A model of cosmology and particle physics at
an intermediate scale

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Abstract

We propose a model of cosmology and particle physics in which all relevant scales arise in
a natural way from an intermediate string scale. We are led to assign the string scale to
the intermediate scale $M_* \sim 10^{13}$GeV by four independent pieces of physics: electroweak
symmetry breaking; the $\mu$ parameter; the axion scale; and the neutrino mass scale. The
model involves hybrid inflation with the waterfall field $N$ being responsible for generating
the $\mu$ term, the right-handed neutrino mass scale, and the Peccei-Quinn symmetry breaking
scale. The large scale structure of the Universe is generated by the lightest right-handed
sneutrino playing the rôle of a coupled curvaton. We show that the correct curvature
perturbations may be successfully generated providing the lightest right-handed neutrino
is weakly coupled in the see-saw mechanism, consistent with sequential dominance.

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1 Introduction

WMAP [1] has provided an unprecedented glimpse into the early universe at the time of radiation decoupling, which strengthens the case for a period of cosmological inflation [2]. With inflation becoming increasingly established, the need for a synthesis between cosmology and particle physics becomes ever more pressing. Such a synthesis should provide a successful cosmological model of inflation, and cosmological perturbations which can provide the seed of large scale structure. It should give successful baryogenesis, for example via leptogenesis, and should generate the required cold dark matter abundance. Ultimately it should also explain the dark energy, but this is more ambitious since it first requires a solution to the cosmological constant problem, which from our present perspective seems very far away.

To achieve such a synthesis the theory should also give a successful description of particle physics phenomena such as right-handed neutrino masses $M_{RR}$ and a solution to the strong CP problem such as provided for example by the Peccei-Quinn mechanism involving an intermediate scale axion at a scale $f_a$. The theory should also be supersymmetric, to stabilise the hierarchy and provide flat directions for inflation, in which case it should also provide an origin of the Higgs $\mu$ mass parameter. Ideally the theory should provide a complete explanation of electroweak symmetry breaking, not only for example in terms of radiative breaking, but also an explanation for the origin of the weak scale $M_W$. In fact from the point of view of string theory there is only one fundamental parameter namely the string scale $M_s$, from which the Planck scale $M_P$ should be derived in terms of the compactification scales. From this single scale $M_s$ one must be able to derive all the relevant scales in physics above, such as the axion scale $f_a$, the scales of right-handed neutrino masses $M_{RR}$, the $\mu$ parameter, and the weak scale $M_W$, which in the framework of supersymmetric theories is related to the soft supersymmetry breaking masses $m_{soft}$. The successful synthesis would therefore also be expected to provide an explanation of all these scales in terms of a single mass scale $M_s$.

Recently we proposed a very promising model of inflation closely related to the supersymmetric standard model [3]. The explicit model was an extra-dimensional model [3] with an intermediate string scale $M_s \sim 10^{13}$ GeV. This model in turn was based on an earlier model without extra dimensions [4], and the purpose of embedding the model in extra dimensions is to provide a natural explanation for the small Yukawa couplings and various mass scales appearing in the model. The lightest natural mass scale in the model turns out to be an MeV, and this requires that the cosmological perturbations in the model to be generated from a new mechanism.
which depend on isocurvature perturbations in the slowly rolling Higgs field to be transferred to curvature perturbations during reheating [5]. This mechanism [5] in which isocurvature perturbations of a flat direction in hybrid inflation become converted to curvature perturbations during reheating, may be called a coupled curvaton scenario to distinguish it from the weakly coupled late-decaying curvaton scenario [6]. We recently showed that in the coupled curvaton scenario, preheating plays a crucial rôle in the conversion of the isocurvature perturbations to the curvature perturbations [7].

The purpose of the present paper is two-fold. Firstly by providing a careful analysis of mass scales in the model we are led to the remarkable conclusion that all mass scales follow from a single input physical mass scale, namely the string scale $M_*$ which is uniquely fixed by physics to be of order $10^{13}$ GeV. This intermediate scale manifestly determines the (heaviest) right-handed neutrino mass scale $M_{RR}$ as well as the Peccei-Quinn symmetry breaking or axion scale $f_a$. In addition we show how the assumption of a small cosmological constant fixes the value of the supersymmetry breaking scale $F_S$ in terms of the string scale $M_*$ and the Planck scale $M_P$. This requirement reduces the number of free parameters in the model to just one, the string scale $M_*$, together with the Planck scale $M_P$ which may in principle be determined by the string dynamics by a ratio of $M_*/M_c \sim 60$ where $M_c$ is the compactification scale. From $M_*$ and the (in principle) string determined $M_P$ we show how the model then determines all the physical scales of interest, as well as the dimensionless couplings.

The second main purpose of the paper is to show that the large scale structure in the Universe may be generated by the lightest right-handed sneutrino playing the rôle of a coupled curvaton\(^1\). Having a sneutrino rolling along a flat direction during inflation and playing a special rôle in determining the large scale structure of the universe only strengthens the synthesis of particle physics and cosmology in this approach. An oscillating sneutrino may also allow efficient leptogenesis to take place during the reheating process.

The layout of the remainder of the paper is as follows. In section 2 we review the model, including the superpotential, and then discuss qualitatively how all the parameters which enter these potentials can be obtained from a single mass scale $M_*$. In section 3 we discuss the potential and the symmetry breaking aspects of the model. In section 4 we fix the mass scale of the theory. We also discuss the subtle interplay between the high energy gauge group symmetry breaking, supersymmetry and electroweak symmetry breaking, and show how the weak scale may be derived in terms of the string scale. We also show how the axion scale, the right-

\(^1\)In [8] the right-handed sneutrino plays the rôle of the weakly coupled late-decaying curvaton.
handed neutrino mass scale and the origin of the $\mu$ parameter are all associated with the same string scale. Section 5 discusses the physics during the inflationary epoch including the special rôle played by the sneutrino as a coupled curaton. We also discuss the physics of preheating and reheating which is expected to give rise to curvature perturbations. Section 7 contains our conclusions. There are two Appendices detailing the calculation of Yukawa couplings (Appendix A) and soft masses (Appendix B) in the presence of extra dimensions.

2 The model

Let us consider two four dimensional boundaries $^2$ spatially separated along $d$ extra dimensions with a common radius $R = 1/M_*$. These extra dimensions are compactified on some orbifold that leads at least two fixed points at $\{0, \pi R\}$ where the two D3-branes (Brane I and Brane II) are located. All the families of quarks and leptons are localized in one of the fixed point ($y = 0$), Brane I, while SUSY is broken by the $F$-term of a gauge singlet field $S$ localized in the parallel brane ($y = \pi R$), Brane II. The gauge group is $G_A \times G_B$ where $G_A$ and $G_B$ are localized in the bulk and the SUSY breaking brane respectively. The rest of the matter are localized in the bulk, namely the inflaton field, $\phi$, the waterfall field, $N$, the MSSM Higgs fields $h_u$ and $h_d$, and the massive Higgs fields $H_1$, $H_2$ which mediate the breakdown of the gauge group $G_A \times G_B$ to the Standard Model gauge group $G_{SM} = SU(3) \times SU(2) \times U(1)$. This is depicted in Fig. (1).

One of the underlying assumption of this model is that the radii of the extra dimensions are stabilized before inflation takes place, for example, by one of the mechanism proposed in the literature (see for example [9]). Done this, firstly we are going to discuss the size of the effective four dimensional parameters (Yukawa couplings and soft masses) in a general model with two parallel branes and then discuss particular issues of our model.

The four dimensional superpotential of the model in Fig. (1) is given by

$$W_4 = -\kappa \phi N^2 + \lambda_N N h_u h_d + \lambda_H \phi H_1 H_2 + \lambda_{\nu,1} \frac{N^2 (\nu^c_{R,1})^2}{M_*} + W_{MSSM},$$

(1)

where $W_{MSSM}$ defines the Yukawa couplings for the MSSM

$$W_{MSSM} = y_u^{ij} Q_i h_u U_j + y_d^{ij} Q_i h_d D_j + y_e^{ij} L_i h_d e^c_{Rj} + y_{\nu}^{ij} L_i h_u \nu^c_{Rj}.$$  

(2)

$^2$From now on let us use the word D3-brane instead of four dimensional boundary. In spite of the abuse of language in this choice, schematically it may represent the string connection of our model and also it provides a simplification of the English in this paper.
Figure 1: The model showing the parallel 3-branes spatially separated along $d$ extra dimensions with coordinates $y = (y_1, \ldots, y_d)$ and a common radius $R$. The index $i$ in the matter fields represents the family index, $i = 1, 2, 3$.

Following the relationship between higher dimensional couplings ($\hat{\lambda}_{i,j}$) and four dimensional one ($\lambda_{i,j}$) given in Appendix A in Eq. (34), we get

$$
\kappa = \left( \frac{M_*}{M_P} \right)^3 \hat{\kappa}, \quad 
\lambda_N = \left( \frac{M_*}{M_P} \right)^3 \hat{\lambda}_N, \quad 
\lambda_H = \left( \frac{M_*}{M_P} \right)^3 \hat{\lambda}_H
$$

$$
\lambda_{\nu} \equiv \lambda_{\nu,1} \sim \left( \frac{M_*}{M_P} \right)^2 \hat{\lambda}_{\nu}, \quad 
y^{ij}_u \sim y^{ij}_d \sim y^{ij}_e \sim y^{ij}_{\nu} \sim \left( \frac{M_*}{M_P} \right) \hat{\lambda}_{\nu}, \quad (3)
$$

where for the last equation we supposed that all the higher dimensional couplings present in $W_{MSSM}$ are equal to $\hat{\lambda}_{\nu}$ for all families. Phenomenologically the four dimensional MSSM couplings have to be of the order one $^3$. Therefore it turns out that the higher dimensional coupling is non-perturbative and it has to be $\hat{\lambda}_{\nu} \sim M_P/M_*$. The size of the higher dimensional couplings are completely meaningless. They could be either in the perturbative or in the non-perturbative regime. From the effective field theory point of view what is important is the size of the four dimensional couplings. However for naturalness we impose that in the higher dimensional theory all the couplings are of the same order

$$
\hat{\kappa} \sim \hat{\lambda}_N \sim \hat{\lambda}_H \sim \hat{\lambda}_{\nu}, \quad (4)
$$

$^3$The family hierarchy in the Yukawa sector can be generated through some family symmetry very well explored in the literature.
and additionally if we also require, as experimental fact, that at least the Yukawa coupling for the third generation (defined in $W_{MSSM}$) have to be of order one, one gets

$$\kappa \sim \lambda_N \sim \lambda_H \sim \lambda \equiv \left( \frac{M_s}{M_P} \right)^2, \quad \lambda_\nu \sim \left( \frac{M_s}{M_P} \right).$$

(5)

Before getting into some more technical details of the model, let us explain the physical motivations for considering each term of the superpotential (1).

1. **Inflation**: The term $\kappa \phi N^2$ will define the hybrid inflationary potential where $\phi$ is the inflaton which slow rolls in a semi-flat potential while the waterfall field $N$ is set to zero. Once the inflaton field takes some value below a critical point given by the supersymmetric breaking sector of the model, inflation would end and the waterfall field develops an expectation value $\langle N \rangle$.

2. **The $\mu$ problem**: The term $\lambda N h_u h_d$ will provide a higgsino mass once inflation ends given by $\mu = \lambda_N \langle N \rangle$ like in the Next to Minimal Supersymmetric Standard Model (NMSSM).

3. **Right handed neutrino masses**: We assume that the lightest right-handed neutrino gets Majorana mass through the non-renormalizable operator in Eq.(1) $\lambda_\nu 1^2 (\nu^c_{R,1})^2 / M_*$, where the Yukawa coupling is suppressed due to the fact that the operator contains two bulk fields $N$, and is given by $\lambda_\nu = \lambda_\nu \sim O(M_s/M_P)$. It is this lightest right-handed neutrino $\nu^c_{R,1}$, that we shall henceforth simply refer to as $\nu^c_1$, that will play the role of the coupled curvaton, although it may be completely subdominant in the see-saw mechanism, only contributing to the lightest physical neutrino mass $m_1$, which may be vanishingly small, leading only to an upper bound on its Yukawa couplings.

4. **$G_A \times G_B$ symmetry breaking**: The vevs of $H_1$ and $H_2$ which transform under $G_A \times G_B$ as $H_1 = (\bar{R}, R)$ and $H_2 = (R, \bar{R})$ mediate the breaking of the group $G_A \times G_B$ down to the SM gauge group. With $G_A \times G_B$ unbroken, any $F(D)$-flat direction would be protected against radiative corrections during inflation arising from either Yukawa or gauge interactions. For example, during inflation the Brane I soft masses will be smaller than the Hubble constant (see Appendix B), this means any $F(D)$-flat direction would satisfy automatically the slow roll conditions. However, well after inflation ends the Higgses $H_1$ and $H_2$ get a vev and the Brane I soft masses turn out to be of the order $M_{SUSY}/(4\pi)$ due to bulk particles propagating inside a loop with $M_{SUSY}$ masses, i.e. gaugino mediation [10].

In order to specify completely the superpotential (1) we have to impose a global $U(1)_{PQ}$ Peccei-Quinn symmetry in such a way undesirable terms like $N^3$, $\phi^3$, $\phi h_u h_d$ and so on are
forbidden. Under this global symmetry the fields have the following charges:

\[ Q_N + Q_{h_u} + Q_{h_d} = 0, \quad Q_\phi + 2Q_N = 0, \quad Q_\phi + Q_{H_1} + Q_{H_2} = 0, \quad Q_N + Q_{\nu_{R,1}} = 0. \] (6)

The global symmetry, \( U(1)_PQ \) also forbids explicit RH Majorana neutrino masses in the superpotential, but \( B - L \) symmetry is broken by the non-renormalizable term \( \lambda_{\nu,1}N^2(\nu_{R,1})^2/M_s \). The global symmetry is broken at the scale of the scalar singlet VEV’s releasing a very light axion and providing as consequence an axionic solution to the CP-strong problem. In the next section will discuss what is precisely the axion scale \( f_a \).

3 The potential

Now we are ready to study in detail the scalar potential of our model. We will write the potential along the D-flat directions in both Higgs sectors, \( h_u = h_d = h \) and \( H_1 = H_2 = H \), and comment later on symmetry breaking. Also we will take the coefficients \( c_i \) for the vacuum energy Eq. (36) and for the soft bulk masses Eq. (37) given in Appendix B equal to one for simplicity, with

\[ V_0 \sim F_S^2, \]  
\[ m^2 \equiv \left( \frac{F_S}{M_P} \right)^2 \sim m_\phi^2 \sim m_h^2 \sim m_N^2 \sim m_H^2, \]  
\[ A \equiv \frac{F_S}{M_s} = \frac{A_{\lambda_H}}{c_H} = \frac{A_{\lambda_N}}{c_N} = \frac{A_\kappa}{c_\kappa}. \] (9)

With this simplifications the scalar potential for the real components of the fields, at energy below \( M_s \) looks like

\[ V = V_0 + m^2 \left( \phi^2 + h^2 + N^2 + H^2 \right) + 2\lambda A \left( c_H \phi H^2 - c_\kappa \phi N^2 + c_N Nh^2 \right) \]
\[ + \lambda^2(h^2 - 2\phi N)^2 + 2\lambda^2 N^2 h^2 + \lambda^2(H^2 - N^2)^2 + 2\lambda^2 \phi^2 H^2 \]
\[ + 4\lambda^2 \frac{N^4 H^4}{M^2_s} + 4\lambda^2 \frac{N^2 v_R^4}{M^2_s} + 4\lambda^2 \lambda \frac{N^2 v_R^4}{M_s} (h^2 - 2\phi N), \] (10)

where the couplings \( \lambda \) and \( \lambda_{\nu} \) are those given by (5), and the first line of the above equation are the soft susy breaking terms.

Neglecting the \( m^2 \)-term since \( m \ll A \), the global minimum of the potential (10) is given by

\[ \langle \phi \rangle = \frac{c_\kappa A}{4\lambda}, \]  
\[ \langle N \rangle = \frac{c_\kappa A}{2\sqrt{2}\lambda}, \]  
\[ \langle H \rangle = \langle h \rangle = \langle \tilde{\nu}_R \rangle = 0. \] (13)
where all \( c_i \approx \mathcal{O}(1) \). These coefficients should satisfy \( c_H > c_N/4 \) \( (c_N > (2 - \sqrt{2})c_N/8) \), in order to stabilize the vev for \( h \) \( (H) \) at zero. On the other hand, there is no minimum with \( \langle H \rangle \neq 0 \) and \( \langle H \rangle \sim \langle \phi \rangle \sim \langle N \rangle \). The solutions of the the minimization equations with \( H \neq 0 \) are a maximum of the potential instead.

Using Eqs. (5) and (9) the VEV are approximately

\[
\langle \phi \rangle \sim \langle N \rangle \sim \frac{A}{\lambda} = \frac{F_S}{M_*} \left( \frac{M_P}{M_*} \right)^2.
\]  

(14)

All parameters (couplings and mass terms) of our potential (10) are functions of just two free parameters\(^4\), \( M_* \) and \( F_S \). However, one of them can be eliminated by imposing zero cosmological constant around the global minima of the scalar potential, \( V(\langle \phi \rangle, \langle N \rangle) = 0 \), and therefore \( F_S \) can be expressed as a function of \( M_* \),

\[
F_S = \frac{M_*^4}{M_P^2}.
\]  

(15)

With this choice for \( F_S \) all the different scales involved in our model are function of just one scale \( M_* \). The soft masses (Eqs. (8) and (9)) can be rewritten as

\[
m^2 = M_*^2 \left( \frac{M_*}{M_P} \right)^6, \quad A = M_* \left( \frac{M_*}{M_P} \right)^2.
\]  

(16)

In general, since \( M_* < M_P \) we have \( m/A = M_*/M_P \ll 1 \), and \( \lambda \simeq (M_*/M_P)^2 \ll 1 \). Plugging Eq. (15) into (14) we found that the VEV of the scalar fields are degenerated at the higher dimensional cutoff scale \( M_* \)

\[
\langle \phi \rangle \sim \langle N \rangle \sim M_*.
\]  

(17)

The \( H_1, H_2 \) fields will develop later a much smaller vev by a similar mechanism to the radiative electroweak symmetry breaking in the Higgs sector. Extending the matter content of the model on Brane I by two pairs of fermions \( F_1, \bar{F}_1 \) and \( F_2, \bar{F}_2 \), in conjugate representations, they will couple respectively to the Higgses as \( H_1 F_1 \bar{F}_2 \) and \( H_2 F_2 \bar{F}_1 \) with Yukawa couplings of order one (same order of magnitude as the top Yukawa coupling). Radiative corrections due to this large Yukawa coupling will render one of the \( H_1, H_2 \) masses negative, lifting the D-flat direction and allowing them to get a vev. We notice that up to this point all the matter fields on Brane I are massless. The only massive fields are those living in the bulk.

\(^4\)The reduced four dimensional Planck scale is fixed at \( M_P = 2.4 \times 10^{18} \) GeV by gravity.
4 The question of scale

In this section we shall address the numerical question: what are the correct sizes of $m^2$ and $A$ in order to reproduce a good phenomenology? In other words, how large is the single mass scale $M_*$ in the theory? The physical requirement that one of the scales $m$ or $A$ is precisely the Electroweak scale will fix uniquely the value of $M_*$, and consequently the remaining scales. We shall see that we are led to the conclusion that $M_*$ must be identified with an intermediate mass scale.

4.1 The electroweak scale

Chiral matter do not directly feel the breaking of SUSY which takes place in the “hidden” 3-brane sector. The effects for the chiral matter of SUSY-breaking are only transmitted through the influence of bulk fields, which are the only ones which can move into the bulk spacetime and couple to both kinds of 3-brane sectors. As we have seen from the Fig. (1) the bulk fields are the inflaton, the Higgses, the waterfall field and gauginos belonging to the gauge group, $G_A$. Their soft masses are equal to what we have called $m$ in (8). So far we have not mentioned gravity in this paper. It is widely believed that gravity is propagating in the bulk in which case the gravitino mass would be $m_{3/2} \sim m$. With this information in mind we could think that the most natural selection for $m$ would be the Electroweak Scale, $m_{3/2} \sim m \sim M_{SUSY} \sim$ TeV. However this is not possible because the other scale involved, $A$, would be $A = (M_P / M_*)$ TeV and it can be as much as $A \sim 10^3$ TeV. On the other hand, $A$ is the scale associated with gauginos living in the SUSY breaking brane. When the full group $G_A \times G_B$ breaks down to the SM group, it turns out that SM gauginos would be as heavy as $10^3$ TeV which we regard as phenomenologically unacceptable.

The other possibility (the only one) is choose the scale $A$ as the Electroweak scale. Fixing $A \sim$ TeV and using Eq. (16) we have that the scale $M_*$ (the fundamental scale in higher dimensions) has to be

$$M_* \approx 10^{13} \text{GeV},$$

(18)

As a consequence the SUSY breaking scale is $F_S^{1/2} \sim 10^8$ GeV and the $m$-term is $m \sim 10$ MeV. In the next section we will see that $m$ will give us the inflaton mass. Using Eq. (39) in Appendix B, the Hubble expansion parameter during inflation turns out to be of the order MeV.

Some of the phenomenological benefits of an intermediate scale have been noted in [12,13]. Below we carefully examine these issues relevant to the present model.
4.2 The $\mu$-scale

A problem of the Minimal version of the Supersymmetric Standard Model (MSSM) is why the $\mu$ term which is a supersymmetric mass has to be of the same order of the soft terms, as required to get an acceptable phenomenology. In the other words, what is the origin of the $\mu$-scale? There are many solutions to this problem \footnote{Note that the Giudice-Masiero mechanism \cite{14} presents also a solution of the $\mu$ problem within the MSSM by generating the $\mu$ term via a non-minimal Kahler potential.}, for example in the Next to Minimal model (NMSSM) the $\mu$-parameter is replaced by a trilinear coupling involving an extra field $N$ and the higgses, $\lambda N H_u H_d$. Once $N$ gets a vev, the $\mu$-term is generated.

The solution of the $\mu$-problem in our model relies on the same mechanism as in the NMSSM. However, there are many features that make our model different to the usual NMSSM model. The usual NMSSM involves a term like $\kappa N^3$ in the superpotential so that the model has an exact $Z_3$ symmetry \cite{15,16} which is broken at the weak scale (at the scale of $\langle N \rangle$) leading to a serious domain wall problem\footnote{See for example \cite{19} for an alternative solution to the domain wall problem based on a $Z_2$ $R$-symmetry.} \cite{18}. In our model a global $U(1)_{PQ}$ symmetry forbids such cubic terms so there is no domain wall problem. As we will see in the next subsection the global $U(1)_{PQ}$ symmetry is linked with the solution to the CP - strong problem. In fact the singlet field $N$ in our model plays three rôles. It switches on the $\mu$-term once it gets a vev

\[
\mu = \lambda \langle N \rangle \approx A \approx M_* \left( \frac{M_*}{M_P} \right)^2 \approx \text{TeV},
\]

It plays the rôle of a waterfall field of hybrid inflation, ending inflation through a phase transition, as discussed in Sec. 5. And its vev is responsible for generating the right-handed neutrino mass scale.

4.3 The axion scale

The most elegant explanation of the strong CP problem is provided by the Peccei-Quinn (PQ) mechanism \cite{20}, in which the CP violating angle $\bar{\Theta}$ is set to zero dynamically as a result of a global, spontaneously broken $U(1)_{PQ}$ Peccei-Quinn symmetry. The corresponding Goldstone mode of this symmetry is the axion field and the static $\bar{\Theta}$ parameter is substituted by a dynamical one, $a(x)/f_a$, where $a(x)$ is the axion field and $f_a$ is a dimensionful constant known as the axion decay constant.

In our model the $U(1)_{PQ}$ Peccei-Quinn symmetry is spontaneously broken once the scalar fields charged under $U(1)_{PQ}$ (see Eq. (6)) get a VEV of the order $M_*$. This implies automatically
that the axion decay constant is
\[ f_a \sim M_* \sim 10^{13} \text{GeV} \, . \] (20)

On the other hand, the axion also has interesting cosmological implications, especially as a cold dark matter candidate. Indeed coherent oscillations around the minimum of its potential may dominate the energy density of the universe if its potential is very flat. This puts an upper bound for \( f_a \) of order \( f_a \leq 10^{12} \text{GeV} \). It seems that our prediction is a little bit higher that the allowed by experiments. However, as has been pointed out in [21], \( f_a \) can be as big as \( 10^{15} \text{GeV} \) in models where the reheating temperature is below a GeV, that is, below the temperature at which the axion field begins to oscillate. The point is that during inflation the PQ symmetry is broken and the axion field is displaced at some arbitrary angle, and it relaxes to zero only after reheating and only below the QCD phase transition when its potential is tilted. At this point the dangerous energy stored in the axion field is released, but if the reheating temperature is of order a GeV then the resulting axion density from the displaced axion field will be diluted by the entropy release produced by the inflaton decay. In the Refs. [4,17] has been showed that the reheating temperature for the model under consideration is around the GeV scale and therefore an axion decay constant of the order \( 10^{13} \text{GeV} \) may be consistent cosmological constraints.

4.4 The right-handed neutrino mass scale

Neutrino oscillation phenomenology requires that there must be two further heavy right-handed neutrinos with a Majorana mass arising from the renormalisable operator \( \lambda_{\nu,i} N (\nu_{R,i})^2 \) where \( \lambda_{\nu,i} \sim 1 \), and \( i = 2, 3 \). We have not included these operators in the superpotential in Eq. (1) because these heavy right-handed neutrinos play no rôle in cosmology, but such operators may readily be included by suitable choice of PQ charges for the second and third right-handed neutrinos \( \nu_{R,i}, \, i = 2, 3 \). The heaviest right-handed neutrinos of mass \( \lambda_{\nu,2} \langle N \rangle \) and \( \lambda_{\nu,3} \langle N \rangle \) will give the dominant contribution to the solar and atmospheric neutrino masses of the order of \( m_2 \approx y_{\nu,2}^2 v^2 / \lambda_{\nu,2} \langle N \rangle \), and \( m_3 \approx y_{\nu,3}^2 v^2 / \lambda_{\nu,3} \langle N \rangle \), respectively, where the mild hierarchy \( m_2 \ll m_3 \) can be achieved by suitable choices of Yukawa couplings above [11]. On the other hand the lightest right-handed neutrino which plays an important rôle during inflation, and will explain the amplitude for the curvature perturbation, will play no part in the see-saw generation of atmospheric and solar neutrino masses, but will generate the lightest physical neutrino mass \( m_1 \). The lightest right-handed neutrino mass given by \( \lambda_{\nu,1} \langle N \rangle \), where now \( \lambda_{\nu,1} \sim M_* / M_P \), will contribute to the lightest physical neutrino mass \( m_1 = y_{\nu,1}^2 v^2 / (\lambda_{\nu,1} \langle N \rangle) \). A hierarchy in the
neutrino sector \((m_1 \ll m_2)\) is trivially achieved when \(y_{\nu,1} \ll (M_*/M_P)^{1/2} y_{\nu,2} \sim 3 \times 10^{-3} y_{\nu,2} \). As we will see in the next section we need such small Yukawa coupling for \(y_{\nu,1}\) in order to stabilize the preheating effect due to the oscillations of the lightest right-handed neutrino. Since \(m_1\) can be arbitrarily light, the lightest right-handed neutrino responsible for its mass can be effectively decoupled from the see-saw mechanism due to its highly suppressed Yukawa coupling. The scenario described above is familiar in neutrino phenomenology and is known as the sequential dominance mechanism [11].

To summarize, a hierarchical neutrino mass scheme, where \(m_3 \sim 0.05\) eV, assuming \(y_{\nu,3} \sim 0.1 - 0.5\), led to a right-handed neutrino mass scale of the same order as the axion scale,

\[
\langle N \rangle \sim M_* \sim 10^{13}\text{GeV}
\]  

Therefore in our model we are led to assign the string scale to the intermediate scale \(M_* \sim 10^{13}\) GeV by four independent pieces of physics: electroweak symmetry breaking; the \(\mu\) parameter; the axion scale; and the atmospheric neutrino mass scale.

## 5 The lightest right-handed sneutrino as a coupled curvaton

In a previous paper we suggested that the Higgs fields of the supersymmetric standard model could play the rôle of a coupled curvaton [5] within this class of models, and we discussed how Higgs perturbations could be converted into the total curvature perturbations during the first stages of reheating. At first we assumed that the curvature Higgs contribution does not change after horizon crossing, and we obtained the desired curvature perturbation for a Higgs VEV value \(h_* \sim 1\) TeV. However we later found that, once the fields get coupled during the phase transition, the evolution of the field fluctuations will be affected suppressing the amplitude of curvature perturbation of the Higgs field relative to its value at horizon crossing. Subsequently we showed that this suppression could be compensated by taking into account preheating or parametric resonance effects [34] which can enhance the value of the curvature perturbation to the desired value [7].

In this section we propose and study the possibility that the the rôle of the coupled curvaton is instead played by the lightest right-handed sneutrino, with the inflaton identified as the field \(\phi\), as before. The \(N\) field and the Higgs fields will here be assumed to be zero during inflation. The associated lightest right-handed neutrino \(\nu_{R,1}\) shall simply be referred to as \(\nu_R\), its Majorana
Yukawa coupling as $\lambda_\nu = \lambda_\nu,1$, and its Yukawa coupling to left-handed leptons as $y_\nu = y_\nu,1$, for ease of notation.

According to the potential in Eq. 10, there is a flat direction for both $\tilde{\nu}_R$ and $\phi$, while $N$ field is held at zero for values of the inflaton field $\phi$ larger than the critical value

$$\phi_c = \frac{c_\kappa A + \sqrt{c_\kappa^2 A^2 - 16m^2}}{4\lambda} \approx c_\kappa \frac{A}{2\lambda}.$$  \hspace{1cm} (22)

The inflationary epoch is therefore described by a slowly rolling inflation field $\phi$ and a slowly rolling light right-handed sneutrino field $\tilde{\nu}_R$ (recall that the large mass of the right-handed sneutrino is generated by the VEV of the $N$ field, which is zero during inflation). As long as $\phi > \phi_c$, the $N$ field dependent squared mass is positive and then $N$ (as well as $H$) is trapped at the origin; the potential energy in Eq. (23) is then dominated by the vacuum energy $V_0$, and the potential (10) simplifies to:

$$V = V_0 + m_\phi^2 \phi^2 + m_\nu^2 \tilde{\nu}_R^2,$$  \hspace{1cm} (23)

with $V_0 \approx F_S^2 \approx (10^8 \text{GeV})^2$, and $m_\alpha \sim c_\alpha m$. The slow roll conditions are given by:

$$\epsilon_N = \frac{M_P^2 m_\phi^2 \phi_N^2}{V_0^2} < 1,$$  \hspace{1cm} (24)

$$|\eta_N| = \frac{M_P^2 m_\phi^2 \eta_N^2}{V_0} < 1,$$  \hspace{1cm} (25)

where the subscript $N$ means $N$ e-folds before the end of inflation. Using Eqs. (36) and (37) we then get $|\eta_N| = c_\phi/c_V$, and $\epsilon_N \ll \eta_N$, and slow-roll only requires $c_\phi/c_V < 1/3$.

The amplitude of the spectrum of the (comoving) curvature perturbation $\mathcal{R}$, generated by the inflaton field is given by [24]:

$$P_R^{1/2} \approx \left( \frac{H_*}{\phi_*} \right) \left( \frac{H_*}{2\pi} \right) \approx \left( \frac{H_*}{2\pi \eta_N \phi_*} \right),$$  \hspace{1cm} (26)

where the subscript “*” denotes the time of horizon exit, say 60 e-folds before the end of inflation. The value of the inflaton field during inflation is around the cutoff of the theory, $\phi_* \sim \phi_c \sim M_*$, as usual in SUSY inflation [38], while the Hubble parameter is of the order of $H \sim M_* (M_*/M_P)^3$. Therefore

$$P_R^{1/2} \approx \left( \frac{M_*}{M_P} \right)^3,$$  \hspace{1cm} (27)

which for $M_* \approx 10^{13}$ GeV is quite below the COBE value $P_R^{1/2} = 5 \times 10^{-5}$ [23].
However, the quantum fluctuations of any light field during inflation, i.e., the Higgs $h$ and the lightest right-handed sneutrino $\nu_R$, will contribute to the total curvature perturbation$^7$, $\mathcal{R}$, such that [25,26]

$$\mathcal{R} = \sum_{\alpha} \frac{\rho_\alpha + P_\alpha}{\rho + P} \mathcal{R}_\alpha,$$

(28)

$\rho_\alpha$ and $P_\alpha$ being respectively the energy density and the pressure for each component, with $\rho_\alpha + P_\alpha = \dot{\phi}_\alpha^2$, $\rho$ and $P$ the total energy density and pressure, $\mathcal{R}_\alpha$ the curvature perturbation generated by each field,

$$\mathcal{R}_\alpha \simeq H \frac{Q_\alpha}{\dot{\phi}_\alpha},$$

(29)

and $Q_\alpha$ the gauge invariant quantum fluctuations of the field [27]. Given the model parameters, we know that the inflaton field has a background value of the order of the cut-off scale $M_*$, both during inflation and at the global minimum, but the value of the sneutrino field is arbitrary during inflation. At the global minimum it will relax to zero. Given that, we may assume that even during inflation the sneutrino field is not far from its global minimum value, and therefore $\phi \gg \nu_R$. From this condition it follows that $\dot{\phi} \gg \dot{\nu}_R$, but $\mathcal{R}_\phi \ll \mathcal{R}_\nu$. The total curvature perturbation in Eq. (28) is dominated by the field with the largest kinetic energy, and during inflation this is just the inflaton field, which as we have seen gives rise to a too small contribution. Nevertheless, the COBE normalization of the spectrum constraints its value at the onset of the radiation dominated era, after inflation and the reheating process is complete. In single field models of inflation, the total curvature perturbation on large scale remains practically constant after horizon crossing, and it is enough to estimate the spectrum at that point. On the other hand, in a multifield inflationary model (or in general in a multi-component Universe) we have both adiabatic (total) and entropy or isocurvature perturbations. The latter are given by the relative contributions between different components, $\mathcal{R}_\alpha - \mathcal{R}_\beta$. Entropy modes can seed the the adiabatic one, i.e., the total curvature perturbation, when their contribution to the total energy density becomes comparable [28–33]. This is what we expect at the end of inflation, when all the fields move fast toward the global minimum. At this point, the energy densities of the fields become comparable, and the total curvature perturbation, Eq. (28) may become of the order of $\mathcal{R}_\nu$.

Here we would like to consider instead the lightest right-handed sneutrino as the main source of the isocurvature perturbation during inflation. The Yukawa coupling of the lightest right-

$^7$We have dropped the reference to the wavenumber in the curvature perturbation but it is implicitly assumed that we only refer to large scale perturbations, with $k \ll H a$. 

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handed sneutrino to the $N$ field is a factor $(M_P/M_*)$ larger than the coupling $\kappa$ between $\phi$ and $N$. Hence, the right handed sneutrino acquires a mass in the global minimum which is larger than the other particles by the same factor. Nevertheless, due to the coupling with the other fields, it may oscillate together with $N$ and $\phi$ but with a smaller amplitude $\nu_R \sim (\kappa/\lambda_N)M_*$. In any case, during the oscillations the energy density of the 3 fields become comparable, with $M_{RR}\nu_R \sim (\kappa N) N$.

Previous to that we have to consider the effect of the phase transition on the perturbations of the fields. The tachyonic instability in the $N$ direction makes the field grow until it reaches the straight-line trajectory in field space, $N = \sqrt{2}(\phi - \phi_c)$. On the other hand the sneutrino still follows undisturbed its inflationary trajectory for a while. With respect to $Q_N$ perturbations, they also feel the tachyonic instability, they grow as $Q_N \propto \dot N$ and we end up with $Q_N = -\sqrt{2}Q_\phi$ along the straight-line trajectory. At this point no change has been induced in the total curvature perturbation. The fact that $Q_N$ is non negligible after the phase transition does not mean that we are generating additional entropy modes. Along the straight line trajectory we still have only one degree of freedom, the adiabatic mode [26].

Soon after that, the lightest right-handed sneutrino starts feeling the presence of the other fields. The effect on its perturbation is the same than we found for the Higgs in Ref. [7]: when the field gets coupled the $R_\nu$ perturbation is dragged toward $R$, and by the time they first reach the global minimum all the contributions $R_\alpha$ are comparable but some orders of magnitude smaller than the initial $R_\nu$.

From this point of view, during the phase transition the initial entropy perturbation is only partially converted into the adiabatic one. In other words, entropy perturbations are suppressed due to the tachyonic instability during the phase transition. And in particular for our model values, we are still some orders of magnitude below the COBE normalization. Again, once the oscillations begin, the presence of a third field, in this case the sneutrino, will curve the initial straight-line trajectory in the $N\phi$ plane, giving rise to preheating of the large scale perturbations. In this scenario this will happen before the fields reach the global minimum. We have the inflaton field decreasing from its value at the critical point, while $N$ and $\tilde{\nu}_R$ are growing and moving faster than $\phi$. First $N$ gets destabilized, and then $\tilde{\nu}_R$. But for the values of the fields such that $\lambda_N \tilde{\nu}_R^2/M_* > \kappa(\phi_c - \phi)$, there is an approximate minimum with $N \simeq \tilde{\nu}_R$ and $\lambda_N \tilde{\nu}_R^2/M_* \simeq \phi \simeq \phi_c$, around which $N$ and the RH sneutrino oscillates. It is already during this period that the large scale perturbations are preheated. First, $N$ and $\tilde{\nu}_R$ do not oscillate in phase, and in addition there is a tachyonic instability in the $N$ and $\tilde{\nu}_R$ squared masses during
the oscillations, which enhances the resonance [35]. This effect only lasts a short period of time, until the inflaton field is close enough to the global minimum to drag towards it the other fields.

At the same time we may also preheat fluctuations on smaller scales\(^8\), which will back-react on the system shutting down the resonance. But we do not expect this to happen until the inflaton is also oscillating, due to the small amplitude of the previous \(N - \tilde{\nu}_R\) oscillations. Another effect to be considered is the decay of the heavy sneutrino, which happens earlier than the decay of the singlets \(\phi\) and \(N\). Once the lightest right-handed sneutrino (the source of the entropy perturbation) disappears the resonance on the large scales will stop.

As an example, in Fig. (2) we have plotted the evolution of the background fields (LHS plot), and the different field contributions to the curvature perturbation \(\mathcal{R}_\alpha\) (RHS plot), during the first oscillations of the fields, after they have passed the critical point. The value of the lightest right-handed sneutrino field during inflation is \(\nu_R = 1\) TeV, which gives \(\mathcal{R}_{\nu_R} \simeq 10^{-5}\)

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\(^8\)Without the sneutrino and/or the Higgs, there is indeed a strong parametric resonance (tachyonic preheating [35]) for the fluctuations with a wavenumber up to \(O(\kappa N_0)\) [36] during the first 3 oscillations of the fields, before it enters in a narrow resonance regime.
at horizon-crossing. Nevertheless, as it can be seen in the plot, after the phase transition all contributions are roughly equal. For the numerical integration, we have chosen $\lambda_\nu = 2 \times 10^5 \kappa$, and we have also introduced a decay rate for the lightest right-handed sneutrino, $\Gamma_\nu \simeq y_\nu^2 M_{RR}$, with $y_\nu \simeq 7 \times 10^{-3}$. This gives rise to the right order of magnitude for the amplitude of the curvature perturbation (slightly larger than COBE). The main enhancement is produced during the oscillations of the fields $N$ and $\tilde{\nu}_R$, and the resonance ends when $\tilde{\nu}_R$ goes to zero and we are left only with $\phi$ and $N$ oscillating.

The decay of the lightest right-handed sneutrino gives rise to radiation, with $\rho_R \simeq 0.2 \rho_\phi$, where $\rho_\phi$ refers to the energy density in the oscillating singlets. Having converted a fraction of the initial vacuum energy into radiation, the background fields still oscillate (in phase) but with smaller amplitudes compared to the case without the sneutrino, which may not be large enough to allow the preheating of the small scale fluctuations. Nevertheless one has to bear in mind that larger decay rates for the sneutrino means a shorter resonance for the curvature perturbation, or no resonance at all. On the other hand, for a smaller decay rate the resonance will then be shut down by the backreaction of the small scale modes.

## 6 Conclusion

We have proposed a model of cosmology and particle physics in which all relevant scales are derived from an intermediate string scale $M_\ast \sim 10^{13}$ GeV, identified with both the Peccei-Quinn symmetry breaking axion scale $f_a$ and the heaviest right-handed neutrino mass scale $M_{RR}$. A supersymmetry breaking scale is derived from the constraint of having a small cosmological constant leading to $F_S^{1/2} \sim M_\ast^2 / M_P \sim 10^8$ GeV. The $\mu$ parameter of the MSSM and the electroweak breaking scale are then given by $F_S / M_\ast \sim 10^3$ GeV. The model involves hybrid inflation, with the inflaton mass given by $F_S / M_\ast \sim \text{MeV}$, and their Yukawa couplings given by $(M_\ast / M_P)^2 \sim 10^{-10}$. In our model we were led to assign the string scale to the intermediate scale $M_\ast \sim 10^{13}$ GeV by four independent pieces of physics: electroweak symmetry breaking; the $\mu$ parameter; the axion scale; and the neutrino mass scale. In recent years models with an intermediate string scale ($10^{11} < M_\ast < 10^{14}$ GeV) have been seen of great interest because contain many phenomenological issues (see for example [12,13]). The novelty of our construction is we give an explicit potential (10) where the link between particle physics and cosmology is explicit. Also

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9The decay of the lightest right-handed sneutrino is also expected to generate leptogenesis, and cold dark matter can arise from preheating of neutralinos as discussed in [40].
the intermediate scale appears as a consequence of requiring zero cosmological constant at the
global minimum of the potential together with requiring one of the mass parameters involved
in the potential to be the electroweak scale. In addition, if all the dimensionless couplings are
equal in the higher dimensional model, we found that the effective couplings are just powers of
the ratio $M_*/M_P$.

The large scale structure in the Universe is generated by the lightest right-handed sneutrino
playing the rôle of a coupled curvaton. We showed that it is possible to obtain the correct
curvature perturbation by including the effects of preheating and adjusting neutrino Yukawa
couplings to be rather small, leading to a (lightest) right-handed neutrino mass of about $10^{8}$ GeV.
This is possible because of sequential dominance [11] since in this case the lightest right-handed
neutrino plays no essential rôle in the see-saw mechanism. The sneutrino as a coupled curvaton
only strengthens the synthesis between cosmology and particle physics at the intermediate scale.

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Appendix A: Yukawa Couplings

In higher dimensions the superpotential is well defined just in one of the fixed points (Brane I
or Brane II) \(^{10}\). In general the Lagrangian in higher dimensions is given by

$$\mathcal{L}_{4+d} = \int d^2 \theta \, \hat{W}_{4+d} \left( \delta^d(0) + \delta^d(y_i - \pi R) \right),$$

(30)

where the superpotential $\hat{W}_{4+d}$ is a function of bulk fields ($\hat{\Psi}_i$) and brane fields ($\Phi_j$), given by:

$$W_{4+d} \equiv \hat{\lambda}_{i,j} \frac{\hat{\Psi}_i^a \Phi_j^\beta}{M_*^{a+\beta-3}}.$$

(31)

Here we have introduced a mass scale, $M_*$, in such a way the couplings $\hat{\lambda}_{i,j}$ remain dimensionless.
This scale is actually the Planck scale in higher dimensions, or the string scale in string theories,
which is related with the four dimensional Planck scale, $M_P$, through the well known formula

$$M_P^2 = M_*^{2+d} R^d.$$

(32)

\(^{10}\)Notice that this is always the case if the coupling involves bulk and brane fields. However by construction
also couplings which involve just bulk fields have to be defined in one of the four dimensional boundary due to
an enhancement of the number of supersymmetries contained in the bulk.
Using the fact that the bulk fields contain a volume suppression factor with respect to their zero mode (the effective four dimensional field), i.e., $\hat{\Psi}_i = \Psi_i/R^{d/2}$, we can express the four dimensional effective superpotential, after integrating out the extra $d$ dimensions, as

$$W_4 = \int d^d y \hat{W}_{4+d} \left( \delta^d(0) + \delta^d(y_i - \pi R) \right)$$

$$= \hat{\lambda}_{i,j} \left( \frac{M_s}{M_P} \right)^\alpha \frac{\Psi^\alpha \Phi^\beta}{M_s^{\alpha + \beta - 3}}. \tag{33}$$

Redefining the effective four dimensional coupling as

$$\lambda_{i,j} = \left( \frac{M_s}{M_P} \right)^\alpha \hat{\lambda}_{i,j}, \tag{34}$$

we notice that in general we get an important suppression factor if $M_s \ll M_P$. In other words, if there is a fundamental intermediate scale, $M_s$, defined in the higher dimensional theory, we always get small couplings, $\lambda_{i,j} \ll 1$, even though the higher dimensional couplings, $\hat{\lambda}_{i,j}$, can be of order one or even non-perturbative as it would be in our particular case.

Finally note that following the same procedure in the Kahler potential will lead to no additional volume effects. For example consider a canonically normalized $4 + d$-dimensional Kahler potential term of the form $\hat{\Psi}_i \hat{\Psi}_i^\dagger$ then the corresponding 4-dimensional Kahler potential will contain the term $\Psi_i \Psi_i^\dagger$ which maintains its canonical form due to a cancellation of the volume factors. There may be some small additional corrections due to canonical normalization effects [39], but these will not affect the analysis here.

### Appendix B: Soft Masses

We shall suppose that SUSY is broken by the $F_S$-term of a four dimensional gauge-singlet field $S$ localized at the Brane II ($y_i = \pi R$) and mediated across the extra dimensional space to the Brane I by bulk fields propagating in a loop correction like gaugino mediation [10]. The SUSY breaking Lagrangian contains six terms that lead to the bulk gaugino mass ($M_A$), Brane II gaugino mass ($M_B$), vacuum energy ($V_0$), soft masses for bulk fields ($M_{\Psi_i}$), soft masses for Brane II fields ($M_{\Phi_i}$) and trilinear soft terms ($A_{\lambda_{i,j}}$). In general we have

$$L_{4+d}^{\text{soft}} = \delta^d(y_i - \pi R) \left( c_{\lambda_{i,j}} \int d^2 \theta \frac{\hat{W}_{4+d} S}{M_s} + c_{\Psi_i} \int d^4 \theta \frac{\hat{\Psi}_i \hat{\Psi}_i S^\dagger S}{M_s^{2+d}} + c_{\Phi_i} \int d^4 \theta \frac{\Phi^\dagger \Phi_i S^\dagger S}{M_s^2} \right)$$

$$+ c_V \int d^4 \theta S^\dagger S + c_A \int d^2 \theta \frac{\hat{W}_{(A)}^\dagger \hat{W}_{(A)} S}{M_s^{1+d}} + c_B \int d^2 \theta \frac{W_{(B)}^\dagger W_{(B)} S}{M_s}, \tag{35}$$

11This suppression factor only would depend on the number $\alpha$ of fields living in the bulk and it will be completely independent of the number $d$ of extra dimensions.
where $W^{(A)}_a$ ($W^{(B)}_a$) is the field strength of the gauge group $G_A$ ($G_B$) and the constants $c_i = \{c_{\lambda_{ij}}, c_{\Psi_i}, c_{\Phi_i}, c_V, c_A, c_B\}$ are of the order one. These constants are completely model-dependent, for example, the non-renormalizable terms (35) might come from integrating out some modes of mass $M_*$ propagating in a loop process such that $c_i = \mathcal{N}/(4\pi)^2$, being $\mathcal{N}$ the massive mode’s degree of freedom. It is straightforward to show that in the effective four dimensional theory we get different mass scales associated with bulk fields and branes fields as along as $M_* \ll M_P$. These are given by

\begin{align}
V_0 &= c_V F_S^2 \\
M_A &= c_A F_S/M_P \\
M_B &= c_B F_S/M_* \\
M_{\Psi_i}^2 &= c_{\Psi_i} \left( F_S/M_P \right)^2 \\
M_{\Phi_i}^2 &= c_{\Phi_i} \left( F_S/M_* \right)^2 \\
A_{\lambda_{ij}} &= c_{\lambda_{ij}} F_S/M_*.
\end{align}

(36)\hspace{1cm} (37)\hspace{1cm} (38)

The vacuum energy ($V_0$) would dominate the total energy density during inflation providing the typical expansion rate (the Hubble constant) as

$$H^2 = \frac{V_0}{3M_P^2} = \frac{c_V}{3} \left( \frac{F_S}{M_P} \right)^2 = \frac{c_V}{3c_{\Psi_i}} M_{\Psi_i}^2.$$  \hspace{1cm} (39)

The bulk mass scale ($M_{\Psi_i}$) would give us, for example, the inflaton mass. Therefore, in order to satisfy the slow roll condition during inflation ($H < M_{\Psi_i}$) it turns out from (39) that we just need some tuning on the constants $c_i$, i.e. $c_V < 3c_{\Psi_i}$. Finally, the Brane II mass scale would define the typical MSSM soft term, $M_{SUSY} \sim M_{\Phi_i} \sim \mathcal{O}$(TeV). Notice that in general we would have $H \ll M_{SUSY}$ which is completely different what happen in the normal four dimensional supersymmetric hybrid inflationary models where $H \sim M_{SUSY}$ [38].

References


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\footnote{As we have said at the beginning of the section one important assumption of our model is that the location of the two branes are already stabilized before inflation takes place which means that only the four dimensional space are inflated away. In the case that all de 4 + d-dimensions feel inflation at the same time, the Hubble constant would be given by $H^2 = V_0/(3M_{\Psi_i}^2)$ instead [37].}


