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Moduli stabilization, de Sitter vacua and supersymmetry breaking

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We describe the phenomenology of a model of supersymmetry breaking in the presence of a tiny (tunable) positive cosmological constant. It utilises a single chiral multiplet with a gauged shift symmetry, that can be identified with the string dilaton (or an appropriate compactification modulus). The model is coupled to the MSSM, leading to calculable soft supersymmetry breaking masses and a distinct low energy phenomenology that allows to differentiate it from other models of supersymmetry breaking and mediation mechanisms. We also study the question if this model can lead to inflation by identifying the dilaton with the inflaton. We find that this is possible if the Kähler potential is modified by a term that has the form of NS5-brane instantons, leading to an appropriate inflationary plateau around the maximum of the scalar potential, depending on two extra parameters.

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1. Introduction

If String Theory is a fundamental theory of Nature and not just a tool for studying systems with strongly coupled dynamics, it should be able to describe at the same time particle physics and cosmology, which are phenomena that involve very different scales from the microscopic four-dimensional (4d) quantum gravity length of 10^{-33} cm to large macroscopic distances of the size of the observable Universe $\sim 10^{28}$ cm spanned a region of about 60 orders of magnitude. In particular, besides the 4d Planck mass, there are three very different scales with very different physics corresponding to the electroweak, dark energy and inflation. These scales might be related via the scale of the underlying fundamental theory, such as string theory, or they might be independent in the sense that their origin could be based on different and independent dynamics. An example of the former constraint and more predictive possibility is provided by TeV strings with a fundamental scale at low energies due for instance to large extra dimensions transverse to a four-dimensional braneworld forming our Universe [1]. In this case, the 4d Planck mass is emergent from the fundamental string scale and inflation should also happen around the same scale [2].

Here, we will adopt the second more conservative approach, assuming that all three scales have an independent dynamical origin. Moreover, we will assume the presence of low energy supersymmetry that allows for an elegant solution of the mass hierarchy problem, a unification of fundamental forces as indicated by low energy data and a natural dark matter candidate due to an unbroken R-parity. The assumption of independent scales implies that supersymmetry breaking should be realized in a metastable de Sitter vacuum with an infinitesimally small (tunable) cosmological constant independent of the supersymmetry breaking scale that should be in the TeV region. In a recent work [3], we studied a simple $N = 1$ supergravity model having this property and motivated by string theory. Besides the gravity multiplet, the minimal field content consists of a chiral multiplet with a shift symmetry promoted to a gauged R-symmetry using a vector multiplet. In the string theory context, the chiral multiplet can be identified with the string dilaton (or an appropriate compactification modulus) and the shift symmetry associated to the gauge invariance of a two-index antisymmetric tensor that can be dualized to a (pseudo)scalar. The shift symmetry fixes the form of the superpotential and the gauging allows for the presence of a Fayet-Iliopoulos (FI) term, leading to a supergravity action with two independent parameters that can be tuned so that the scalar potential possesses a metastable de Sitter minimum with a tiny vacuum energy (essentially the relative strength between the F- and D-term contributions). A third parameter fixes the Vacuum Expectation Value (VEV) of the string dilaton at the desired (phenomenologically) weak coupling regime. An important consistency constraint of our model is anomaly cancellation which has been studied in [5] and implies the existence of additional charged fields under the gauged R-symmetry.

In a more recent work [6], we analyzed a small variation of this model which is manifestly anomaly free without additional charged fields and allows to couple in a straight forward way a visible sector containing the minimal supersymmetric extension of the Standard Model (MSSM) and studied the mediation of supersymmetry breaking and its phenomenological consequences. It turns out that an additional 'hidden sector' field z is needed to be added for the matter soft scalar masses to be non-tachyonic; although this field participates in the supersymmetry breaking and is similar to the so-called Polonyi field, it does not modify the main properties of the metastable de Sitter (dS) vacuum. All soft scalar masses, as well as trilinear A-terms, are generated at the tree level and are universal under the assumption that matter kinetic terms are independent of the 'Polonyi' field, since matter fields are neutral under the shift symmetry and supersymmetry breaking is driven by a combination of the $U(1)$ D-term and the dilaton and z-field F-term. Alternatively, a way to avoid the tachyonic scalar masses without adding the extra field z is to modify the matter kinetic terms by a dilaton dependent factor.

A main difference of the second analysis from the first work is that we use a field representation in which the gauged shift symmetry corresponds to an ordinary $U(1)$ and not an R-symmetry. The two representations differ by a Kähler transformation that leaves the classical supergravity action invariant. However, at the quantum level, there is a Green-Schwarz term generated that amounts an extra dilaton dependent contribution to the gauge kinetic terms needed to cancel the anomalies of the R-symmetry. This creates an apparent puzzle with the gaugino masses that vanish in the first representation but not in the latter. The resolution to the puzzle is based to the so called anomaly mediation contributions [7, 8] that explain precisely the above apparent discrepancy. It turns out that gaugino masses are generated at the quantum level and are thus suppressed compared to the scalar masses (and A-terms).

This model has the necessary ingredients to be obtained as a remnant of moduli stabilisation within the framework of internal magnetic fluxes in type I string theory, turned on along the compact directions for several abelian factors of the gauge group. All geometric moduli can in principle be fixed in a supersymmetric way, while the shift symmetry is associated to the 4d axion and its gauging is a consequence of anomaly cancellation [9, 10].

We then make an attempt to connect the scale of inflation with the electroweak and supersymmetry breaking scales within the same effective field theory, that at the same time allows the existence of an infinitesimally small (tuneable) positive cosmological constant describing the present dark energy of the universe. We thus address the question whether the same scalar potential can provide inflation with the dilaton playing also the role of the inflaton at an earlier stage of the universe evolution [11]. We show that this is possible if one modifies the Kähler potential by a correction that plays no role around the minimum, but creates an appropriate plateau around the maximum. In general, the Kähler potential receives perturbative and non-perturbative corrections that vanish in the weak coupling limit. After analysing all such corrections, we find that only those that have the form of (Neveu-Schwarz) NS5-brane instantons can lead to an inflationary period compatible with cosmological observations. The scale of inflation turns out then to be of the order of low energy supersymmetry breaking, in the TeV region. On the other hand, the predicted tensor-to-scalar ratio is too small to be observed.

2. Conventions

Throughout this paper we use the conventions of [12]. A supergravity theory is specified (up to Chern-Simons terms) by a Kähler potential $\mathscr{K}$, a superpotential W, and the gauge kinetic functions $f_{AB}(z)$. The chiral multiplets z^α, χ^α are enumerated by the index α and the indices A, B indicate the different gauge groups. Classically, a supergravity theory is invariant under Kähler tranformations, viz.

$$\mathscr{K}(z,\bar{z}) \longrightarrow \mathscr{K}(z,\bar{z}) + J(z) + \bar{J}(\bar{z}),$$

$$W(z) \longrightarrow e^{-\kappa^2 J(z)} W(z), \tag{2.1}$$

where κ is the inverse of the reduced Planck mass, $m_p = \kappa^{-1} = 2.4 \times 10^{15}$ TeV. The gauge transformations of chiral multiplet scalars are given by holomorphic Killing vectors, i.e. $\delta z^\alpha = \theta^A k_A^\alpha(z)$, where θ^A is the gauge parameter of the gauge group A. The Kähler potential and superpotential need not be invariant under this gauge transformation, but can change by a Kähler transformation

$$\delta \mathscr{K} = \theta^A \left[r_A(z) + \bar{r}_A(\bar{z}) \right], \tag{2.2}$$

provided that the gauge transformation of the superpotential satisfies $\delta W = -\theta^A \kappa^2 r_A(z) W$. One then has from $\delta W = W_\alpha \delta z^\alpha$

$$W_\alpha k_A^\alpha = -\kappa^2 r_A W, \tag{2.3}$$

where $W_\alpha = \partial_\alpha W$ and α labels the chiral multiplets. The supergravity theory can then be described by a gauge invariant function

$$\mathscr{G} = \kappa^2 \mathscr{K} + \log(\kappa^6 W \bar{W}). \tag{2.4}$$

The scalar potential is given by

$$\begin{aligned} V &= V_F + V_D \\ V_F &= e^{\kappa^2 \mathscr{K}} \left(-3\kappa^2 W \bar{W} + \nabla_\alpha W g^{\alpha\bar{\beta}} \bar{\nabla}_{\bar{\beta}} \bar{W} \right) \\ V_D &= \frac{1}{2} (\mathrm{Re} f)^{-1\, AB} \mathscr{P}_A \mathscr{P}_B, \end{aligned} \tag{2.5}$$

where W appears with its Kähler covariant derivative

$$\nabla_\alpha W = \partial_\alpha W(z) + \kappa^2 (\partial_\alpha \mathscr{K}) W(z). \tag{2.6}$$

The moment maps $\mathscr{P}_A$ are given by

$$\mathscr{P}_A = i(k_A^\alpha \partial_\alpha \mathscr{K} - r_A). \tag{2.7}$$

In this paper we will be concerned with theories having a gauged R-symmetry, for which $r_A(z)$ is given by an imaginary constant $r_A(z) = i\kappa^{-2}\xi$. In this case, $\kappa^{-2}\xi$ is a Fayet-Iliopoulos [13] constant parameter.

3. The model

The starting point is a chiral multiplet S and a vector multiplet associated with a shift symmetry of the scalar component s of the chiral multiplet S

$$\delta s = -ic\theta\,, \tag{3.1}$$

and a string-inspired Kähler potential of the form $-p\log(s+\bar{s})$. The most general superpotential is either a constant $W = \kappa^{-3}a$ or an exponential superpotential $W = \kappa^{-3}ae^{bs}$ (where a and b are constants). A constant superpotential is (obviously) invariant under the shift symmetry, while an exponential superpotential transforms as $W \to We^{-ibc\theta}$, as in eq. (2.3). In this case the shift symmetry becomes a gauged R-symmetry and the scalar potential contains a Fayet-Iliopoulos term.

Note however that by performing a Kähler transformation (2.1) with $J = \kappa^{-2} bs$, the model can be recast into a constant superpotential at the cost of introducing a linear term in the Kähler potential $\delta K = b(s+\bar{s})$. Even though in this representation, the shift symmetry is not an R-symmetry, we will still refer to it as $U(1)_R$. The most general gauge kinetic function has a constant term and a term linear in s, $f(s) = \delta + \beta s$.

To summarise,[1]

$$\begin{aligned} \mathscr{K}(s,\bar{s}) &= -p\log(s+\bar{s}) + b(s+\bar{s}), \\ W(s) &= a, \\ f(s) &= \delta + \beta s, \end{aligned} \tag{3.2}$$

where we have set the mass units $\kappa = 1$. The constants a and b together with the constant c in eq. (3.1) can be tuned to allow for an infinitesimally small cosmological constant and a TeV gravitino mass. For $b > 0$, there always exists a supersymmetric AdS (anti-de Sitter) vacuum at $\langle s+\bar{s}\rangle = b/p$, while for $b = 0$ (and $p < 3$) there is an AdS vacuum with broken supersymmetry. We therefore focus on $b < 0$. In the context of string theory, S can be identified with a compactification modulus or the universal dilaton and (for negative b) the exponential superpotential may be generated by non-perturbative effects.

The scalar potential is given by:

$$\begin{aligned} V &= V_F + V_D \\ V_F &= a^2 e^{\frac{b}{l}} l^{p-2} \left\{ \frac{1}{p}(pl-b)^2 - 3l^2 \right\} \qquad l = 1/(s+\bar{s}) \\ V_D &= c^2 \frac{l}{\beta + 2\delta l}(pl-b)^2 \end{aligned} \tag{3.3}$$

In the case where S is the string dilaton, V_D can be identified as the contribution of a magnetized D-brane, while V_F for $b = 0$ and $p = 2$ coincides with the tree-level dilaton potential obtained by considering string theory away its critical dimension [14]. For $p \geq 3$ the scalar potential V is positive and monotonically decreasing, while for $p < 3$, its F-term part V_F is unbounded from below when $s+\bar{s} \to 0$. On the other hand, the D-term part of the scalar potential V_D is positive and diverges when $s+\bar{s} \to 0$ and for various values for the parameters an (infinitesimally small) positive (local) minimum of the potential can be found.

If we restrict ourselves to integer p, tunability of the vacuum energy restricts $p = 2$ or $p = 1$ when $f(s) = s$, or $p = 1$ when the gauge kinetic function is constant. For $p = 2$ and $f(s) = s$, the minimization of V yields:

$$b/l = -\rho_0 \approx -0.183268 \quad , \quad p = 2 \tag{3.4}$$

$$\frac{a^2}{bc^2} = A_2(-\rho_0) + B_2(-\rho_0)\frac{\Lambda}{b^3c^2} \approx -50.6602 + \mathscr{O}(\Lambda), \tag{3.5}$$

where Λ is the value of V at the minimum (i.e. the cosmological constant), $-\rho_0$ is the negative root of the polynomial $-x^5 + 7x^4 - 10x^3 - 22x^2 + 40x + 8$ compatible with (3.5) for $\Lambda = 0$ and $A_2(\alpha)$,

[1] In superfields the shift symmetry (3.1) is given by $\delta S = -ic\Lambda$, where Λ is the superfield generalization of the gauge parameter. The gauge invariant Kähler potential is then given by $\mathscr{K}(S,\bar{S}) = -p\kappa^{-2}\log(S+\bar{S}+cV_R) + \kappa^{-2}b(S+\bar{S}+cV_R)$, where V_R is the gauge superfield of the shift symmetry.

$B_2(\alpha)$ are given by

$$A_2(\alpha) = 2e^{-\alpha}\frac{-4+4\alpha-\alpha^2}{\alpha^3-4\alpha^2-2\alpha} \quad ; \quad B_2(\alpha) = 2\frac{\alpha^2 e^{-\alpha}}{\alpha^2-4\alpha-2} \tag{3.6}$$

It follows that by carefully tuning a and c, Λ can be made positive and arbitrarily small independently of the supersymmetry breaking scale. A plot of the scalar potential for certain values of the parameters is shown in figure 1.

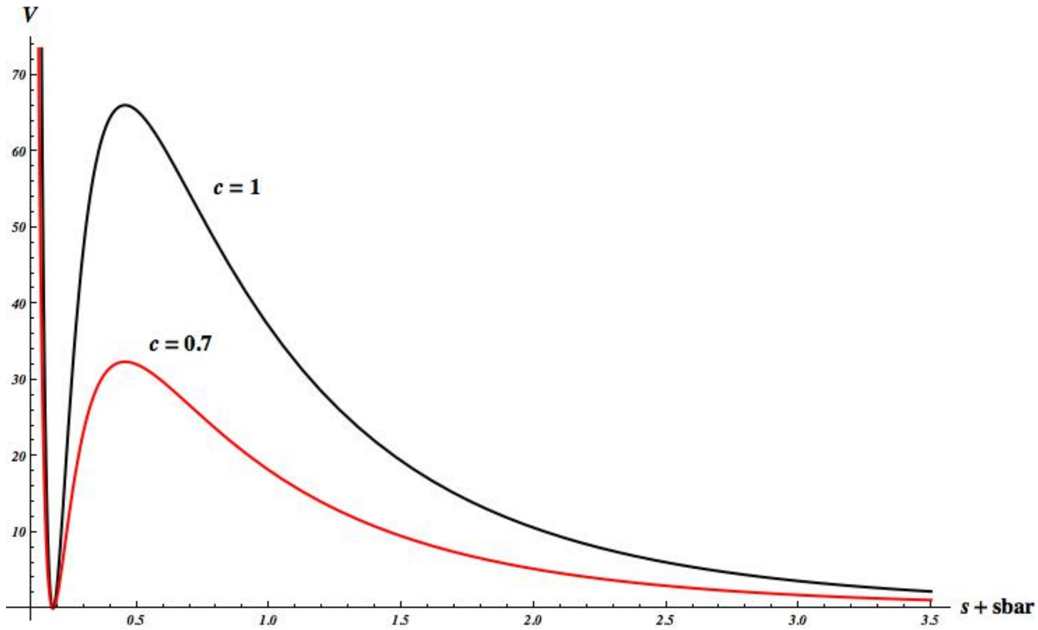

Figure 1: A plot of the scalar potential for $p=2$, $b=-1$, $\delta=0$, $\beta=1$ and a given by equation (3.5) for $c=1$ (black curve) and $c=0.7$ (red curve).

At the minimum of the scalar potential, for nonzero a and $b<0$, supersymmetry is broken by expectation values of both an F and D-term. Indeed the F-term and D-term contributions to the scalar potential are

$$\begin{aligned} V_F|_{s+\bar{s}=\frac{-\rho_0}{b}} &= \frac{1}{2}a^2b^2e^-\rho_0\left(1+\frac{2}{\rho_0}\right)^2 > 0, \\ V_D|_{s+\bar{s}=\frac{-\rho_0}{b}} &= -\frac{b^3c^2}{\rho_0}\left(1+\frac{2}{\rho_0}\right)^2 > 0\,. \end{aligned} \tag{3.7}$$

The gravitino mass term is given by

$$(m_{3/2})^2 = e^{\mathscr{G}} = \frac{a^2b^2}{\rho_0^2}e^{-\rho_0}\,. \tag{3.8}$$

Due to the Stueckelberg coupling, the imaginary part of s (the axion) gets eaten by the gauge field, which acquires a mass. On the other hand, the Goldstino, which is a linear combination of the fermion of the chiral multiplet χ and the gaugino λ gets eaten by the gravitino. As a result, the physical spectrum of the theory consists (besides the graviton) of a massive scalar, namely the dilaton, a Majorana fermion, a massive gauge field and a massive gravitino. All the masses are of the same order of magnitude as the gravitino mass, proportional to the same constant a (or c related by eq. (3.5) where b is fixed by eq. (3.4)), which is a free parameter of the model. Thus, they vanish in the same way in the supersymmetric limit $a \to 0$.

The local dS minimum is metastable since it can tunnel to the supersymmetric ground state at infinity in the s-field space (zero coupling). It turns out however that it is extremely long lived for realistic perturbative values of the gauge coupling $l \simeq 0.02$ and TeV gravitino mass and, thus, practically stable; its decay rate is [5]:

$$\Gamma \sim e^{-B} \quad \text{with} \quad B \approx 10^{300}\,. \tag{3.9}$$

4. Coupling a visible sector

The guideline to construct a realistic model keeping the properties of the toy model described above is to assume that matter fields are invariant under the shift symmetry (3.1) and do not participate in the supersymmetry breaking. In the simplest case of a canonical Kähler potential, MSSM-like fields ϕ can then be added as:

$$\begin{aligned} \mathscr{K} &= -\kappa^{-2}\log(s+\bar{s}) + \kappa^{-2} b(s+\bar{s}) + \sum \varphi\bar{\varphi}, \\ W &= \kappa^{-3} a + W_{MSSM}, \end{aligned} \tag{4.1}$$

where $W_{MSSM}(\phi)$ is the usual MSSM superpotential. The squared soft scalar masses of such a model can be shown to be positive and close to the square of the gravitino mass (TeV^2). On the other hand, for a gauge kinetic function with a linear term in s, $\beta \neq 0$ in eq. (3.2), the Lagrangian is not invariant under the shift symmetry

$$\delta\mathscr{L} = -\theta\frac{\beta c}{8}\,\varepsilon^{\mu\nu\rho\sigma}F_{\mu\nu}F_{\rho\sigma}. \tag{4.2}$$

and its variation should be canceled. As explained in Ref. [5], in the 'frame' with an exponential superpotential the R-charges of the fermions in the model can give an anomalous contribution to the Lagrangian. In this case the 'Green-Schwarz' term $\mathrm{Im} sF\tilde{F}$ can cancel quantum anomalies. However as shown in [5], with the minimal MSSM spectrum, the presence of this term requires the existence of additional fields in the theory charged under the shift symmetry.

Instead, to avoid the discussion of anomalies, we focus on models with a constant gauge kinetic function. In this case the only (integer) possibility[2] is $p = 1$. The scalar potential is given by (3.3) with $\beta = 0$, $\delta = p = 1$. The minimization yields to equations similar to (3.4), (3.5) and (3.6) with a different value of ρ_0 and functions A_1 and B_1 given by:

$$\begin{aligned} b\langle s+\bar{s}\rangle &= -\rho_0 \approx -0.233153 \\ \frac{bc^2}{a^2} &= A_1(-\rho_0) + B_1(-\rho_0)\frac{\Lambda}{a^2 b} \approx -0.359291 + \mathscr{O}(\Lambda) \\ A_1(\alpha) &= 2e^{\alpha}\alpha\frac{3-(\alpha-1)^2}{(\alpha-1)^2} \quad , \quad B_1(\alpha) = \frac{2\alpha^2}{(\alpha-1)^2}\,, \end{aligned} \tag{4.3}$$

[2] If $f(s)$ is constant, the leading contribution to V_D when $s+\bar{s} \to 0$ is proportional to $1/(s+\bar{s})^2$, while the leading contribution to V_F is proportional to $1/(s+\bar{s})^p$. It follows that $p < 2$; if $p > 2$, the potential is unbounded from below, while if $p = 2$, the potential is either positive and monotonically decreasing or unbounded from below when $s+\bar{s} \to 0$ depending on the values of the parameters.

where $-\rho_0$ is the negative root of $-3+(\rho-1)^2(2-\rho^2/2)=0$ close to -0.23, compatible with the second constraint for $\Lambda=0$. However, this model suffers from tachyonic soft masses when it is coupled to the MSSM, as in (4.1). To circumvent this problem, one can add an extra hidden sector field which contributes to (F-term) supersymmetry breaking. Alternatively, the problem of tachyonic soft masses can also be solved if one allows for a non-canonical Kähler potential in the visible sector, which gives an additional contribution to the masses through the D-term.

Let us discuss first the addition of an extra hidden sector field z (similar to the so-called Polonyi field [15]). The Kähler potential, superpotential and gauge kinetic function are given by

$$\begin{aligned}
\mathcal{K} &= -\kappa^{-2}\log(s+\bar{s})+\kappa^{-2}b(s+\bar{s})+z\bar{z}+\sum\varphi\bar{\varphi}\,,\\
W &= \kappa^{-3}a(1+\gamma\kappa z)+W_{\mathrm{MSSM}}(\varphi)\,,\\
f(s) &= 1 \quad , \quad f_A=1/g_A^2\,,
\end{aligned}\tag{4.4}$$

where A labels the Standard Model gauge group factors and γ is an additional constant parameter. The existence of a tuneable dS vacuum with supersymmetry breaking and non-tachyonic scalar masses implies that γ must be in a narrow region:

$$0.5\lesssim\gamma\lesssim 1.7\,.\tag{4.5}$$

In the above range of γ the main properties of the toy model described in the previous section remain, while Rez and its F-auxiliary component acquire non vanishing VEVs. All MSSM soft scalar masses are then equal to a universal value m_0 of the order of the gravitino mass, while the B_0 Higgs mixing parameter is also of the same order:

$$\begin{aligned}
m_0^2 &= m_{3/2}^2\left[(\sigma_s+1)+\frac{(\gamma+t+\gamma t)^2}{(1+\gamma t)^2}\right],\\
A_0 &= m_{3/2}\left[(\sigma_s+3)+t\frac{(\gamma+t+\gamma t^2)}{1+\gamma t}\right],\\
B_0 &= m_{3/2}\left[(\sigma_s+2)+t\frac{(\gamma+t+\gamma t^2)}{(1+\gamma t)}\right],
\end{aligned}\tag{4.6}$$

where $\sigma_s=-3+(\rho+1)^2$ with $\rho=-b(s+\bar{s})$ and $t\equiv\langle\mathrm{Re}\,z\rangle$ determined by the minimization conditions as functions of γ. Also, A_0 is the soft trilinear scalar coupling in the standard notation, satisfying the relation [16]

$$A_0=B_0+m_{3/2}\,.\tag{4.7}$$

On the other hand, the gaugino masses appear to vanish at tree-level since the gauge kinetic functions are constants (see (4.4)). However, as mentioned in Section 3, this model is classically equivalent to the theory[3]

$$\begin{aligned}
\mathcal{K} &= -\kappa^{-2}\log(s+\bar{s})+z\bar{z}+\sum_{\varphi}\varphi\bar{\varphi},\\
W &= \left(\kappa^{-3}a(1+z)+W_{\mathrm{MSSM}}(\varphi)\right)e^{bs}\,,
\end{aligned}\tag{4.8}$$

[3] This statement is only true for supergravity theories with a non-vanishing superpotential where everything can be defined in terms of a gauge invariant function $G=\kappa^2\mathcal{K}+\log(\kappa^6W\bar{W})$ [17].

obtained by applying a Kähler transformation (2.1) with $J = -\kappa^{-2} bs$. All classical results remain the same, such as the expressions for the scalar potential and the soft scalar masses (4.6), but now the shift symmetry (3.1) of s became a gauged R-symmetry since the superpotential transforms as $W \longrightarrow We^{-ibc\theta}$. Therefore, all fermions (including the gauginos and the gravitino) transform[4] as well under this $U(1)_R$, leading to cubic $U(1)_R^3$ and mixed $U(1) \times G_{\mathrm{MSSM}}$ anomalies. These anomalies are cancelled by a Green-Schwarz (GS) counter term that arises from a quantum correction to the gauge kinetic functions:

$$f_A(s) = 1/g_A^2 + \beta_A s \quad \text{with} \quad \beta_A = \frac{b}{8\pi^2}\left(T_{R_A} - T_{G_A}\right), \tag{4.9}$$

where T_G is the Dynkin index of the adjoint representation, normalized to N for $SU(N)$, and T_R is the Dynkin index associated with the representation R of dimension d_R, equal to $1/2$ for the $SU(N)$ fundamental. An implicit sum over all matter representations is understood. It follows that gaugino masses are non-vanishing in this representation, creating a puzzle on the quantum equivalence of the two classically equivalent representations. The answer to this puzzle is based on the fact that gaugino masses are present in both representations and are generated at one-loop level by an effect called Anomaly Mediation [7, 8]. Indeed, it has been argued that gaugino masses receive a one-loop contribution due to the super-Weyl-Kähler and sigma-model anomalies, given by [8]:

$$M_{1/2} = -\frac{g^2}{16\pi^2}\left[(3T_G - T_R)m_{3/2} + (T_G - T_R)\mathcal{K}_\alpha F^\alpha + 2\frac{T_R}{d_R}(\log\det\mathcal{K}|_R{}'')_{,\alpha}F^\alpha\right]. \tag{4.10}$$

The expectation value of the auxiliary field F^α, evaluated in the Einstein frame is given by

$$F^\alpha = -e^{\kappa^2\mathcal{K}/2} g^{\alpha\bar\beta}\bar\nabla_{\bar\beta}\bar W. \tag{4.11}$$

Clearly, for the Kähler potential (4.4) or (4.8) the last term in eq. (4.10) vanishes. However, the second term survives due to the presence of Planck scale VEVs for the hidden sector fields s and z. Since the Kähler potential between the two representations differs by a linear term $b(s+\bar s)$, the contribution of the second term in eq. (4.10) differs by a factor

$$\delta m_A = \frac{g_A^2}{16\pi^2}(T_G - T_R) b e^{\kappa^2\mathcal{K}/2} g^{\alpha\bar\beta}\bar\nabla_{\bar\beta}\bar W, \tag{4.12}$$

which exactly coincides with the 'direct' contribution to the gaugino masses due to the field dependent gauge kinetic function (4.9) (taking into account a rescaling proportional to g_A^2 due to the non-canonical kinetic terms).

We conclude that even though the models (4.4) and (4.8) differ by a (classical) Kähler transformation, they generate the same gaugino masses at one-loop. While the one-loop gaugino masses for the model (4.4) are generated entirely by eq. (4.10), the gaugino masses for the model (4.8) after a Kähler transformation have a contribution from eq. (4.10) as well as from a field dependent gauge kinetic term whose presence is necessary to cancel the mixed $U(1)_R \times G$ anomalies due to the fact that the extra $U(1)$ has become an R-symmetry giving an R-charge to all fermions in the theory. Using (4.10), one finds:

$$M_{1/2} = -\frac{g^2}{16\pi^2} m_{3/2}\left[(3T_G - T_R) - (T_G - T_R)\left((\rho+1)^2 + t\frac{\gamma + t + \gamma t^2}{1+\gamma t}\right)\right]. \tag{4.13}$$

[4]The chiral fermions, the gauginos and the gravitino carry a charge $bc/2$, $-bc/2$ and $-bc/2$ respectively.

For $U(1)_Y$ we have $T_G = 0$ and $T_R = 11$, for $SU(2)$ we have $T_G = 2$ and $T_R = 7$, and for $SU(3)$ we have $T_G = 3$ and $T_R = 6$, such that for the different gaugino masses this gives (in a self-explanatory notation):

$$\begin{aligned} M_1 &= 11\frac{g_Y^2}{16\pi^2}m_{3/2}\left[1-(\rho+1)^2-\frac{t(\gamma+t+\gamma t)}{1+\gamma t}\right], \\ M_2 &= \frac{g_2^2}{16\pi^2}m_{3/2}\left[1-5(\rho+1)^2-5\frac{t(\gamma+t+\gamma t^2)}{1+\gamma t}\right], \\ M_3 &= -3\frac{g_3^2}{16\pi^2}m_{3/2}\left[1+(\rho+1)^2+\frac{t(\gamma+t+\gamma t^2)}{1+\gamma t}\right]. \end{aligned} \tag{4.14}$$

5. Phenomenology

The results for the soft terms calculated in the previous section, evaluated for different values of the parameter γ are summarised in Table 1. For every γ, the corresponding t and ρ are calculated by imposing a vanishing cosmological constant at the minimum of the potential. The scalar soft masses and trilinear terms are then evaluated by eqs. (4.6) and the gaugino masses by eqs. (4.14). Note that the relation (4.7) is valid for all γ. We therefore do not list the parameter B_0.

γ	t	ρ	m_0	A_0	M_1	M_2	M_3	$\tan\beta(\mu>0)$	$\tan\beta(\mu<0)$
0.6	0.446	0.175	0.475	1.791	0.017	0.026	0.027		
1	0.409	0.134	0.719	1.719	0.015	0.025	0.026		
1.1	0.386	0.120	0.772	1.701	0.015	0.024	0.026	46	29
1.4	0.390	0.068	0.905	1.646	0.014	0.023	0.026	40	23
1.7	0.414	0.002	0.998	1.588	0.013	0.022	0.025	36	19

Table 1: The soft terms (in terms of $m_{3/2}$) for various values of γ. If a solution to the RGE exists, the value of $\tan\beta$ is shown in the last columns for $\mu > 0$ and $\mu < 0$ respectively.

In most phenomenological studies, B_0 is substituted for $\tan\beta$, the ratio between the two Higgs VEVs, as an input parameter for the renormalization group equations (RGE) that determine the low energy spectrum of the theory. Since B_0 is not a free parameter in our theory, but is fixed by eq. (4.7), this corresponds to a definite value of $\tan\beta$. For more details see [18] (and references therein). The corresponding $\tan\beta$ for a few particular choices for γ are listed in the last two columns of table 1 for $\mu > 0$ and $\mu < 0$ respectively. No solutions were found for $\gamma \lesssim 1.1$, for both signs of μ. The lighest supersymmetric particle (LSP) is given by the lightest neutralino and since $M_1 < M_2$ (see table 1) the lightest neutralino is mostly Bino-like, in contrast with a typical mAMSB (minimal anomaly mediation supersymmetry breaking) scenario, where the lightest neutralino is mostly Wino-like [19].

To get a lower bound on the stop mass, the sparticle spectrum is plotted in Figure 2 as a function of the gravitino mass for $\gamma = 1.1$ and $\mu > 0$ (for $\mu < 0$ the bound is higher). The experimental limit on the gluino mass forces $m_{3/2} \gtrsim 15$ TeV. In this limit the stop mass can be as low as 2 TeV. To conclude, the lower end mass spectrum consists of (very) light charginos (with a lightest chargino between 250 and 800 GeV) and neutralinos, with a mostly Bino-like neutralino as LSP ($80-230$

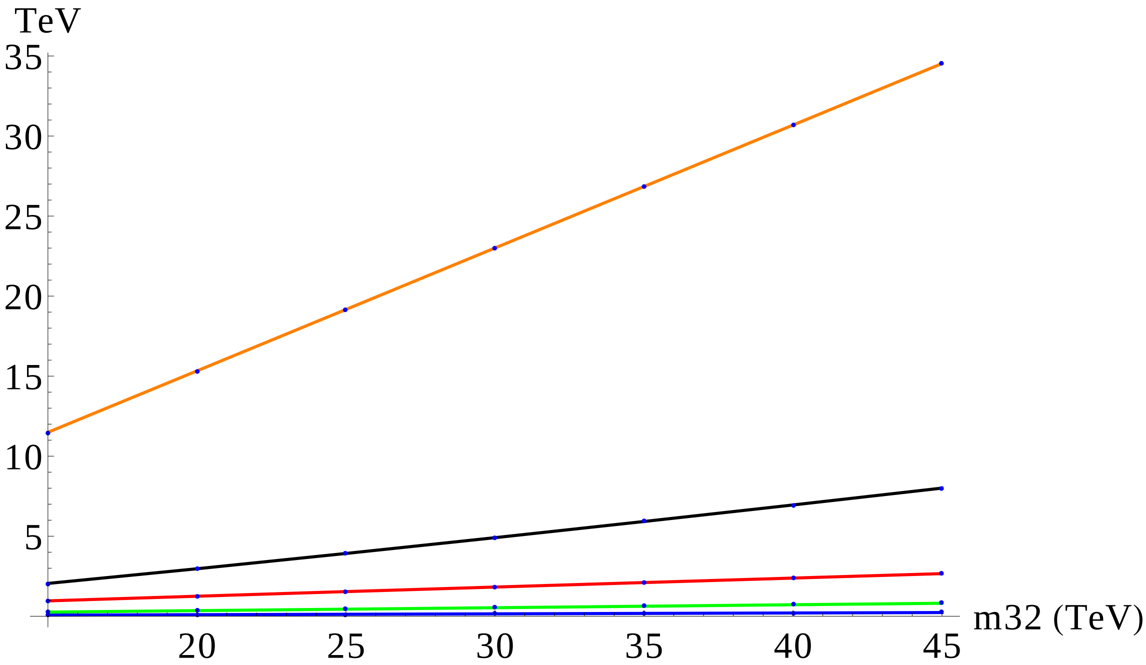

Figure 2: The masses (in TeV) of the sbottom (yellow), stop (black), gluino (red), lightest chargino (green) and lightest neutralino (blue) as a function of $m_{3/2}$ for $\gamma = 1.1$ and for $\mu > 0$. No solutions to the RGE were found when $m_{3/2} \gtrsim 45$ TeV. The lower bound corresponds to a gluino mass of 1 TeV.

GeV), which would distinguish this model from the mAMSB where the LSP is mostly Wino-like. These upper limits on the LSP and the lightest chargino imply that this model could in principle be excluded in the next LHC run. In order for the gluino to escape experimental bounds, the lower limit on the gravitino mass is about 15 TeV. The gluino mass is then between 1-3 TeV. This however forces the squark masses to be very high $(10-35$ TeV), with the exception of the stop mass which can be relatively light $(2-15$ TeV).

6. Non-canonical Kähler potential for the visible sector

As mentioned already in Section 4, an alternative way to avoid tachyonic soft scalar masses for the MSSM fields in the model (4.1), instead of adding the extra Palonyi-type field z in the hidden sector, is by introducing non-canonical kinetic terms for the MSSM fields, such as:

$$\begin{aligned} \mathcal{K} &= -\kappa^{-2}\log(s+\bar{s}) + \kappa^{-2}b(s+\bar{s}) + (s+\bar{s})^{-\nu}\sum\varphi\bar{\varphi}, \\ W &= \kappa^{-3}a + W_{MSSM}, \\ f(s) &= 1, \qquad f_A(s) = 1/g_A^2, \end{aligned} \tag{6.1}$$

where ν is an additional parameter of the theory, with $\nu = 1$ corresponding to the leading term in the Taylor expansion of $-\log(s+\bar{s}-\varphi\bar{\varphi})$. Since the visible sector fields appear only in the combination $\varphi\bar{\varphi}$, their VEVs vanish provided that the scalar soft masses squared are positive. Moreover, for vanishing visible sector VEVs, the scalar potential and is minimization remains the same as in eqs. (refbsalpha). Therefore, the non-canonical Kähler potential does not change the fact that the F-term contribution to the soft scalar masses squared is negative. On the other hand, the visible fields enter in the D-term scalar potential through the derivative of the Kähler potential with respect to s. Even though this has no effect on the ground state of the potential, the φ-dependence of the D-term scalar potential does result in an extra contribution to the scalar masses squared which become positive

$$\nu > -\frac{e^{\alpha}(\sigma_s+1)\alpha}{A(\alpha)(1-\alpha)} \approx 2.6\,. \tag{6.2}$$

The soft MSSM scalar masses and trilinear couplings in this model are:

$$\begin{aligned} m_0^2 &= \kappa^2 a^2 \left(\frac{b}{\alpha}\right)\left(e^{\alpha}(\sigma_s+1)+\nu\frac{A(\alpha)}{\alpha}(1-\alpha)\right) \\ A_0 &= m_{3/2}(s+\bar{s})^{\nu/2}(\sigma_s+3) \\ B_0 &= m_{3/2}(s+\bar{s})^{\nu/2}(\sigma_s+2) \end{aligned} \tag{6.3}$$

where σ_s is defined as in (4.6), eq. (4.4) has been used to relate the constants a and c, and corrections due to a small cosmological constant have been neglected. A field redefinition due to a non-canonical kinetic term $g_{\varphi\bar{\varphi}} = (s+\bar{s})^{-\nu}$ is also taken into account. The main phenomenological properties of this model are not expected to be different from the one we analyzed in section 5 with the parameter ν replacing γ. Gaugino masses are still generated at one-loop level while mSUGRA applies to the soft scalar sector. We therefore do not repeat the phenomenological analysis for this model.

7. Identifying the dilaton with the inflaton

In the following, we study the possibility to identify the dilaton with the inflaton. We will show first that the above model does not allow slow roll inflation.

Indeed, the kinetic terms in the model (3.2-3.3) for the scalar $\phi \equiv s+\bar{s} = 1/l$ are given by

$$\mathcal{L}_s/e = -g_{s\bar{s}}\partial_\mu s\partial^\mu \bar{s} = -\frac{p\kappa^{-2}}{4}\frac{1}{\phi^2}\partial_\mu\phi\partial^\mu\phi. \tag{7.1}$$

The canonically normalised field χ therefore satisfies $\chi = \kappa^{-1}\sqrt{\frac{p}{2}}\log\phi$, where we re-introduce the gravitational coupling κ.

The slow roll parameters are given by

$$\begin{aligned} \varepsilon &= \frac{1}{2\kappa^2}\left(\frac{dV/d\chi}{V}\right)^2 = \frac{1}{2\kappa^2}\left[\frac{1}{V}\frac{dV}{d\phi}\left(\frac{d\chi}{d\phi}\right)^{-1}\right]^2, \\ \eta &= \frac{1}{\kappa^2}\frac{V''(\chi)}{V} = \frac{1}{\kappa^2}\frac{1}{V}\left[\frac{d^2V}{d\phi^2}\left(\frac{d\chi}{d\phi}\right)^{-2} - \frac{dV}{d\phi}\frac{d^2\chi}{d\phi^2}\left(\frac{d\chi}{d\phi}\right)^{-3}\right], \end{aligned} \tag{7.2}$$

It can be shown that, when the conditions (3.4) and (3.5) are satisfied, the slow roll parameters and the potential depend only on $\rho = -b\phi$; indeed

$$\frac{\kappa^4 V(\rho)}{b^3c^2} = \frac{e^{-\rho}\left(A_2(\alpha)\rho\left(\rho^2+4\rho-2\right)-2e^{\rho}(\rho+2)^2\right)}{2\rho^3}, \tag{7.3}$$

where $A_2(\alpha) \approx -50.66$ as in eq. (3.5). In Fig. 3, a plot is shown of $\frac{\kappa^4 V(\rho)}{|b|^3 c^2}$ as a function of ρ. The minimum of the potential is at $\rho_{\text{min}} \approx 0.1832$ (see eq. (3.4)), while the potential has a local maximum at $\rho_{\text{max}} \approx 0.4551$. A plot of the slow roll parameter η (also in Fig. 3) shows that $|\eta| \ll 1$ is not satisfied. This result holds for any parameters a, b, c satisfying eqs. (3.4) and (3.5). A similar

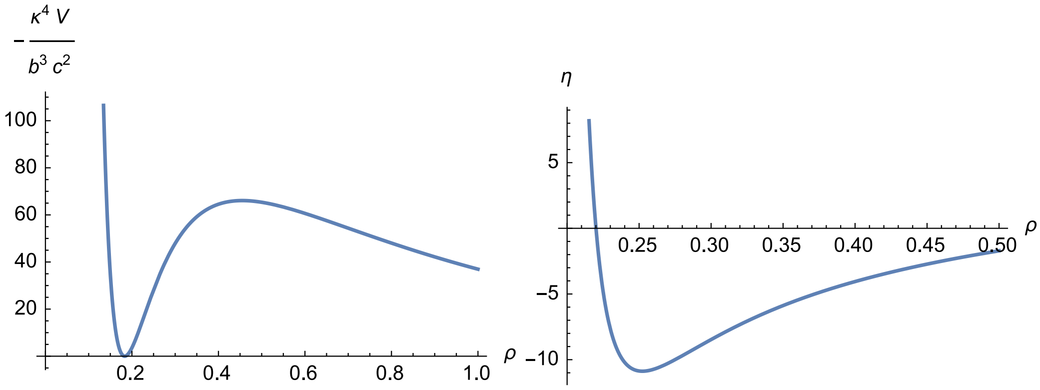

Figure 3: A plot of $-\frac{\kappa^4 V(\rho)}{b^3 c^2}$ as a function of $\rho = -b\phi$ (left), and a plot of the slow roll parameter η as a function of ρ (right). The slow roll condition $|\eta| \ll 1$ is not satisfied for any value of the parameters a, b, c.

analysis to the one above can be performed for $p = 1$, showing that the slow roll condition $\eta \ll 1$ can not be satisfied.

8. Extensions of the model that satisfy the slow roll conditions

In the previous section we showed that the slow roll conditions can not be satisfied in the minimal versions of the model. In this section we modify the above model by modifying the Kähler potential. While the superpotential is uniquely fixed (up to a Kähler transformation), the Kähler potential admits corrections that can always be put in the form

$$\mathcal{K} = -p\kappa^{-2} \log\left(s + \bar{s} + \frac{\xi}{b} F(s+\bar{s})\right) + \kappa^{-2} b(s+\bar{s}), \tag{8.1}$$

while the superpotential, the gauge kinetic function and moment map are given by

$$\begin{aligned} W &= \kappa^{-3} a, \\ f(s) &= \delta + \beta s, \\ \mathcal{P} &= \kappa^{-2} c \left(b - p \frac{1 + \frac{\xi}{b} F_s}{s + \bar{s} + \frac{\xi}{b} F} \right), \end{aligned} \tag{8.2}$$

where $\mathcal{P}$ is the $U(1)$ moment map (2.7) and $F_s = \partial_s F(s+\bar{s})$. The scalar potential is given by ($\phi = s + \bar{s}$)

$$\begin{aligned} V &= V_F + V_D, \\ V_F &= \kappa^{-4} \frac{|a|^2 e^{b\phi}}{(\phi + \frac{\xi}{b} F)^p} \left[-3 - \frac{1}{p} \frac{\left(b(b\phi + \xi F) - p(b + \xi F_\phi)\right)^2}{\xi F_{\phi\phi}(b\phi + \xi F) - (b + \xi F_\phi)^2} \right], \\ V_D &= \kappa^{-4} \frac{b^2 c^2}{2\delta + \beta\phi} \left[1 - p \frac{1 + \frac{\xi}{b} F_\phi}{b\phi + \xi F} \right]^2. \end{aligned} \tag{8.3}$$

As was discussed above, we take $\delta = 1$, $\beta = 0$ for $p = 1$ and $\delta = 0$, $\beta = 1$ for $p = 2$.

Identifying $\mathrm{Re}(s)$ with the inverse string coupling, the function F may contain perturbative contributions that can be expressed as power series of $1/(s+\bar{s})$, as well as non-perturbative corrections which are exponentially suppressed in the weak coupling limit. The later can be either of the form $e^{-\lambda(s+\bar{s})}$ for $\lambda > 0$ in the case of D-brane instantons, or of the form $e^{-\lambda(s+\bar{s})^2}$ in the case of (Neveu-Schwarz) NS5-brane instantons (since the closed string coupling is the square of the open string coupling). We have considered a generic contribution of these three different types of corrections and we found that only the last type of contributions can lead to an inflationary plateau providing sufficient inflation. The other corrections can be present but do not modify the main properties of the model (as long as weak coupling description holds). In the following section, we analyse in detailed a function F describing a generic NS5-brane instanton correction to the Kähler potential.

9. Slow-roll Inflation

9.1 p=2 case

We now consider the case with

$$F(\phi) = \exp(\alpha b^2\phi^2), \tag{9.1}$$

where $b < 0$ and $\alpha < 0$. $F(\phi)$ vanishes asymptotically at large ϕ. In this case, we obtain

$$V_D = \frac{\kappa^{-4}b^3c^2}{b\phi}\left[\frac{b\phi - 2 + \xi e^{\alpha b^2\phi^2}(1-4\alpha b\phi)}{b\phi + \xi e^{\alpha b^2\phi^2}}\right]^2, \tag{9.2}$$

and

$$V_F = -\frac{\kappa^{-4}|a|^2b^2e^{b\phi}}{2\left(\xi e^{\alpha b^2\phi^2} + b\phi\right)^2}\left[\frac{\left(b\phi + \xi e^{\alpha b^2\phi^2}(1-4\alpha b\phi) - 2\right)^2}{2\alpha\xi e^{\alpha b^2\phi^2}\left(2\alpha b^3\phi^3 + \xi e^{\alpha b^2\phi^2} - b\phi\right) - 1} + 6\right]. \tag{9.3}$$

There are four parameters in this model namely α, ξ, b and c. The first two parameters α and ξ control the shape of the potential. There are some regions in the parameter space of α and ξ that the potential satisfies the slow-roll conditions i.e. $\varepsilon \ll 1$ and $|\eta| \ll 1$. In order to obtain the potential with flat plateau shape which is suitable for inflation and in agreement with Planck '15 data, we choose

$$\alpha \simeq -4.84 \quad \text{and} \quad \xi \simeq 0.025 \tag{9.4}$$

Note that in the case of $\xi = 0$ and $b < 0$, we can find the Minkowski minimum by solving the equations $V(\phi_{min}) = 0$ and $dV(\phi_{min})/d\phi = 0$, where $\phi_{min} = s_{min} + \bar{s}_{min}$ is the value of ϕ at the minimum of the potential. In the case of $\xi \neq 0$, we can not solve the equations analytically and the relations (3.4), (3.5) are not valid. We can always assume that they are modified into

$$b\phi_{min} = -\rho(\xi,\alpha) \quad \text{and} \quad \frac{a^2}{bc^2} = -50.66\times\lambda(\xi,\alpha,\Lambda)^2, \tag{9.5}$$

where λ takes positive values and satisfies $|\lambda - 1| \ll 1$. For any given value of parameters ξ, α and the cosmological constant Λ, one can numerically fix the value of ρ and λ. By fine-tuning the cosmological constant Λ to be very close to zero, we can numerically solve the equations $V = 0$ and $dV/d\phi = 0$ for the value of ρ and λ in (9.5) as:

$$\rho \approx 0.18, \tag{9.6}$$

$$\lambda \approx 1.017 \tag{9.7}$$

From eq. (9.5), we can see that the third parameter, b, controls the vacuum expectation value ϕ_{min}. This can be shown in Fig. 4 where we compare the scalar potential for different values of b. Motivated by string theory, we have the identification $\phi \sim 1/g_s$. We can choose the value of the parameter b such that ϕ_{min} is of the order of 10 to make sure that we are in the perturbative regime in g_s. The last parameter, c, controls the overall scale of the potential but does not change its minimum and its shape. In the following, we will fix b and c by using the cosmological data.

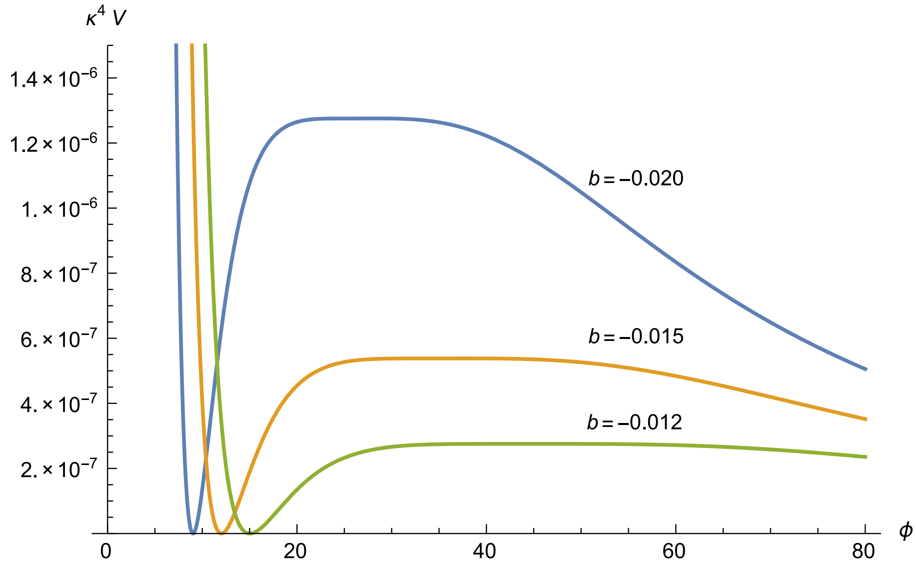

Figure 4: A plot of the scalar potential for $p = 2$, with $b = -0.020$, $b = -0.015$ and $b = -0.012$. Note that we choose the parameters α and ξ as in eq. (9.4) with $c = 0.06$.

In order to compare the predictions of our models with Planck '15 data, we choose the following boundary conditions:

$$\phi_{int} = 27.32 \qquad \phi_{end} = 22.68 \tag{9.8}$$

The initial conditions are chosen very near the maximum on the (left) side, so that the field rolls down towards the electroweak minimum. Any initial condition on the right of the maximum may produce also inflation, but the field will roll towards the SUSY vacuum at infinity. The results are therefore very sensitive to the initial conditions (9.8) of the inflaton field.

The slow roll parameters are given as in equation (7.2). The total number of e-folds N can be determined by

$$N = \kappa^2 \int_{\chi_{end}}^{\chi_{int}} \frac{V}{\partial_\chi V} d\chi = \kappa^2 \int_{\phi_{end}}^{\phi_{int}} \frac{V}{\partial_\phi V} \left(\frac{d\chi}{d\phi}\right)^2 d\phi. \tag{9.9}$$

n_s	r	A_s
0.965	2.969×10^{-23}	2.259×10^{-9}

Table 2: The theoretical predictions for $p = 2$, with $b = -0.0182$ and $c = 0.61 \times 10^{-13}$, where α and ξ are given in eq. (9.4).

Note that we choose $|\eta(\chi_{end})| = 1$. We can compare the theoretical predictions of our model to the experimental results via the power spectrum of scalar perturbations of the CMB, namely the amplitude A_s and tilt n_s, and the relative strength of tensor perturbations, i.e. the tensor-to-scalar ratio r. In terms of slow roll parameters, these are given by

$$A_s = \frac{\kappa^4 V_*}{24\pi^2 \varepsilon_*}, \tag{9.10}$$

$$n_s = 1 + 2\eta_* - 6\varepsilon_*, \tag{9.11}$$

$$r = 16\varepsilon_*, \tag{9.12}$$

where all parameters are evaluated at the field value χ_{int}.

In order to satisfy Planck '15 data, we choose the parameters $b = -0.0182$, $c = 0.61 \times 10^{-13}$. The value of the slow-roll parameters at the beginning of inflation are

$$\varepsilon(\phi_{int}) \simeq 1.86 \times 10^{-24} \quad \text{and} \quad \eta(\phi_{int}) \simeq -1.74 \times 10^{-2}. \tag{9.13}$$

The total number of e-folds N, the scalar power spectrum amplitude A_s, the spectral index of curvature perturbation n_s and the tensor-to-scalar ratio r are calculated and summarised in Table 2, in agreement with Planck '15 data [20]. Fig. 5 shows that our predictions for n_s and r are within 1σ C.L. of Planck '15 contours with the total number of e-folds $N \approx 1075$. Note that N is the total number of e-folds from ϕ_{int} to ϕ_{end}. However the number of e-folds associated with the CMB observation corresponds to a period between the time of horizon crossing and the end of inflation, which is much smaller than 1075. According to general formula in [20], the number of e-folds between the horizon crossing and the end of inflation is roughly estimated to be around 50-60.

We would like to remark that the parameter c also controls the gravitino mass at the minimum of the potential around $O(10)$ TeV. Indeed, the gravitino mass is written as

$$m_{3/2} = \kappa^2 e^{\kappa^2 \mathscr{K}/2} W = \frac{1}{\kappa}\left(\frac{a b e^{b\phi/2}}{b\phi + \xi F(\phi)}\right). \tag{9.14}$$

For $b = -0.0182$, we get $\phi_{\min} \approx 9.91134$ and the gravitino mass at the minimum of the potential

$$\langle m_{3/2} \rangle \approx 14.98 \ \text{TeV}. \tag{9.15}$$

The Hubble parameter during inflation (evaluated at $\phi_* = \phi_{int}$) is

$$H_* = \kappa\sqrt{V_*/3} = 1.38 \ \text{TeV}. \tag{9.16}$$

This shows that our predicted scale for inflation is of the order of TeV. The mass of gravitino during the inflation $m^*_{3/2} = 4.15$ TeV is higher than the inflation scale, and the gauge boson mass is $M^*_{A_\mu} = 3.12$ TeV.[5] In fact, the gauge boson acquires a mass due to a Stueckelberg mechanism by eating the imaginary component of s, where its mass at the minimum of the potential is given by

$$\langle M_{A_\mu} \rangle = 15.48\text{TeV}. \tag{9.17}$$

As a result, the model essentially contains only one scalar field $\mathrm{Re}(s)$, which is the inflaton. This is in contrast with other supersymmetric models of inflation, which usually contain at least two real scalars [21].[6]

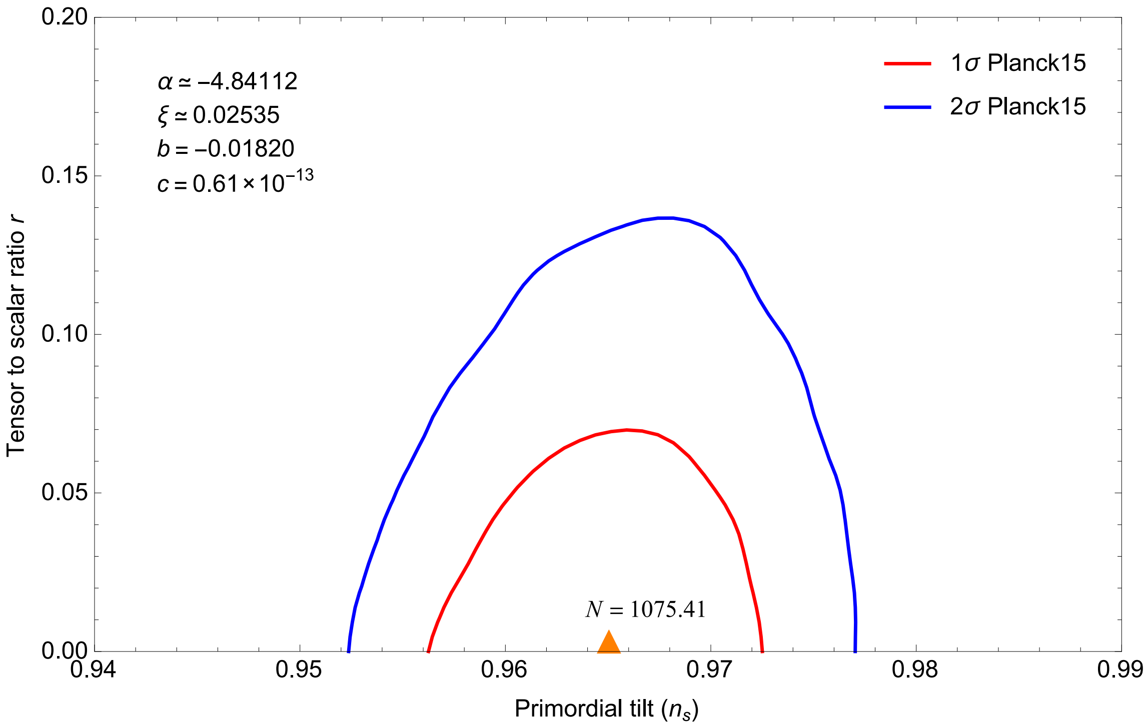

Figure 5: We plot the theoretical predictions for the case $p = 2$, shown in Table 2, in the n_s - r plane together with the Planck '15 results for TT, TE, EE, + lowP and assuming ΛCDM + r [20].

9.2 p=1 case

In this case, we obtain

$$V_D = \frac{\kappa^{-4} b^2 c^2}{2} \left[\frac{b\phi - 1 + \xi e^{\alpha b^2 \phi^2}(1 - 2\alpha b\phi)}{b\phi + \xi e^{\alpha b^2 \phi^2}} \right]^2, \tag{9.18}$$

and

$$V_F = -\frac{\kappa^{-4} |a|^2 b e^{b\phi}}{\xi e^{\alpha b^2 \phi^2} + b\phi} \left[\frac{\left(b\phi + \xi e^{\alpha b^2 \phi^2}(1 - 2\alpha b\phi) - 1 \right)^2}{2\alpha \xi e^{\alpha b^2 \phi^2} \left(2\alpha b^3 \phi^3 + \xi e^{\alpha b^2 \phi^2} - b\phi \right) - 1} + 3 \right]. \tag{9.19}$$

[5] The gauge boson mass is given by $m_{A_\mu} = \sqrt{2 g_{s\bar{s}} c^2 / \mathrm{Re}(s)}$.
[6] This is because a chiral multiplet contains a complex scalar.

n_s	r	A_s
0.959	4.143×10^{-22}	2.205×10^{-9}

Table 3: The theoretical predictions for $p=1$ case with $b=-0.0234$, $c=1\times10^{-13}$, $\alpha=-0.781$ and $\xi=0.3023$.

The potential has similar properties with the $p=2$ case although it may give different phenomenological results at low energy. Similar to the previous case, the relations (4.4) are not valid when $\xi\neq0$ and we assume that they are modified into

$$b\phi_{min}=-\rho(\xi,\alpha)\quad\text{and}\quad\frac{bc^2}{a^2}\simeq-0.359\times\lambda(\xi,\alpha,\Lambda)^{-2}. \tag{9.20}$$

By choosing $\alpha=-0.781$ and $\xi=0.3023$ and tuning the cosmological constant Λ to be very close to zero, we can numerically fix $\rho\approx0.56$ and $\lambda\approx1.29$ for this case. The gravitino mass for $p=1$ case can be written as

$$m_{3/2}=\kappa^2e^{\kappa^2\mathscr{K}/2}W=\frac{1}{\kappa}\left(\frac{a\sqrt{b}e^{b\phi/2}}{\sqrt{b\phi+\xi F(\phi)}}\right). \tag{9.21}$$

By choosing the parameters $b=-0.0234$, $c=1\times10^{-13}$, the gravitino mass at the minimum of the potential is

$$\langle m_{3/2}\rangle=18.36\ \text{TeV}. \tag{9.22}$$

with $\phi_{\min}\approx21.53$, and

$$\langle M_{A_\mu}\rangle=36.18\ \text{TeV}. \tag{9.23}$$

By choosing appropriate boundary conditions, we find

$$\phi_{int}=64.53\quad\text{and}\quad\phi_{end}=50.99 \tag{9.24}$$

As summarised in Table 3, the predictions for the $p=1$ case are similar to those of $p=2$, in agreement with Planck '15 data with the total number of e-folds $N\approx888$. In this case, the Hubble parameter during inflation is

$$H_*=\kappa\sqrt{V_*/3}=5.09\ \text{TeV}. \tag{9.25}$$

Note that for the $p=1$ case, the mass of the gauge boson is $M^*_{A_\mu}=6.78$ TeV, and the mass of the gravitino during inflation is $m^*_{3/2}=4.72$ TeV.

9.3 SUGRA spectrum

The above model can be coupled to MSSM, as described in section 4:

$$\begin{aligned}\mathscr{K}&=\mathscr{K}(s+\bar{s})+\sum\varphi\bar{\varphi},\\ W&=W_h(s)+W_{\text{MSSM}}.\end{aligned}$$

(9.27)

The soft supersymmetry breaking terms can then be calculated as follows

$$\begin{aligned} m_0^2 &= e^{\kappa^2 \mathscr{K}} \left(-2\kappa^4 W_h(s)\bar{W}_h(s) + \kappa^2 g^{s\bar{s}} |\nabla_s W_h|^2 \right), \\ A_0 &= \kappa^2 e^{\kappa^2 \mathscr{K}/2} g^{s\bar{s}} K_s \left(\bar{W}_{\bar{s}} + \kappa^2 K_s \bar{W} \right), \\ B_0 &= \kappa^2 e^{\kappa^2 \mathscr{K}/2} \left(g^{s\bar{s}} K_s \left(\bar{W}_{\bar{s}} + \kappa^2 K_s \bar{W} \right) - \bar{W} \right). \end{aligned} \tag{9.28}$$

For $p = 2$ the Lagrangian contains a Green-Schwarz term eq. (4.2), and the theory is not gauge invariant (without the inclusion of extra fields that are charged under the $U(1)$). We therefore focus on $p = 1$. The soft terms can be written in terms of the gravitino mass (see eq. (9.14))

$$\begin{aligned} m_0^2 &= m_{3/2}^2 [-2 + \mathscr{C}], \\ A_0 &= m_{3/2}\, \mathscr{C}, \\ B_0 &= A_0 - m_{3/2}, \end{aligned} \tag{9.29}$$

where

$$\mathscr{C} = - \left. \frac{\left(-\xi e^{\alpha b^2 \phi^2} + b\phi \left(4\alpha\xi e^{\alpha b^2 \phi^2} - 1 \right) + 2 \right)^2}{4\alpha\xi^2 e^{2\alpha b^2 \phi^2} - 4\alpha b \xi \phi e^{\alpha b^2 \phi^2} + 8\alpha^2 b^3 \xi \phi^3 e^{\alpha b^2 \phi^2} - 2} \right|_{\phi = \phi_{\min}} . \tag{9.30}$$

Using the parameters presented in section 9.2, we find $m_{3/2} = 18.36$ TeV and $\mathscr{C} = 1.53$. For $\xi = 0$ the model reduces to the one analysed in section 4, where one has $\mathscr{C} = 1.52$ and $m_{3/2} = 17.27$ TeV (with $\phi_{\min} = 9.96$). Moreover, the scalar soft mass is tachyonic. This can be solved either by introducing an extra Polonyi-like field, or by allowing a non-canonical Kähler potential for the MSSM-like fields φ. The resulting low energy spectrum is expected to be similar to the one described in sections 4 and 5. We do not perform this analysis, but only summarise the results.

Since the tree-level contribution to the gaugino masses vanishes, their mass is generated at one-loop by the so-called 'Anomaly Mediation' contribution (4.10). As a result, the spectrum consists of very light neutralinos ($O(10^2)$ GeV), of which the lightest (a mostly Bino-like neutralino) is the LSP dark matter candidate, slightly heavier charginos and a gluino in the $1-3$ TeV range. The squarks are of the order of the gravitino mass (~ 10 TeV), with the exception of the stop squark which can be as light as 2 TeV.

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