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Supergravity origin of the MSSM inflation

Kari Enqvist[1,2,*], **Lotta Mether**[1,2,†] **and Sami Nurmi**[2,‡]

[1] *Helsinki Institute of Physics, P.O. Box 64, FIN-00014 University of Helsinki, Finland*
[2] *Department of Physical Sciences, P.O. Box 64, FIN-00014 University of Helsinki, Finland*

Abstract: We consider the supergravity origin of the recently proposed MSSM inflationary model, which relies on the existence of a saddle point along a dimension six flat direction. We derive the conditions that the Kähler potential has to satisfy for the saddle point to exist irrespective of the hidden sector vevs. We show that these conditions are satisfied by a simple class of Kähler potentials, which we find to have a similar form as in various string theory compactifications. For these potentials, slow roll MSSM inflation requires no fine tuning of the soft supersymmetry breaking parameters.

Keywords: Cosmology, Inflation, Supergravity, MSSM.

*E-mail: kari.enqvist@helsinki.fi
†E-mail: lotta.mether@helsinki.fi
‡E-mail: sami.nurmi@helsinki.fi

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1. Introduction

Although conventionally inflation is assumed to take place at a very high energy scale and be driven by the slow roll motion of an unknown singlet field, the inflaton, recently it has been pointed out [1, 2] that inflation can in fact be realized already within the Minimally Supersymmetric Standard Model (MSSM). In this case the inflaton is a gauge invariant combination of squark or slepton fields. The flatness of the inflaton potential is provided by supersymmetry and the gauge symmetries, which together give rise to about 300 flat directions in the space of scalar fields [3] (for a review of the physics of the MSSM flat directions, see [4]). Along these flat directions, in the limit of exact supersymmetry (susy), the scalar potential vanishes identically. However, the flat directions are lifted by non-renormalizable superpotential terms, as well as by soft susy breaking [5]; of particular importance for MSSM inflation are the non-renormalizable A-terms $\propto W_n$, where W_n is the non-renormalizable superpotential of dimension n. The dimensionality of the non-renormalizable terms depends on the particular direction, as does the existence of the A-terms, which are absent for some directions.

As discussed in [1], phenomenologically acceptable slow roll MSSM inflation can arise along the dimension six flat directions **udd** and **LLe**, which we denote by the field Φ. The flat direction field is complex, and in the complex plane there exists a set of discrete directions along which the contribution of the A-term is most negative. Along these directions the MSSM inflaton potential reads

$$V = \frac{1}{2}m^2\phi^2 - \frac{A\lambda}{6}\phi^6 + \lambda^2\phi^{10} \ , \tag{1.1}$$

where ϕ is the absolute value of the field, m and A are the soft susy breaking terms, λ is an effective coupling constant and we have set $M_P \equiv 1$. Generically, the potential Eq. (1.1) does not as such give rise to inflation. However, one may notice that it has a secondary minimum at

$$\phi_0 = \left(\frac{A}{20\lambda}\right)^{1/4} \ll 1 \ , \tag{1.2}$$

which becomes a saddle point if the condition

$$A^2 = 40m^2 \tag{1.3}$$

holds. In that case the potential is extremely flat with $V'(\phi_0) = V''(\phi_0) = 0$. In the vicinity of the saddle point the potential is given by

$$V(\phi) \approx V(\phi_0) + \frac{1}{6}V'''(\phi_0)(\phi - \phi_0)^3 = V(\phi_0) + \frac{16}{3}\frac{m^2}{\phi_0}(\phi - \phi_0)^3 \ . \tag{1.4}$$

If in the initial state $\phi \simeq \phi_0$, there follows a period of slow roll inflation with a very low scale of $H_{inf} \sim 1 - 10$ GeV, assuming $\lambda \sim O(1)$, and a spectral index of $n \simeq 0.92$ [1]. Slight deviations from the saddle point condition Eq. (1.3) modify the spectral index somewhat (see [2]). Because of the low inflationary scale, there are no observable tensor perturbations.

The great virtue of the MSSM inflation is that the inflaton couplings to Standard Model particles are known and, at least in principle, measurable in laboratory experiments such as LHC or a future Linear Collider. The inflaton mass is directly related to the slepton or squark masses and the model can thus be tested in the laboratory.

However, the obvious disadvantage is the fine tuning implicit in the saddle point condition Eq. (1.3). Slow roll inflation[1] requires that the ratio A/m should be tuned to the saddle point with an accuracy of about 10^{-16}; otherwise the slow roll properties of the potential Eq. (1.1) would be spoiled [2]. Since in the MSSM the soft susy breaking parameters are put in by hand, there can be no explanation for the saddle point condition other than simple finetuning. Thus the relation Eq. (1.3) must reflect physics that is beyond the MSSM and in particular the mechanism of supersymmetry breaking. Hence the values of the soft susy breaking parameters reflect the properties of the hidden sector. The question then is: is it possible to realize the saddle point condition naturally in some supergravity model as defined by the Kähler potential? This means that the condition Eq. (1.3) should not be just an accidental coincidence that emerges when the hidden sector fields settle in their vevs, but rather a generic condition that holds irrespective of the hidden sector field values. In the present paper we demonstrate that this is indeed the case. Moreover, the form of the Kähler potential turns out to be rather suggestive, with features that can be found in certain string theoretical compactification schemes.

[1] For a discussion of the inflationary properties of a potential that has the generic form of Eq. (1.1), see [6]; for a discussion on dark matter and MSSM inflation, see [7].

2. The scalar potential

Our aim is to identify a class of Kähler potentials that generate such soft susy breaking terms for the flat direction Φ that the saddle point condition (1.3) is identically satisfied. It is obvious that the simplest, flat Kähler potential will not do the job; instead, one has to consider more complicated possibilities. We focus on soft terms generated through F-term susy breaking (recall that the flat directions are D-flat also in supergravity). In this case the scalar potential is determined solely by the function

$$G(\Phi_M, \Phi_M^*) = K(\Phi_M, \Phi_M^*) + \log |W(\Phi_M)|^2 \ , \tag{2.1}$$

where K and W are respectively the Kähler potential and the superpotential. Here Φ_M, which includes both the hidden sector fields h_m and the flat direction inflaton field Φ, denotes the scalar part of the corresponding chiral superfield. Assuming vanishing D-terms also in the hidden sector, the tree-level scalar potential reads

$$V = e^G \Big(G^{M\bar{N}} G_M G_{\bar{N}} - 3 \Big) \ , \tag{2.2}$$

where the lower indices $_M$ and $_{\bar{M}}$ refer to derivatives with respect to Φ_M and Φ_M^*, and the matrix $G^{M\bar{N}} = K^{M\bar{N}}$ is the inverse of the Kähler metric $G_{M\bar{N}} = K_{M\bar{N}}$.

For the dimension 6 flat directions that we are considering as the inflaton, the superpotential is of the form

$$W = \hat{W} + \frac{\hat{\lambda}_6}{6} \Phi^6 \equiv \hat{W} + I \ , \tag{2.3}$$

where I is the lowest order non-renormalizable term that lifts the flat direction. The superpotential may also contain all possible higher order terms allowed by symmetries but these will not affect our analysis and have therefore been suppressed. In Eq. (2.3) and elsewhere in the text, we use the hat to denote quantities that are independent of Φ, but are functions of the hidden sector fields. This is in general the case for the Φ-independent term $\hat{W}$ of the superpotential, as well as for the coupling constant $\hat{\lambda}_6$ of the non-renormalizable term. However, since our focus is on finding a Kähler potential that satisfies the relation Eq. (1.3), we will neglect the hidden sector dependence of the superpotential, and hence treat these quantities as constants throughout this paper. In this context, it is worth noting that, in order to ensure the validity of the MSSM inflation scenario, we are assuming the flat direction to be the only dynamical variable during inflation. Thus we are implicitly requiring that the hidden sector fields are stabilized before the beginning of inflation either by the neglected superpotential terms or through some other mechanism.[2]

[2]The inclusion of the hidden sector fields, even when stabilized before the onset of inflation, may in general lead to additional fine-tuning conditions on the inflationary potential. A detailed analysis of these effects however requires precise knowledge of the nature and dynamics of the hidden sector fields and is beyond the scope of this paper. A qualitative discussion of this issue can be found in Ref. [8], even if the results therein as such are not applicable here.

Given the superpotential Eq. (2.3), the scalar potential Eq. (2.2) can be written as

$$V = |\hat{W}|^2 f + \hat{W} I^* g + \hat{W}^* I g^* + |I_\phi|^2 k \ , \tag{2.4}$$

where

$$f = e^K \Big(K^{M\bar{N}} K_M K_{\bar{N}} - 3 \Big) \ , \tag{2.5}$$

$$g = e^K \Big(\frac{6}{\Phi^*} K^{M\bar{\phi}} K_M + K^{M\bar{N}} K_M K_{\bar{N}} - 3 \Big) \ , \tag{2.6}$$

$$k = e^K \Big(K^{\phi\bar{\phi}} + \frac{\Phi}{6} K^{M\bar{\phi}} K_M + \frac{\Phi^*}{6} K^{\phi\bar{M}} K_{\bar{M}} + \frac{\Phi\Phi^*}{36} (K^{M\bar{N}} K_M K_{\bar{N}} - 3) \Big) \ . \tag{2.7}$$

To find the explicit expression for the potential Eq. (2.4), one needs to determine the Kähler potential. Here we consider Kähler potentials of the generic perturbative form

$$K = \hat{K} + \hat{Z}_2 \phi^2 + \hat{Z}_4 \phi^4 + \hat{Z}_6 \phi^6 + \dots \ , \tag{2.8}$$

where ϕ denotes the absolute value, $\Phi = \phi \exp(i\theta)$. Using the Kähler potential Eq. (2.8) to expand the coefficients f, g and k in Eq. (2.4) in powers of ϕ and keeping only the lowest order terms, the scalar potential Eq. (2.4) becomes

$$V = V_0 + V_2 \phi^2 + V_6 \phi^6 + V_{10} \phi^{10} \ , \tag{2.9}$$

where

$$V_0 = e^{\hat{K}} |\hat{W}|^2 \Big(\hat{K}^m \hat{K}_m - 3 \Big) \ , \tag{2.10}$$

$$V_2 = e^{\hat{K}} |\hat{W}|^2 \hat{Z}_2 \Big(\hat{K}^m \hat{K}_m + \hat{K}^m \hat{K}^{\bar{n}} (\hat{Z}_2^{-2} \hat{Z}_{2m} \hat{Z}_{2\bar{n}} - \hat{Z}_2^{-1} \hat{Z}_{2m\bar{n}}) - 2 \Big) \ , \tag{2.11}$$

$$V_6 = e^{\hat{K}} |\hat{W}| |\hat{\lambda}_6| \cos(\xi - 6\theta) \Big| \frac{1}{3} \hat{K}^m \hat{K}_m - 2 \hat{Z}_2^{-1} \hat{K}^{\bar{m}} \hat{Z}_{2\bar{m}} + 1 \Big| \ , \tag{2.12}$$

$$V_{10} = e^{\hat{K}} |\hat{\lambda}_6|^2 \hat{Z}_2^{-1} \ , \tag{2.13}$$

the phase ξ in V_6 reads

$$\xi \equiv \arg \Big(\frac{1}{6} \hat{K}^m \hat{K}_m - \hat{Z}_2^{-1} \hat{K}^{\bar{m}} \hat{Z}_{2\bar{m}} + \frac{1}{2} \Big) + \arg(\hat{W}) - \arg(\hat{\lambda}_6) \tag{2.14}$$

and indices are raised and lowered with $\hat{K}^{M\bar{N}}$ and $\hat{K}_{M\bar{N}}$ respectively. Here V_0, V_6 and V_{10} result from the leading order expansion of f, g and k respectively, whereas V_2 is obtained by expanding f to next to leading order. The expansion is performed in this manner since the constant V_0, which would give rise to a cosmological constant, will be neglected henceforth[3]. Thus V_2 becomes the leading non-trivial term in the

[3] We assume the cosmological constant to be adjusted to the observationally required value either by terms arising from the hidden sector dependent superpotential, or by some other (yet unknown) mechanism.

expansion of f and Eq. (2.9) with V_0 removed then constitutes the leading order potential. In the following, the term leading order will be understood precisely in this sense, i.e. as leading non-trivial order.

By expanding f, g and k in Eq. (2.4) to next to leading order, one finds a first order correction $\Delta_1 V$ to the potential Eq. (2.9), at next to next to leading order one finds a second order correction $\Delta_2 V$, and so on. In carrying out this sort of expansion, we implicitly restrict our analysis to the values of ϕ for which all the terms in Eq. (2.4) are comparable, which will certainly be the case in the vicinity of the eventual saddle point. Using Eqs. (2.4) – (2.8), an order of magnitude approximation of the n-th order correction to the leading order potential is then given by

$$\Delta_n V \sim e^{\hat{K}} |\hat{W}|^2 \hat{Z}_2^{n+1} \phi^{2n+2} \ . \tag{2.15}$$

3. The saddle point condition

In this Section we consider the restrictions placed by the saddle point condition Eq. (1.3) on the leading order potential. The role of higher order corrections will be discussed in the next Section.

By choosing the phase θ of Φ such that $\cos(\xi - 6\theta) = -1$ in Eq. (2.12), the θ dependent part of the leading order potential Eq. (2.9) is minimized and we recover Eq. (1.1), where

$$m^2 = 2e^{\hat{K}} |\hat{W}|^2 \hat{Z}_2 \Big(\hat{K}^m \hat{K}_m + \hat{K}^m \hat{K}^{\bar{n}} (\hat{Z}_2^{-2} \hat{Z}_{2m} \hat{Z}_{2\bar{n}} - \hat{Z}_2^{-1} \hat{Z}_{2m\bar{n}}) - 2 \Big) \ , \tag{3.1}$$

$$A = e^{\hat{K}/2} |\hat{W}| \hat{Z}_2^{1/2} \Big| 2\hat{K}^m \hat{K}_m - 12 \hat{Z}_2^{-1} \hat{K}^{\bar{m}} \hat{Z}_{2\bar{m}} + 6 \Big| \ , \tag{3.2}$$

$$\lambda^2 = e^{\hat{K}} |\hat{\lambda}_6|^2 \hat{Z}_2^{-1} \ . \tag{3.3}$$

The saddle point condition Eq. (1.3) then becomes

$$|\hat{K}^m \hat{K}_m - 6 \hat{Z}_2^{-1} \hat{K}^{\bar{m}} \hat{Z}_{2\bar{m}} + 3|^2 = 20 (\hat{K}^m \hat{K}_m + \hat{K}^m \hat{K}^{\bar{n}} (\hat{Z}_2^{-2} \hat{Z}_{2m} \hat{Z}_{2\bar{n}} - \hat{Z}_2^{-1} \hat{Z}_{2m\bar{n}}) - 2) \ , \tag{3.4}$$

which is a partial differential equation for two unknown functions, $\hat{K}$ and $\hat{Z}_2$.

As a simple example we first consider a scenario in which there is only one hidden sector field h. Treating the functions $\hat{K}$ and $\hat{Z}_2$ as independent variables, Eq. (3.4) implies

$$\partial_h \hat{K} \partial_{\bar{h}} \hat{K} = -\beta \partial_h \partial_{\bar{h}} \hat{K} \ , \tag{3.5}$$

where β is a constant. An analogous equation appears in no-scale supergravity models [9] and is solved for

$$\hat{K} = \beta \log(h + h^*) \ . \tag{3.6}$$

Using this result and assuming $\hat{Z}_2 = \hat{Z}_2(h+h^*)$, the saddle point condition Eq. (3.4) in the one-dimensional case becomes

$$(3 - \beta + 6(h + h^*)\partial_h \log Z_2)^2 = 20(-\beta - 2 - (h + h^*)^2 \partial_h^2 \log \hat{Z}_2) \ , \tag{3.7}$$

whose general solution is

$$\hat{Z}_2 = (h + h^*)^{-2/9+\beta/6} \left[c_1 (h + h^*)^{\omega(\beta)} + c_2 (h + h^*)^{-\omega(\beta)} \right]^{5/9} \ , \tag{3.8}$$

where c_1, c_2 are constants, and $\omega(\beta) = 1/2\sqrt{-17 - 6\beta}$. The solution takes a particularly simple form if one of the constants c_1, c_2 is zero.

To find a solution of Eq. (3.4) in the general case with several hidden sector fields, we make an Ansatz motivated by the one-dimensional case and write

$$K = \sum_m \beta_m \log(h_m + h_m^*) + \kappa \prod_m (h_m + h_m^*)^{\alpha_m} \phi^2 + \mathcal{O}(\phi^4) \ , \tag{3.9}$$

where α_m, β_m and κ are constants. Kähler potentials of this type are found e.g. in Abelian orbifold compactifications of heterotic string theory [10] as well as in intersecting D-brane models [11]. In both cases the moduli fields play the role of the hidden sector fields h_m. Here we will, however, treat the parameters in Eq. (3.9) from a phenomenological point of view, without any particular string scenario in mind.

The Ansatz Eq. (3.9) solves Eq. (3.4) provided the parameters are related by

$$\alpha(36\alpha + 16 - 12\beta) + (\beta + 7)^2 = 0 \ , \tag{3.10}$$

where

$$\alpha = \sum_m \alpha_m \ , \tag{3.11}$$

$$\beta = \sum_m \beta_m \ . \tag{3.12}$$

In Table 1 we list solutions to Eq. (3.10) for which the soft susy breaking terms are nonzero and the α_m are rational numbers. In the string context $-\beta$ generically measures the number of hidden sector fields and therefore we restrict ourselves to the lowest values of β.

To summarize, for the leading order potential the saddle point condition Eq. (1.3) is satisfied identically with the class of Kähler potentials determined by Eq. (3.9) and the conditions on the parameters β_m, α_m as given in Table 1. While there definitely exist other solutions of Eq. (3.4) as well, Eq. (3.9) represents the only class of solutions for which $\hat{Z}_2$ is separable and the hidden sector dependence is of similar functional form for all the fields, provided that $K = K(h_m + h_m^*)$ and the hidden sector metric $\hat{K}_{m\bar{n}}$ is diagonal. Since we are not making any specific assumptions about the physical nature of the hidden sector fields, these seem to be quite natural conditions to impose on $\hat{Z}_2$.

Table 1: Values of β and α in the Kähler potential Eq. (3.9) for which the saddle point condition is satisfied identically.

$\beta = \sum_m \beta_m$	$\alpha = \sum_m \alpha_m$
-3	$-\frac{4}{9}$
-7	0
-7	$-\frac{25}{9}$
-11	$-\frac{1}{9}$
-11	-4

4. Higher order corrections

In the vicinity of the saddle point as given by Eq. (1.2), the slope of the leading order potential Eq. (1.1) is extremely small. Therefore, one may ask whether the corrections arising from the expansion of the potential Eq. (2.4) to higher orders will destroy this flatness. In this Section, we show that the required flatness [1, 2] is maintained if, in addition to the leading order potential, also the first and second order corrections satisfy certain conditions. Analogously to the leading order results, we find a form of the Kähler potential for which all these conditions are satisfied identically, i.e. irrespective of the vevs of hidden sector fields.

Within the slow roll approximation, the dynamics are determined by the first derivative of the potential. Therefore the condition for the n-th order correction, $\Delta_n V$, not to alter the leading order results can be expressed as

$$\Delta_n V'(\phi) \ll V'(\phi) \sim 10^{-3} \frac{m^2 \hat{Z}_2^2}{N(\phi)^2} \phi_0^5 \;, \tag{4.1}$$

where the derivative of the leading order potential, $V'(\phi)$, has been written in terms of the e-foldings $N(\phi)$ remaining until the end of inflation [1, 2]. Using Eq. (2.15), this condition becomes

$$\hat{Z}_2^{n-2} \phi_0^{2n-4} \ll 10^{-3} N(\phi)^{-2} \;, \tag{4.2}$$

which is satisfied automatically for $n > 2$ since[4] $\hat{Z}_2^{1/2} \phi_0 \ll 1$. This means that the third and higher order corrections are negligible and require no further attention. The corrections $\Delta_1 V$ and $\Delta_2 V$, on the other hand, can not be made small simply by adjusting parameters, but one needs to set their derivatives to zero identically. To be more precise, Eq. (4.2) is satisfied if $\Delta V_1'(\phi_0) = \Delta V_1''(\phi_0) = 0$ and $\Delta_2 V'(\phi_0) = 0$.

The first order corrections to the leading order potential can be written as

$$\Delta_1 V = V_4 \phi^4 + V_8 \phi^8 + V_{12} \phi^{12} \;, \tag{4.3}$$

[4] Note that it is actually the canonically normalized field $\phi_{\text{can}} \sim \hat{Z}_2^{1/2} \phi$ that is the MSSM inflaton. Therefore $\hat{Z}_2^{1/2} \phi \ll 1$. However, $V(\phi_{\text{can}})$ has a saddle point under the same conditions as $V(\phi)$ and the choice of the field variable plays no role in our analysis.

where the coefficients are obtained from Eqs. (2.4) – (2.7) by retaining only the next to leading order terms. With the leading order Kähler potential given by Eq. (3.9), the conditions $\Delta V_1'(\phi_0) = \Delta V_1''(\phi_0) = 0$ yield a pair of partial differential equations for the coefficient $\hat{Z}_4$ in the Kähler potential Eq. (2.8) whose only solution is

$$\hat{Z}_4 = \mu(\alpha, \beta, \gamma(\alpha, \beta))\hat{Z}_2^2 \ , \tag{4.4}$$

where $\gamma = \sum_m \alpha_m^2/\beta_m$ and we have assumed $\hat{Z}_4 = \hat{Z}_4(h_m + h_m^*)$. The parameters μ and γ are not freely selectable but completely determined by α and β such that $\Delta V_1'(\phi_0) = \Delta V_1''(\phi_0) = 0$. In Table 2 below we give their values for the choices of α and β considered in this work.

In a similar manner, the second order correction reads

$$\Delta_2 V = V_6\phi^6 + V_{10}\phi^{10} + V_{14}\phi^{14} \ , \tag{4.5}$$

with V_6, V_{10} and V_{14} here denoting the next to next to leading order part of Eqs. (2.4) – (2.7). In this case, the conditions to be placed on the Kähler potential are less stringent since one only needs to set $\Delta V_2'(\phi_0) = 0$. Assuming $\hat{K}, \hat{Z}_2, \hat{Z}_4$ in Eq. (2.8) to be given by Eqs. (3.9), (4.4), the condition $\Delta V_2'(\phi_0) = 0$ is satisfied for

$$\hat{Z}_6 = \nu(\alpha, \beta, \delta)\hat{Z}_2^3 \ , \tag{4.6}$$

where $\delta = \sum_m \alpha_m^3/\beta_m^2$, and only the relation between δ and ν is determined by α and β, see Table 2 below. Moreover, in addition to $\hat{Z}_6$ given by Eq. (4.6), the $\mathcal{O}(\phi^6)$ part of the Kähler potential may also contain solutions of the homogeneous part of the partial differential equation arising from $\Delta_2 V'(\phi_0) = 0$.

Thus we have shown that there exists a class of Kähler potentials for which the extreme flatness of the MSSM inflaton potential [1, 2] is generated and maintained also in the presence of higher order corrections irrespective of the hidden sector vevs. This class of Kähler potentials can be written in the form

$$\begin{aligned} K = & \sum_m \beta_m \log(h_m + h_m^*) + \kappa \prod_m (h_m + h_m^*)^{\alpha_m}\phi^2 + \mu\Big(\kappa \prod_m (h_m + h_m^*)^{\alpha_m}\Big)^2\phi^4 + \\ & \nu\Big(\kappa \prod_m (h_m + h_m^*)^{\alpha_m}\Big)^3\phi^6 + \mathcal{O}(\phi^8) \ , \end{aligned} \tag{4.7}$$

where the parameters are subject to the conditions given in Table 2 below.

Table 2: The coefficients of the higher order terms in the Kähler potential that guarantee the flatness of the MSSM inflaton potential.

$\beta = \sum_m \beta_m$	$\alpha = \sum_m \alpha_m$	$\gamma = \sum_m \alpha_m^2/\beta_m$	μ	$\delta = \sum_m \alpha_m^3/\beta_m^2$	ν
-3	$-\frac{4}{9}$	$\frac{1}{9}$	$-\frac{7}{36}$	$\frac{91}{324}$	ν
-7	0	0	$\frac{1}{12}$	δ	ν
-7	$-\frac{25}{9}$	$-\frac{10}{9}$	$\frac{2}{9}$	$-\frac{1654}{1863} + \frac{162}{23}\nu$	ν
-11	$-\frac{1}{9}$	$\frac{1}{21}$	$\frac{13}{126}$	$-\frac{8465}{75411} + \frac{162}{19}\nu$	ν
-11	-4	$-\frac{29}{21}$	$\frac{17}{84}$	$-\frac{2491}{2940} + \frac{36}{5}\nu$	ν

5. Conclusions

In this work we have considered the supergravity origin of the recently proposed MSSM inflationary model [1, 2]. In particular, we have shown that for the simple class of Kähler potentials given by Eq. (4.7), the extremely flat inflaton potential is produced identically in F-term supersymmetry breaking. The desired form of the potential is thus obtained for all hidden sector vevs and not just for some carefully chosen vacua.

The class of Kähler potentials Eq. (4.7) has a number of appealing features. Firstly, although it is necessary to fix the potential up to $\mathcal{O}(\phi^6)$, no new functions need to be introduced in addition to $\hat{K}$ and $\hat{Z}_2$ appearing already in the leading order expression, Eq. (3.9). Moreover, it is interesting to note that the form of Kähler potentials for which the MSSM inflationary scenario happens to be realized, is very common in string theory compactifications. As mentioned in Section 3, Kähler potentials of the form given in Eq. (3.9) arise e.g. in Abelian orbifold compactifications of the heterotic string theory [10] and in intersecting D-brane models [11]. In the heterotic case, the parameters α_m are modular weights, whereas in the intersecting brane models they depend on internal fluxes of the branes. To our knowledge, there is no specific compactification known so far, which would generate exactly the required values given in Table 1. This certainly would be a matter worth further investigation, especially keeping in mind that none of the string theory compactifications known to date produce exactly the actual MSSM. It would be very interesting to find a compactification that generates the saddle point along a d=6 flat direction of the MSSM.

The argumentation may also be turned the other way around. A supergravity model with F-term supersymmetry breaking, an MSSM like visible sector and a Kähler potential of the form given in Eq. (4.7) may naturally lead to an inflationary period driven by the MSSM degrees of freedom and with properties consistent with the observed cosmological data [12], provided the initial condition is such that the flat direction field finds itself in the vicinity of the saddle point. At this point, however, we wish to emphasize that since we are assuming the flat direction to be the only dynamical degree of freedom during inflation, we are also implicitly assuming the moduli fields to be stabilized before the beginning of inflation. Although the exceptionally low scale of inflation gives some justification for this assumption, the validity of it is highly model dependent and non-trivial, and should be discussed separately in the context of any realistic supergravity model. In any case, while inflation may still be possible even if the moduli fields are not stabilized, the resulting inflationary model will in general be different from the MSSM inflation discussed in this paper.

Finally, the supergravity model leading to the MSSM inflation can, at least in principle, be tested not only by cosmological observations but also in particle accelerators. For instance, given the Kähler potential Eq. (4.7), for a non-flat direction ψ with a renormalizable superpotential $W(\psi) = \frac{1}{3}\hat{\lambda}_3\psi^3$ one finds that the trilinear A-term is given by $A_3 = m_\psi\sqrt{2}(\alpha-\beta/3)\cos\xi/\sqrt{\alpha-\beta-2}$, where $\cos\xi$ contains the phase information. Once scaled down to LHC energies by the renormalization group equations, such relations have obvious ramifications for both sparticle phenomenology and inflation.

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