Model of cosmology and particle physics at an intermediate scale

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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 μ parameter; the axion scale; and the neutrino mass scale. The model involves hybrid inflation with the waterfall field N being responsible for generating the μ 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 seesaw mechanism, consistent with sequential dominance.

I. 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 axion 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 μ 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_*, from which the Planck scale M_P should be derived in terms of the compactification scales. From this single scale M_* one must be able to derive all the other relevant scales in physics, such as the axion scale f_a, the scales of right-handed neutrino masses M_{RR}, the μ 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_*.

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 $M_* \sim 10^{13}$ 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 being 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 twofold. First 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}

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as well as the Peccei-Quinn (PQ) 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 Sec. II 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 Sec. III we discuss the potential and the symmetry breaking aspects of the model. In Sec. IV 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-handed neutrino mass scale and the origin of the μ parameter are all associated with the same string scale. Section V 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 VI 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.

II. THE MODEL

Let us consider two four dimensional boundaries[2] spatially separated along d extra dimensions with a common

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[1]In [8] the right-handed sneutrino plays the rôle of the weakly coupled late-decaying curvaton.

[2]From now on let us use the phrase *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.

radius $R = 1/M_c$. These extra dimensions are compactified on some orbifold that leads at least to 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 supersymmetry (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, ϕ, the waterfall field, N, the minimal supersymmetric standard model (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]). Having done this, first 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

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

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

$$W_{MSSM} = y_u^{ij} Q_i h_u U_j + y_d^{ij} Q_i h_d D_j + y_e^{ij} L_i h_d e_{Rj}^c$$
$$+ y_\nu^{ij} L_i h_u \nu_{Rj}^c. \tag{2}$$

Following the relationship between higher dimensional couplings ($\hat{\lambda}_{i,j}$) and four dimensional one ($\lambda_{i,j}$) given in Appendix A in Eq. (A5), we get

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

$$\lambda_\nu \equiv \lambda_{\nu,1} \sim \left(\frac{M_*}{M_P}\right)^2 \hat{\lambda}_\nu, \quad y_u^{ij} \sim y_d^{ij} \sim y_e^{ij} \sim y_\nu^{ij} \sim \left(\frac{M_*}{M_P}\right) \hat{\lambda}_\nu, \tag{3}$$

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

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[3]The family hierarchy in the Yukawa sector can be generated through some family symmetry very well explored in the literature.
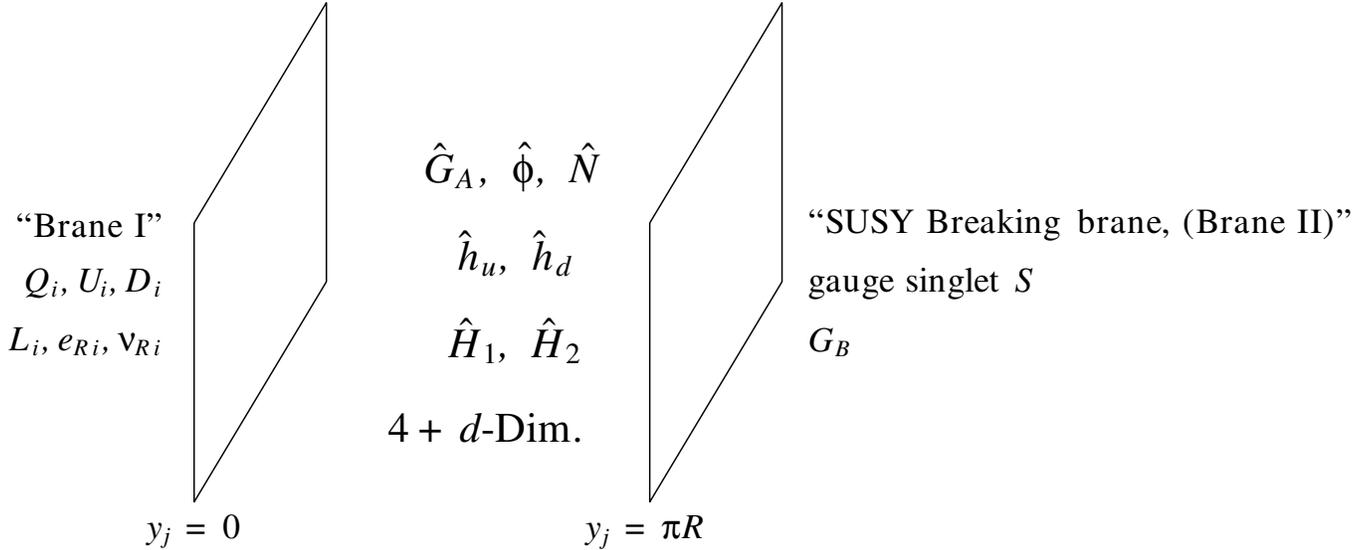
FIG. 1.   The model showing the parallel 3-branes spatially separated along d extra dimensions with coordinates $\mathbf{y} = (y_1, \ldots, y_d)$ and a common radius R. The index i in the matter fields represents the family index, $i = 1, 2, 3$.

turns out that the higher dimensional coupling is nonperturbative 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 nonperturbative 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

$$\hat{\kappa} \sim \hat{\lambda}_N \sim \hat{\lambda}_H \sim \hat{\lambda}_\nu, \tag{4}$$

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

$$\kappa \sim \lambda_N \sim \lambda_H \sim \lambda \equiv \left(\frac{M_*}{M_P}\right)^2, \qquad \lambda_\nu \sim \left(\frac{M_*}{M_P}\right). \tag{5}$$

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 ϕ is the inflaton which slow rolls in a semiflat 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 μ problem*: The term $\lambda_N N h_u h_d$ 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) [10,11].

3. *Right-handed neutrino masses*: We assume that the lightest right-handed neutrino gets Majorana mass through the nonrenormalizable operator in Eq. (1)

$\lambda_{\nu,1} N^2 (\nu_{R,1}^c)^2/M_*$, where the Yukawa coupling is suppressed due to the fact that the operator contains two bulk fields N, and is given by $\lambda_{\nu,1} = \lambda_\nu \sim \mathcal{O}(M_*/M_P)$. This lightest right-handed neutrino $\nu_{R,1}^c$, that we shall henceforth simply refer to as ν_R^c, will play the rôle of the coupled curvaton, although it may be completely subdominant in the seesaw 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), and therefore 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 $M_{\mathrm{SUSY}}/(4\pi)$ due to bulk particles propagating inside a loop with M_{SUSY} masses, i.e. gaugino mediation [12].

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

$$Q_N + Q_{h_u} + Q_{h_d} = 0, \qquad Q_\phi + 2Q_N = 0,$$
$$Q_\phi + Q_{H_1} + Q_{H_2} = 0, \qquad Q_N + Q_{\nu_{R,1}} = 0. \tag{6}$$

The global symmetry, $U(1)_{\mathrm{PQ}}$ also forbids explicit RH

Majorana neutrino masses in the superpotential, but $B - L$ symmetry is broken by the nonrenormalizable term $\lambda_{\nu,1} N^2 (\nu_{R,1}^c)^2 / M_*$. 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.

III. 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. (B2) and for the soft bulk masses Eq. (B3) given in Appendix B equal to one for simplicity, with

$$V_0 \sim F_S^2, \tag{7}$$

$$m^2 \equiv \left(\frac{F_S}{M_P}\right)^2 \sim m_\phi^2 \sim m_h^2 \sim m_N^2 \sim m_H^2, \tag{8}$$

$$A \equiv \frac{F_S}{M_*} = \frac{A_{\lambda_H}}{c_H} = \frac{A_{\lambda_N}}{c_N} = \frac{A_\kappa}{c_\kappa}. \tag{9}$$

With this simplifications the scalar potential for the real components of the fields, at energy below M_* looks like

$$
\begin{aligned}
V = {}& V_0 + m^2(\phi^2 + h^2 + N^2 + H^2) \\
& + 2\lambda A(c_H \phi H^2 - c_\kappa \phi N^2 + c_N N h^2) \\
& + \lambda^2(h^2 - 2\phi N)^2 + 2\lambda^2 N^2 h^2 + \lambda^2 (H^2 - N^2)^2 \\
& + 2\lambda^2 \phi^2 H^2 + 4\lambda_\nu^2 \frac{N^4 \tilde{\nu}_R^2}{M_*^2} + 4\lambda_\nu^2 \frac{N^2 \tilde{\nu}_R^4}{M_*^2} \\
& + 4\lambda_\nu \lambda \frac{N \tilde{\nu}_R^2}{M_*}(h^2 - 2\phi N),
\end{aligned}
\tag{10}
$$

where the couplings λ and λ_ν 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

$$\langle \phi \rangle = c_\kappa \frac{A}{4\lambda}, \tag{11}$$

$$\langle N \rangle = c_\kappa \frac{A}{2\sqrt{2}\lambda}, \tag{12}$$

$$\langle H \rangle = \langle h \rangle = \langle \tilde{\nu}_R \rangle = 0, \tag{13}$$

where all $c_i \approx \mathcal{O}(1)$. These coefficients should satisfy $c_H > c_\kappa/4$ ($c_N > (2 - \sqrt{2})c_\kappa/8$), 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 solu-

tions of the the minimization equations with $H \neq 0$ are a maximum of the potential instead.

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

$$\langle \phi \rangle \sim \langle N \rangle \sim \frac{A}{\lambda} = \frac{F_S}{M_*}\left(\frac{M_P}{M_*}\right)^2. \tag{14}$$

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_*,

$$F_S = \frac{M_*^4}{M_P^2}. \tag{15}$$

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

$$m^2 = M_*^2\left(\frac{M_*}{M_P}\right)^6, \qquad A = M_*\left(\frac{M_*}{M_P}\right)^2. \tag{16}$$

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_*

$$\langle \phi \rangle \sim \langle N \rangle \sim M_*. \tag{17}$$

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 Yuwaka 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.

IV. 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.

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[4]The reduced four dimensional Planck scale is fixed at $M_P = 2.4 \times 10^{18}$ GeV by gravity.

A. 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 Eq. (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, $m_{3/2} \sim m \sim M_{\text{SUSY}} \sim$ 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 to 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

$$M_* \approx 10^{13} \text{ GeV,} \qquad (18)$$

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. (B5) 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 [13,14]. Below we carefully examine the issues relevant to the present model.

B. The μ-scale

A problem of the minimal version of the supersymmetric standard model (MSSM) is why the μ 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 other words, what is the origin of the μ-scale? There are many solutions to this problem,[5] for example, in the NMSSM the μ-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 μ-term is generated.

The solution of the μ-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 κN^3 in the superpotential so that the model has an exact Z_3 symmetry [10,11] which is broken at the weak scale (at the scale of $\langle N \rangle$) leading to a serious domain wall problem[6] [17]. In our model a global $U(1)_{\text{PQ}}$ symmetry forbids such cubic terms so there is no domain wall problem. As we will see in the next subsection the global $U(1)_{\text{PQ}}$ 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 μ-term once it gets a VEV

$$\mu = \lambda \langle N \rangle \approx A \approx M_* \left(\frac{M_*}{M_P}\right)^2 \approx \text{TeV,} \qquad (19)$$

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

C. The axion scale

The most elegant explanation of the strong CP problem is provided by the Peccei-Quinn (PQ) mechanism [18], in which the CP violating angle $\bar{\Theta}$ is set to zero dynamically as a result of a global, spontaneously broken $U(1)_{\text{PQ}}$ 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 $U(1)_{\text{PQ}}$ Peccei-Quinn symmetry is spontaneously broken once the scalar fields charged under $U(1)_{\text{PQ}}$ [see Eq. (6)] get a VEV of the order M_*. This implies automatically that the axion decay constant is

$$f_a \sim M_* \sim 10^{13} \text{ GeV.} \qquad (20)$$

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 on f_a of order $f_a \lesssim 10^{12}$ GeV. It seems that our prediction is a little bit higher that the allowed by experiments. However, as has been pointed out in [19], f_a can be as big as 10^{15} 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

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[5]Note that the Giudice-Masiero mechanism [15] presents also a solution of the μ problem within the MSSM by generating the μ term via a nonminimal Kahler potential.

[6]See for example [16] for an alternative solution to the domain wall problem based on a Z_2 R-symmetry.

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. [4,20] 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} GeV may be consistent with the cosmological constraints.

D. 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 $\lambda_{\nu,i} N (\nu_{R,i}^c)^2$ 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 [21]. 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 seesaw 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 seesaw 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 [21,22].

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,

$$\langle N\rangle \sim M_* \sim 10^{13} \text{ GeV}. \quad (21)$$

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 μ parameter; the axion scale; and the atmospheric neutrino mass scale.

V. 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 [23] 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 rôle of the coupled curvaton is instead played by the lightest right-handed sneutrino, with the inflaton identified as the field ϕ, 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 ν_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 ϕ, while N field is held at zero for values of the inflaton field ϕ larger than the critical value

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

The inflationary epoch is therefore described by a slowly rolling inflation field ϕ 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

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

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

$$\epsilon_N = \frac{M_P^2 m_\phi^4 \phi_N^2}{V_0^2} < 1, \quad (24)$$

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

where the subscript N means N e-folds before the end of inflation. Using Eqs. (B2) and (B3) 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]:

$$P_{\mathcal{R}}^{1/2} \simeq \left(\frac{H_*}{\dot{\phi}_*}\right)\left(\frac{H_*}{2\pi}\right) \simeq \left(\frac{H_*}{2\pi\eta_N\phi_*}\right), \qquad (26)$$

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 [25], while the Hubble parameter is of the order of $H \simeq M_*(M_*/M_P)^3$. Therefore

$$P_{\mathcal{R}}^{1/2} \simeq \left(\frac{M_*}{M_P}\right)^3, \qquad (27)$$

which for $M_* \simeq 10^{13}$ GeV is quite below the COBE value $P_{\mathcal{R}}^{1/2} = 5 \times 10^{-5}$ [26].

However, the quantum fluctuations of any light field during inflation, i.e., the Higgs h and the lightest right-handed sneutrino ν_R, will contribute to the total curvature perturbation[7], $\mathcal{R}$, such that [27,28]

$$\mathcal{R} = \sum_\alpha \frac{\rho_\alpha + P_\alpha}{\rho + P} \mathcal{R}_\alpha, \qquad (28)$$

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

$$\mathcal{R}_\alpha \simeq H \frac{Q_\alpha}{\dot{\phi}_\alpha}, \qquad (29)$$

and Q_α the gauge invariant quantum fluctuations of the field [29]. Given the model parameters, we know that the inflaton field has a background value of the order of the cutoff 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 multicomponent 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 adiabatic one, i.e., the total curvature perturbation, when their contribution to the total energy density becomes comparable [30–35]. 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 this possibility, i.e., the lightest right-handed sneutrino as the main source of the isocurvature perturbation during inflation. The Yukawa coupling of the lightest right-handed sneutrino to the N field is a factor (M_P/M_*) larger than the coupling κ between ϕ 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 ϕ but with a smaller amplitude $\nu_R \sim (\kappa/\lambda_\nu)M_*$. 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 1 degree of freedom, the adiabatic mode [28].

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 $\mathcal{R}_\nu$ perturbation is dragged toward $\mathcal{R}$, and by the time they first reach the global minimum all the contributions $\mathcal{R}_\alpha$ are comparable but some orders of magnitude smaller than the initial $\mathcal{R}_\nu$.

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[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 $k \ll Ha$.
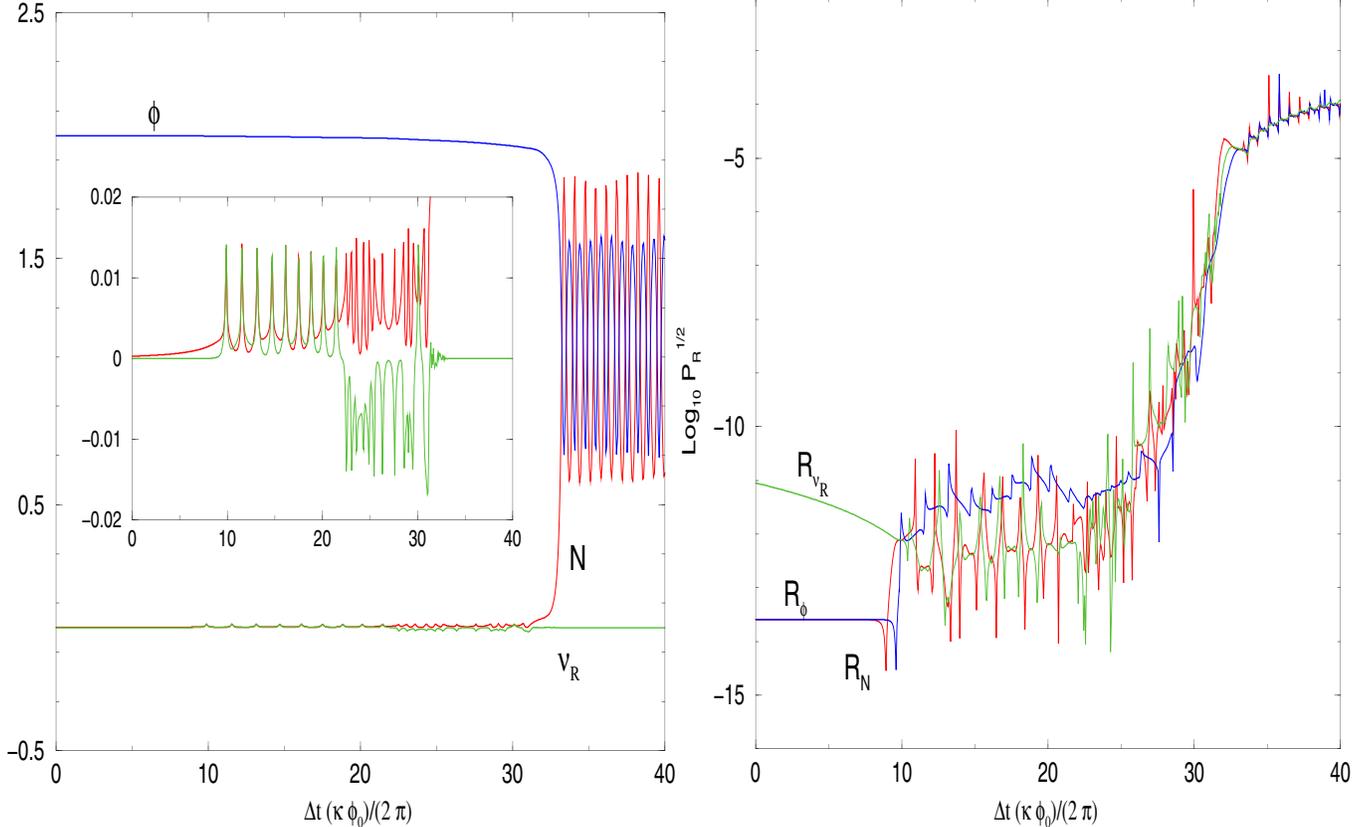
FIG. 2 (color online).   Left plot: Evolution of the background fields, ϕ, N, and $\tilde{\nu}_R$, after they have passed the critical point; inset: oscillations of N and $\tilde{\nu}_R$, previous to that of ϕ. Right plot: Amplitude of the spectrum of the curvature $P_{\mathcal{R}_\alpha}^{1/2}$ for the different components of the model: inflaton ϕ, N field and lightest right-handed sneutrino sneutrino $\tilde{\nu}_R$.

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-ϕ 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 ϕ. First N gets destabilized, and then $\tilde{\nu}_R$. But for the values of the fields such that $\lambda_\nu \tilde{\nu}_R^2/M_* > \kappa(\phi_c - \phi)$, there is an approximate minimum with $N \simeq \tilde{\nu}_R$ and $\lambda_\nu \tilde{\nu}_R^2/M_* \simeq \phi \simeq \phi_c$, 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 [36]. 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 ϕ 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 (left-hand side plot), and the different field contributions to the curvature perturbation $\mathcal{R}_\alpha$ (right-hand side 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 $\mathcal{R}_{\nu*} \simeq 10^{-5}$ at horizon-crossing. Nevertheless, as it can be seen in the plot, after the phase transition all contributions are roughly equal. For

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[8]Without the sneutrino and/or the Higgs, there is indeed a strong parametric resonance (tachyonic preheating [36]) for the fluctuations with a wavenumber up to $O(\kappa N_0)$ [37] during the first 3 oscillations of the fields, before it enters in a narrow resonance regime.

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 ϕ and N oscillating.

The decay of the lightest right-handed sneutrino gives rise to radiation, with $\rho_R \simeq 0.2 \rho_\phi$, where ρ_ϕ refers to the energy density in the oscillating singlets.[9] 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.

VI. CONCLUSION

We have proposed a model of cosmology and particle physics in which all relevant scales are derived from an intermediate string scale $M_* \sim 10^{13}$ 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 $F_S^{1/2} \sim M_*^2/M_P \sim 10^8$ GeV. The μ parameter of the MSSM and the electroweak breaking scale are then given by $F_S/M_* \sim 10^3$ GeV. The model involves hybrid inflation, with the inflaton mass given by $F_S/M_P \sim$ MeV, and their Yukawa couplings given by $(M_*/M_P)^2 \sim 10^{-10}$. In our model we were 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 μ parameter; the axion scale; and the neutrino mass scale. In recent years models with an intermediate string scale ($10^{11} < M_* < 10^{14}$ GeV) have been of great interest because contain many phenomenological issues (see for example [13,14]). The novelty of our construction is we give an explicit potential (10) where the link between particle physics and cosmology is explicit. Also 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 [21] since in this case the lightest right-handed neutrino plays no essential rôle in the seesaw mechanism. The sneutrino as a coupled curvaton only strengthens the synthesis between cosmology and particle physics at the intermediate scale.

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

$$\mathcal{L}_{4+d} = \int d^2\theta \hat{W}_{4+d}(\delta^d(0) + \delta^d(y_i - \pi R)), \quad \text{(A1)}$$

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

$$W_{4+d} \ni \frac{\hat{\lambda}_{i,j}}{M_*^{\frac{\alpha d}{2}}} \frac{\hat{\Psi}_i^\alpha \Phi_j^\beta}{M_*^{\alpha+\beta-3}}. \quad \text{(A2)}$$

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

$$M_P^2 = M_*^{2+d} R^d. \quad \text{(A3)}$$

Using the fact that the bulk fields contain a volume suppression factor with respect 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

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[9]The 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 [38].

[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.

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

$$= \hat{\lambda}_{i,j} \left(\frac{M_*}{M_P} \right)^\alpha \frac{\Psi_i^\alpha \Phi_j^\beta}{M_*^{\alpha+\beta-3}}. \tag{A4}$$

Redefining the effective four dimensional coupling as

$$\lambda_{i,j} = \left(\frac{M_*}{M_P} \right)^\alpha \hat{\lambda}_{i,j}, \tag{A5}$$

we notice that in general we get an important suppression factor if $M_* \ll M_P$[11]. In other words, if there is a fundamental intermediate scale, M_*, 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 nonperturbative 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 [12]. The SUSY-breaking Lagrangian contains six terms that lead respectively to the bulk gaugino mass (M_A), Brane II gaugino mass (M_B), vacuum energy (V_0), soft masses for bulk fields (M_{Ψ_i}), soft masses for Brane II fields (M_{Φ_i}) and trilinear soft terms ($A_{\lambda_{i,j}}$). In general we have

$$\mathcal{L}_{4+d}^{\text{soft}} = \delta^d(y_i - \pi R)(c_{\lambda_{ij}} \int d^2\theta \frac{\hat{W}_{4+d}S}{M_*}$$

$$+ c_{\Psi_i} \int d^4\theta \frac{\hat{\Psi}_i^\dagger \hat{\Psi}_i S^\dagger S}{M_*^{2+d}} + c_{\Phi_i} \int d^4\theta \frac{\Phi_i^\dagger \Phi_i S^\dagger S}{M_*^2}$$

$$+ c_V \int d^4\theta S^\dagger S + c_A \int d^2\theta \frac{\hat{W}_\alpha^{(A)} \hat{W}^{\alpha(A)} S}{M_*^{1+d}}$$

$$+ c_B \int d^2\theta \frac{W_\alpha^{(B)} W^{\alpha(B)} S}{M_*}), \tag{B1}$$

where $W_\alpha^{(A)}$ ($W_\alpha^{(B)}$) 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 nonrenomalizable terms in Eq. (B1) 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

$$V_0 = c_V F_S^2 \tag{B2}$$

$$M_A = c_A \frac{F_S}{M_P} \qquad M_{\Psi_i}^2 = c_{\Psi_i} \left(\frac{F_S}{M_P} \right)^2 \tag{B3}$$

$$M_B = c_B \frac{F_S}{M_*} \qquad M_{\Phi_i}^2 = c_{\Phi_i} (\frac{F_S}{M_*})^2 \qquad A_{\lambda_{i,j}} = c_{\lambda_{ij}} \frac{F_S}{M_*}. \tag{B4}$$

The vacuum energy (V_0) would dominate the total energy density during inflation providing the typical expansion rate (the Hubble constant) as[12]

$$H^2 = \frac{V_0}{3M_P^2} = \frac{c_V}{3} \left(\frac{F_S}{M_P} \right)^2 = \frac{c_V}{3c_{\Psi_i}} M_{\Psi_i}^2. \tag{B5}$$

The bulk mass scale (M_{Ψ_i}) would give us, for example, the inflaton mass. Therefore, in order to satisfy the slow-roll condition during inflation ($M_{\Psi_i} < H$) it turns out from (B5) that we just need some tuning on the constants c_i, i.e. $3c_{\Psi_i} < c_V$. Finally, the Brane II mass scale would define the typical MSSM soft term, $M_{\text{SUSY}} \sim M_{\Phi_i} \sim \mathcal{O}$ (TeV). Notice that in general we would have $H \ll M_{\text{SUSY}}$ which is completely different from what happen in the normal four dimensional supersymmetric hybrid inflationary models where $H \sim M_{\text{SUSY}}$ [25].

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[11]This suppression factor only would depend on the number α of fields living in the bulk and it will be completely independent of the number d of extra dimensions.

[12]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 is inflated away. In the case that all de $4 + d$-dimensions feel inflation at the same time, the Hubble constant would be given by $H^2 = V_0/(3M_*^2)$ instead [40].

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