H_{\rm eq}$ can be obtained by computing the energy density at equality. Recall that the modulus begins to oscillate about its minimum when the Hubble constant becomes of the order of its mass; this implies (making the usual assumption of radiation domination at this time) $\rho^{\rm radiation}(t_{\rm osc}) = 3m_\varphi^2 M_{\rm pl}^2$. Also, the matter density at this time $\rho^{\rm matter}(t_{\rm osc}) = \frac{1}{2}m_\varphi^2\varphi_{\rm in}^2$. From this one obtains

$$\rho_{\rm eq} = m_\varphi^2\varphi_{\rm in}^2(\varphi_{\rm in}^2/6M_{\rm pl}^2)^3. \quad (11)$$

Using the value of $H_{\rm eq}$ derived from (11) in (10) and parametrizing the initial displacement as $\varphi_{\rm in} = YM_{\rm pl}$, we obtain

$$N_{\rm matdom} = -\frac{2}{3}\ln 3 - \frac{5}{3}\ln 2 + \frac{2}{3}\ln m_\varphi\tau_{\rm mod} + \frac{8}{3}\ln Y. \quad (12)$$

Equating the two expressions for $N_{\rm matdom}$ given by (3) and (8), and making use of the slow-roll expression for the Hubble constant, $H_k^2 = \rho_k/3M_{\rm pl}^2 = \frac{1}{2}\pi^2 M_{\rm pl}^2 A_s r$, one finds

$$\frac{1}{6}\ln m_\varphi\tau_{\rm mod} + \frac{1}{4}(1-3w_{\rm re})N_{\rm re} + \frac{2}{3}\ln Y = \frac{1}{4}\ln(\pi^2 g_{\rm re2}/30) + \frac{1}{3}\ln(43/11g_{\rm s,re2}) + \frac{1}{12}\ln(4/3) - N_k - \ln(k/a_0T_0) - \frac{1}{4}\ln(\rho_{\rm end}/\rho_{\rm k}) + \frac{1}{4}\ln(\pi^2 rA_s). \quad (13)$$

The characteristic lifetime of a modulus field in string/supergravity models can be obtained from some generic considerations. Moduli fields are uncharged under the standard model (or any hidden sector) gauge groups. Their decay occurs via nonrenormalizable interactions; for a large class of moduli fields the decays occur by Planck suppressed interactions primarily to radiation, in which case the lifetime is (see e.g [6,9,11])

$$\tau_{\rm mod} \approx \frac{16\pi M_{\rm pl}^2}{m_\varphi^3}. \quad (14)$$

Although, there can be interesting exceptions to this. In the case of brane world constructions the ultraviolet scale Λ which suppresses the interaction strength can be different from the Planck scale for some of the moduli (it can be lower or higher; see e.g. [12]). Also, the above lifetime is for the case when the decay is primarily to massless radiation; its functional form changes can be different if the decay to massive products dominates. For now, we proceed by taking the decay width to be as given by (14) since it is used in arriving at the CMP bound [1] described in the introduction (we will discuss the effects of having the strength of moduli interactions suppressed by a scale different from $M_{\rm pl}$ and the case of decay primarily to massive particles later).

The dependence on the number of degrees of freedom appears as $\ln(g_{\rm re2}^{1/4}/g_{\rm s,re2}^{1/3})$, and hence is quite mild. We use $g_{\rm re2} \approx g_{\rm s,re2} \approx 100$. We use Planck data [10] for quantities that have already been observed with accuracy; the primordial scalar amplitude $A_s = 2.20 \times 10^{-9}$ at the pivot scale $k = 0.05\ {\rm Mpc}^{-1}$ and $T_0 = 2.725\ K$. Plugging in all this, we find

$$\frac{1}{6}\ln\left(\frac{16\pi M_{\rm Pl}^2}{m_\varphi^2}\right) + \frac{1}{4}(1-3w_{\rm re})N_{\rm re} + \frac{2}{3}\ln Y = 55.43 - N_k + \frac{1}{4}\ln r + \frac{1}{4}\ln\left(\frac{\rho_{\rm k}}{\rho_{\rm end}}\right). \quad (15)$$

The above equation is our main result.

There are strong reasons to believe (supported by both analytic and numerical work) that the equation of state during reheating satisfies $w_{\rm re} < 1/3$; see e.g. [10,13] for discussions. Guided by this we will take (for now) $w_{\rm re} < 1/3$; this makes the second term on the right-hand side of (15) positive definite. Equation (15) then gives a bound on the mass of the modulus

$$m_\varphi \gtrsim \sqrt{16\pi}M_{\rm pl}Y^2 e^{-3(55.43 - N_k + \frac{1}{4}\ln(\frac{\rho_{\rm k}}{\rho_{\rm end}}) + \frac{1}{4}\ln r)}. \quad (16)$$

Recall that Y is the initial field displacement of the light modulus in Planck units. As discussed in Sec. II, the generic expectation for the initial displacement is of the order of M_{pl} [2,3]. Thus Y cannot affect the value of the right-hand side of (16) significantly.[2] The effect of including a reheating phase after the decay of the modulus field is a new term in (15) which depends on the number of e-foldings (N_{re2}) and equation of state (w_{re2}) of this reheating phase and has exactly the same form as the second term on the left-hand side of (15); it gets modified to

$$\frac{1}{6}\ln\left(\frac{16\pi M_{\mathrm{Pl}}^2}{m_\varphi^2}\right)+\frac{1}{4}(1-3w_{\mathrm{re}})N_{\mathrm{re}}+\frac{2}{3}\ln Y +\frac{1}{4}(1-3w_{\mathrm{re2}})N_{\mathrm{re2}}=55.43-N_k+\frac{1}{4}\ln r +\frac{1}{4}\ln\left(\frac{\rho_{\mathrm{k}}}{\rho_{\mathrm{end}}}\right). \tag{17}$$

Following the above arguments, the bound is unchanged.

Note that the larger the number of e-foldings during inflation, the stronger the bound. The second parameter in the exponent, $\frac{1}{4}\ln(\rho_{\mathrm{k}}/\rho_{\mathrm{end}})$, is positive definite; on the other hand the third parameter $\frac{1}{4}\ln r$ is negative definite.

We briefly comment on the multiple modulus case. As mentioned earlier, given the general parametrization of the reheating phase, the dynamics of heavier moduli which decay very early on should be captured in the reheating phase. The relevant dynamics involves epochs of matter domination and radiation domination, and should satisfy the bound $w_{\mathrm{re}}<1/3$. If there are N moduli at the same mass scale (with a diagonal Kahler metric or if we make the assumption that the Kahler metric is generic as in [14]) then the energy density at equality (11) scales as N^4 (for fixed φ_{in}); the bound becomes stronger by a factor of N.

We note that even for an exotic reheating phase with $w_{\mathrm{re}}>1/3$, (13) predicts values m_φ to be quite large for $N_k\approx 50$, as long as the number of e-foldings during reheating is not comparable to the number of e-foldings of modulus domination (as is expected for a light modulus). We will discuss this later; focus on (16) for now.

To get a better understanding of the bound, we examine it in detail in the context of small field and large field models.

Small field models: For these models the potential is plateaulike and the change in energy density between horizon exit and the end of inflation is small. It is reasonable to drop the term involving the logarithm of the two energy densities in the exponent of (16). Taking $r=0.01$, we get

$$m_\varphi\gtrsim\sqrt{16\pi}M_{\mathrm{pl}}Y^2e^{-3(54.28-N_{\mathrm{k}})}. \tag{18}$$

To get a feel for the numbers, $N_{\mathrm{k}}=50$ and taking $Y=1/10$ (in what follows, we will always take $\mathrm{Y}=1/10$ while quoting numbers) we get

$$m_\varphi\gtrsim 4.5\times10^8\ \mathrm{TeV}, \tag{19}$$

which is well above the CMP bound ($m_\varphi\gtrsim 30$ TeV).

From the point of view of scalar field potentials that drive inflation,

$$N_{\mathrm{k}}\approx\frac{\beta}{1-n_s}, \tag{20}$$

with β being a model dependent constant. A survey of the values of β associated with various models is given in [15]; for most models $\beta\approx2$. Planck gives the value of n_s at 1σ as $n_s=0.9603\pm0.0073$. For $\beta\approx2$, at the central value of n_s, $N_k\approx50$. The bound is $m_\varphi\gtrsim4.5\times10^8$ TeV. The functional form of the bound makes it highly sensitive to the value of n_s; the Planck upper limit (at 1σ) gives $m_\varphi>M_{\mathrm{pl}}$ (ruling out late time modulus cosmology for small field models with $\beta=2$, $r=0.01$) while the lower limit gives $m_\varphi\gtrsim0.1$ TeV. This sensitivity implies that future experiments [16] will play an important role (the uncertainty in n_s will be reduced by one order of magnitude and this will bring uncertainty in the mass to just two orders of magnitude). Theoretically, the exponential sensitivity of the bound to N_k implies that $1/N_k$ corrections to (20) can be relevant in some models.

Large Field Models: As prototypes of the large fields models, we consider models where the inflation potential is a monomial $V_\chi=\frac{1}{2}m^{4-\alpha}\chi^\alpha$ (keeping the "$m^2\chi^2$" model [17] and axion monodromy [18] models in mind). Given the uncertainty in the measurements of r, we will take α as a model building parameter and take observational input only from n_s. The bound (16) simplifies to

$$m_\varphi\gtrsim\sqrt{16\pi}M_{\mathrm{pl}}Y^2e^{-3\left(55.85-\frac{(2+\alpha)}{2(1-n_s)}+\frac{\alpha}{8}\ln2+\frac{1}{8}(\alpha-2)\ln\left(\frac{2+\alpha}{\alpha(1-n_s)}\right)\right)}. \tag{21}$$

For the $\frac{1}{2}m^2\chi^2$ model and n_s at the central value of Planck the bound is $m_\varphi\gtrsim10^7$ TeV, one order of magnitude below (19). On the other hand for $\alpha=1$ and n_s at the same value ($N_k\approx40$) the bound becomes $m_\varphi\gtrsim10^{-10}$ TeV, which is completely irrelevant.

In summary, the larger the number of e-foldings during inflation, the stronger the bound. For $\beta\approx2$ (which corresponds to $N_k\approx50$ for the central value of Planck data) the bound is strong.

For the nongeneric case of exotic reheating with $w_{\mathrm{re}}>1/3$, it is useful to parametrize the duration of reheating by a parameter λ; $N_{\mathrm{re}}(1-3w_{\mathrm{re}})=-\lambda N_{\mathrm{matdom}}$. Even for $\lambda=1/3$ (which corresponds to a rather long phase of exotic reheating) direct use of (13) gives (for $N_k\approx50$) $m_\varphi\approx10^6$ TeV, which is, again, well above the CMP bound.

As discussed earlier, we have taken moduli interactions to be Planck suppressed in obtaining (16). In string

[2]We note in passing that in deriving the CMP bound one also makes use of the fact that Y is not expected to be significantly less than one.

constructions of brane world models there can be moduli whose interactions are not Planck suppressed but by a scale[3] Λ. If such a modulus (χ) is the last to decay the CMP bound [1] gets modified to the reheat temperature after the decay of a modulus is given by $T_{\rm reheat} \sim \sqrt{\Gamma M_{\rm pl}}$, where Γ is the width of the modulus. The characteristic width of a modulus χ whose interactions in the four-dimensional effective action are suppressed by a scale Λ is given by

$$\Gamma_{\Lambda} \approx \frac{16\pi m_{\chi}^{3}}{\Lambda^{2}}. \quad (22)$$

Combining the above with the requirement of a sufficiently high reheat temperature for nucleosynthesis, one arrives at a generalization of the CMP bound [1] discussed in the introduction,

$$m_{\chi} \gtrsim \eta^{2/3}.30 \text{ TeV}, \quad (23)$$

where $\eta = \Lambda/M_{\rm pl}$.

Following the same steps as in the earlier part of this section [while using the lifetime of the modulus as given by (22)], one can obtain the modification of our bound,

$$m_{\varphi} \gtrsim \sqrt{16\pi} M_{\rm pl} \eta Y^{2} e^{-3(55.43 - N_{k} + \frac{1}{4}\ln(\frac{\rho_{\rm k}}{\rho_{\rm end}}) + \frac{1}{4}\ln r)}. \quad (24)$$

We note that both (23) and (24) scale as a positive power of η. Carrying out an analysis as above, one easily sees that our bound is stronger in a large range of the phenomenologically interesting parameter space.

Finally, we briefly discuss the cases in which the modulus primarily decays to massive particles. Such decay products can be superpartners of standard model particles or additional Higgs bosons. For models in which the primary mode of decay is to massive particles and the lifetime scales as m_{φ}^{p} with $p \leq -1$ our analysis will provide a lower bound on moduli masses. The bound might involve the mass of the decay products [expression for the bound will in general be different from that given in Eqs. (18) and (21)]. In a large number of situations, the lifetime has the same form as (14) or has the form (see e.g [8,19])

$$\tilde{\tau} \approx \frac{16\pi M_{\rm pl}^{2}}{m_{\varphi} \tilde{m}^{2}} \quad (25)$$

(i.e. $p = -1$) where $\tilde{m}$ is the mass of the decay products. Again, following the same steps as in the earlier part of the section one obtains a bound on the mass of the decay products [in the case that the lifetime takes the form (25)] or a bound on the modulus mass [in the case that the lifetime takes the same form as (14)]. But a bound on the mass of the decay products translates to a bound on the mass of the modulus, as the mass of the modulus has to be heavier than the mass of the decay products. Thus, the bound (16) applies equally well for these situations (with $p = -1$). On the other hand, for $p > -1$ our analysis will provide an upper bound for moduli masses (in terms of the mass of the decay products). This can be very interesting, although such models are not generic. We leave the detailed study of specific models for future work.

IV. CONCLUSIONS

We have considered cosmologies in which density perturbations are generated by quantum fluctuations of the inflaton at early times; the late time dynamics involves a modulus which first dominates the energy density of the Universe and then decays to reheat the visible sector. Making use of generic expectations regarding reheating ($w_{\rm re} \leq 1/3$) and initial displacement of the modulus at the end of inflation ($\varphi_{\rm in} \sim M_{\rm pl}$), taking the decay properties of the modulus to be same as in the CMP bound [1] and using CMB data as input, we arrived at a bound on the minimum mass of the modulus; see Eq. (16). For values of the number of e-foldings during $N_{\rm k} \gtrsim 50$, the bound is much stronger than the CMP bound. In the case of instantaneous reheating, the bound becomes an equality and gives a prediction for the mass of the modulus.

The bound should have broad implications for string and supergravity models where it is typical to have scalars interacting with Planck suppressed interactions. It can shed light on the scale of supersymmetry breaking in the context of gravity mediated breaking, where the scale of soft masses can be tied to the moduli masses. The bound is exponentially sensitive to the number of e-foldings during inflation and hence provides a new motivation for precision measurements of the spectral tilt.

ACKNOWLEDGMENTS

K. D. thanks Harish Chandra Research Institute, Max Planck Institute Munich and Abdus Salam International Centre for Theoretical Physics, Trieste for hospitality. K. D. is partially supported by a Ramanujan fellowship and a Department of Science and Technology (India)-Max Planck Society fellowship. A. M. is partially supported by a Ramanujan fellowship.

[3]The scale Λ can be lower than the Planck scale for the modulus which parametrizes the size of the cycle that the branes wrap. In this case Λ is the string scale (see e.g. [12]). We note that there is a large difference between the string and Planck scale only if the volume of the compactification is large.

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