

## MSSM flat direction inflation

Rouzbah Allahverdi<sup>1,2</sup>

<sup>1</sup> *Perimeter Institute for Theoretical Physics, 31 Caroline Street North, Waterloo, ON, N2L 2Y5, Canada.*

<sup>2</sup> *Department of Physics and Astronomy, McMaster University, Hamilton, ON, L8S 4M1, Canada.*

We argue that all necessary ingredients for successful inflation are present in the flat directions of the minimal supersymmetric standard model (MSSM). Out of many gauge invariant combinations of the squarks, sleptons and Higgses, there are two flat directions, **LLe**, and **udd**, which are promising candidates for the inflaton. The model predicts more than  $10^3$  e-foldings with an inflationary scale of  $H_{\text{inf}} \sim \mathcal{O}(1-10)$  GeV, provides a tilted spectrum  $0.92 \leq n_s \leq 1$  with an amplitude of  $\delta_H \sim 10^{-5}$  and a negligible tensor perturbations. It yields a lower bound  $\sim 340$  GeV on the sparticle masses which is within the reach of LHC. There is no gravitino or moduli problem in this model.

### I. INTRODUCTION

The one crucial ingredient still missing in the otherwise highly successful theory of primordial inflation is its natural embedding within particle physics, particularly the standard model (SM) or its extensions. In almost all models of inflation the inflaton is treated as a SM gauge singlet and, sometimes a complete gauge singlet whose origin and couplings are chosen ad-hoc just to fit the observed cosmological data [1].

Recently we have constructed a model of inflation [2, 3] based on the flat directions of MSSM (for a review of MSSM flat directions, see [4]). In this model the inflaton is a gauge invariant combination of either squark or slepton fields. For a choice of the soft SUSY breaking parameters  $A$  and the inflaton mass  $m_\phi$ , the potential along the **udd** and **LLe** directions is such that there is a period of slow roll inflation of sufficient duration to provide the observed spectrum of CMB perturbations [21].

MSSM inflation occurs at a very low scale with  $H_{\text{inf}} \sim 1-10$  GeV and with field values much below the Planck scale. Hence it stands in strong contrast to the conventional inflation models which are based on ad hoc gauge singlet fields and often employ field values close to Planck scale (for a review, see [6]). In such models the inflaton couplings to SM physics are unknown. As a consequence, much of the post-inflationary evolution, such as reheating, thermalization, generation of baryon asymmetry and cold dark matter, which all depend very crucially on how the inflaton couples to the (MS)SM sector [7, 8, 9, 10], is not calculable from first principles. The great virtue of MSSM inflation based on flat directions is that the inflaton couplings to SM particles are known. More importantly, the inflaton mass is related either to squark or slepton masses, it could be measured by LHC or a future Linear Collider.

### II. THE MODEL

Let us recapitulate the main features of MSSM flat direction inflation [2, 3]. As is well known, in the limit

of unbroken SUSY the flat directions have exactly vanishing potential. This situation changes if we take into account soft SUSY breaking and non-renormalizable superpotential terms of the type [4]

$$W_{\text{non}} = \sum_{n>3} \frac{\lambda_n}{n} \frac{\Phi^n}{M^{n-3}}, \quad (1)$$

where  $\Phi$  is a superfield which contains the flat direction. Within MSSM all the flat directions are lifted by non-renormalizable operators with  $4 \leq n \leq 9$  [11]. We expect that quantum gravity effects yield  $M = M_{\text{P}} = 2.4 \times 10^{18}$  GeV and  $\lambda_n \sim \mathcal{O}(1)$  [12].

Let us focus on the lowest order superpotential term in Eq. (1) which lifts the flat direction. Soft SUSY breaking induces a mass term for  $\phi$  and an  $A$ -term so that the scalar potential along the flat direction reads

$$V = \frac{1}{2} m_\phi^2 \phi^2 + A \cos(n\theta + \theta_A) \frac{\lambda_n \phi^n}{n M_{\text{P}}^{n-3}} + \lambda_n^2 \frac{\phi^{2(n-1)}}{M_{\text{P}}^{2(n-3)}}, \quad (2)$$

Here  $\phi$  and  $\theta$  denote respectively the radial and the angular coordinates of the complex scalar field  $\Phi = \phi \exp[i\theta]$ , while  $\theta_A$  is the phase of the  $A$ -term (thus  $A$  is a positive quantity with dimension of mass). Note that the first and third terms in Eq. (2) are positive definite, while the  $A$ -term leads to a negative contribution along the directions whenever  $\cos(n\theta + \theta_A) < 0$ . [22]

The maximum impact from the  $A$ -term is obtained when  $\cos(n\theta + \theta_A) = -1$  (which occurs for  $n$  values of  $\theta$ ).

In the gravity mediated SUSY breaking case, the  $A$ -term and the soft SUSY breaking mass terms are expected to be the same order of magnitude as the gravitino mass, i.e.  $m_\phi \sim A \sim m_{3/2} \sim \mathcal{O}(1)$  TeV. Now if  $A$  and  $m_\phi$  are related by

$$A^2 = 8(n-1)m_\phi^2, \quad (3)$$

then both the first and second derivatives of  $V$  vanish at  $\phi_0$ , i.e.  $V'(\phi_0) = 0$ ,  $V''(\phi_0) = 0$ , where

$$\phi_0 = \left( \frac{m_\phi M_{\text{P}}^{n-3}}{\lambda_n \sqrt{2n-2}} \right)^{1/(n-2)}. \quad (4)$$

The potential near the saddle point Eq. (4) is very flat along the *real direction* but not along the *imaginary direction*. Along the imaginary direction the curvature is determined by  $m_\phi$ . Around  $\phi_0$  the field lies in a plateau with a potential energy

$$V(\phi_0) = \frac{(n-2)^2}{2n(n-1)} m_\phi^2 \phi_0^2. \quad (5)$$

As the result, if initially  $\phi \sim \phi_0$ , a slow roll phase of inflation is driven by the third derivative of the potential. The Hubble expansion rate during inflation which is given by

$$H_{\text{inf}} = \frac{(n-2)}{\sqrt{6n(n-1)}} \frac{m_\phi \phi_0}{M_{\text{P}}}. \quad (6)$$

When  $\phi$  is very close to  $\phi_0$ , the first derivative is extremely small. The field is effectively in a de Sitter background, and we are in self-reproduction (or eternal inflation) regime where the two point correlation function for the flat direction fluctuation grows with time. But eventually classical friction wins and slow roll begins at  $\phi \approx \phi_{\text{self}}$  [2]

$$(\phi_0 - \phi_{\text{self}}) \simeq \left( \frac{m_\phi \phi_0^2}{M_{\text{P}}^3} \right)^{1/2} \phi_0. \quad (7)$$

The slow roll potential in this case reads

$$V(\phi) = V(\phi_0) + \frac{1}{3!} V'''(\phi_0) (\phi - \phi_0)^3 + \dots, \quad (8)$$

$$V'''(\phi_0) = 2(n-2)^2 \frac{m_\phi^2}{\phi_0}.$$

We can now solve the equation of motion for the  $\phi$  field in the slow-roll approximation,

$$3H\dot{\phi} = -\frac{1}{2} V'''(\phi_0) (\phi - \phi_0)^2, \quad (9)$$

assuming initial conditions such that the flat direction starts in the vicinity of  $\phi_0$  with  $\dot{\phi} \approx 0$ . Inflation ends when either of the slow roll parameters,  $\epsilon \equiv (M_{\text{P}}^2/2)(V'/V)^2$  and  $\eta \equiv M_{\text{P}}^2(V''/V)$ , becomes of  $\mathcal{O}(1)$ . It happens that  $|\eta| \sim 1$  when  $\phi \approx \phi_{\text{end}}$ , where

$$(\phi_0 - \phi_{\text{end}}) \sim \frac{\phi_0^3}{4n(n-1)M_{\text{P}}^2}. \quad (10)$$

The number of e-foldings during the slow roll from  $\phi$  to  $\phi_{\text{end}}$  is given by

$$\mathcal{N}_e(\phi) = \int_\phi^{\phi_{\text{end}}} \frac{H_{\text{inf}} d\phi}{\dot{\phi}} \simeq \frac{\phi_0^3}{2n(n-1)M_{\text{P}}^2(\phi_0 - \phi)}, \quad (11)$$

where we have used  $V'(\phi) \sim (\phi - \phi_0)^2 V'''(\phi_0)$  (this is justified since  $V'(\phi_0) \sim 0$ ,  $V''(\phi_0) \sim 0$ ), and Eq. (9).

The total number of e-foldings in the slow roll regime is then found from Eq.(7)

$$\mathcal{N}_{\text{tot}} \simeq \frac{1}{2n(n-1)} \left( \frac{\phi_0^2}{m_\phi M_{\text{P}}} \right)^{1/2}. \quad (12)$$

The observationally relevant perturbations are generated when  $\phi \approx \phi_{\text{COBE}}$ . The number of e-foldings between  $\phi_{\text{COBE}}$  and  $\phi_{\text{end}}$ , denoted by  $\mathcal{N}_{\text{COBE}}$  follows from Eq. (11)

$$\mathcal{N}_{\text{COBE}} \simeq \frac{\phi_0^3}{2n(n-1)M_{\text{P}}^2(\phi_0 - \phi_{\text{COBE}})}. \quad (13)$$

The amplitude of perturbations thus produced is given by

$$\delta_H \equiv \frac{1}{5\pi} \frac{H_{\text{inf}}^2}{\dot{\phi}} \simeq \frac{1}{5\pi} \sqrt{\frac{2}{3}} n(n-1)(n-2) \left( \frac{m_\phi M_{\text{P}}}{\phi_0^2} \right) \mathcal{N}_{\text{COBE}}^2, \quad (14)$$

where we have used Eqs.(6,8,13). Again after using these equations, the spectral tilt of the power spectrum and its running are found to be

$$n_s = 1 + 2\eta - 6\epsilon \simeq 1 - \frac{4}{\mathcal{N}_{\text{COBE}}}, \quad (15)$$

$$\frac{dn_s}{d \ln k} = -\frac{4}{\mathcal{N}_{\text{COBE}}^2}.$$

### III. PROPERTIES AND PREDICTIONS

As discussed in [2], among the about 300 flat directions there are two that can lead to a successful inflation along the lines discussed above.

One is **udd** which, up to an overall phase factor, is parameterized by

$$u_i^\alpha = \frac{1}{\sqrt{3}} \phi, \quad d_j^\beta = \frac{1}{\sqrt{3}} \phi, \quad d_k^\gamma = \frac{1}{\sqrt{3}} \phi. \quad (16)$$

Here  $1 \leq \alpha, \beta, \gamma \leq 3$  are color indices, and  $1 \leq i, j, k \leq 3$  denote the quark families. The flatness constraints require that  $\alpha \neq \beta \neq \gamma$  and  $j \neq k$ .

The other direction is **LLe**, parameterized by (again up to an overall phase factor)

$$L_i^a = \frac{1}{\sqrt{3}} \begin{pmatrix} 0 \\ \phi \end{pmatrix}, \quad L_j^b = \frac{1}{\sqrt{3}} \begin{pmatrix} \phi \\ 0 \end{pmatrix}, \quad e_k = \frac{1}{\sqrt{3}} \phi, \quad (17)$$

where  $1 \leq a, b \leq 2$  are the weak isospin indices and  $1 \leq i, j, k \leq 3$  denote the lepton families. The flatness constraints require that  $a \neq b$  and  $i \neq j \neq k$ . Both these flat directions are lifted by  $n = 6$  non-renormalizable operators,

$$W_6 \supset \frac{1}{M_{\text{P}}^3} (LLe)(LLe), \quad W_6 \supset \frac{1}{M_{\text{P}}^3} (udd)(udd) \quad (18)$$

The reason for choosing either of these two flat directions is twofold: (i) a non-trivial  $A$ -term arises, at the lowest order, only at  $n = 6$ ; and (ii) we wish to obtain the correct COBE normalization of the CMB spectrum [for details, see [3]].

According to the arguments presented above, successful MSSM flat direction inflation has the following model parameters:

$$m_\phi \sim 1 - 10 \text{ TeV}, \quad n = 6, \quad A = \sqrt{40}m_\phi, \quad \lambda \sim \mathcal{O}(1). \quad (19)$$

Here we assume that  $\lambda$  (we drop the subscript "6") is of order one, which is the most natural assumption when  $M = M_P$ .

The Hubble expansion rate during inflation and the VEV of the saddle point are

$$H_{\text{inf}} \sim 1 - 10 \text{ GeV}, \quad \phi_0 \sim (1 - 3) \times 10^{14} \text{ GeV}. \quad (20)$$

Note that both the scales are sub-Planckian. The total energy density stored in the inflaton potential is  $V_0 \sim 10^{36} - 10^{38} \text{ GeV}^4$ . The fact that  $\phi_0$  is sub-Planckian guarantees that the inflationary potential is free from the uncertainties about physics at super-Planckian VEVs. The total number of e-foldings during the slow roll evolution is large enough to dilute any dangerous relic away, see Eq. (12):

$$\mathcal{N}_{\text{tot}} \sim 10^3, \quad (21)$$

Domains which are initially closer than  $\phi_{\text{self}}$  to  $\phi_0$ , see Eq. (7), can enter self-reproduction in eternal inflation, with no observable consequences.

At such low scales as in MSSM inflation the number of e-foldings,  $\mathcal{N}_{\text{COBE}}$ , required for the observationally relevant perturbations, is much less than 60 [14]. If the inflaton decays immediately after the end of inflation, we obtain  $\mathcal{N}_{\text{COBE}} \sim 50$ . Despite the low scale, the flat direction can generate adequate density perturbations as required to explain the COBE normalization. This is due to the extreme flatness of the potential (recall that  $V' = 0$ ), which causes the velocity of the rolling flat direction to be extremely small. From Eq. (14) we find an amplitude of

$$\delta_H \simeq 1.91 \times 10^{-5}. \quad (22)$$

There is a constraint on the mass of the flat direction from the amplitude of the CMB anisotropy:

$$m_\phi = (340 \text{ GeV}) \times \lambda^{-1} \left( \frac{\mathcal{N}_{\text{COBE}}}{50} \right)^{-4}. \quad (23)$$

We get a lower limit on the mass parameter when  $\lambda \leq 1$ . For smaller values of  $\lambda \ll 1$ , the mass of the flat direction must be larger. Note that the above bound on the inflaton mass arises at high scales, i.e.  $\phi = \phi_0$ . However, through renormalization group flow, it is connected to the low scale mass [3].

The spectral tilt of the power spectrum is not negligible because, although the first slow roll parameter is  $\epsilon \sim 1/\mathcal{N}_{\text{COBE}}^4 \ll 1$ , the other slow roll parameter is given by  $\eta = -2/\mathcal{N}_{\text{COBE}}$  and thus, see Eq. (15)[23]

$$n_s \sim 0.92, \quad (24)$$

$$\frac{dn_s}{d \ln k} \sim -0.002, \quad (25)$$

where we have taken  $\mathcal{N}_{\text{COBE}} \sim 50$  (which is the maximum value allowed for the scale of inflation in our model). In the absence of tensor modes, this agrees with the current WMAP 3-years' data within  $2\sigma$  [1]. Note that MSSM inflation does not produce any large stochastic gravitational wave background during inflation. Gravity waves depend on the Hubble expansion rate, and in our case the energy density stored in MSSM inflation is very small.

#### IV. SENSITIVITY OF MSSM INFLATION

The dynamics of the flat direction inflaton has been discussed assuming the saddle point condition Eq. (3) is satisfied exactly. The question then is, how large a deviation can be allowed for before slow roll inflation will be spoiled. To facilitate the discussion, let us define

$$\delta \equiv \frac{A^2}{40m_\phi^2} \equiv 1 \pm 4\alpha^2, \quad (26)$$

where  $\alpha \ll 1$ . There are two distinct possibilities: either  $\delta > 1$  or  $\delta < 1$  (when  $\delta = 1$  we recover the saddle point condition). In the former case there is a barrier which separates the global minimum  $\phi = 0$  and the false minimum at  $\phi \simeq \phi_0$ . If the barrier is too high, the field remains trapped and there is no slow-roll inflation. Therefore one must require that tunneling is effective so that the field can jump to the top of the barrier. In the latter case there is no minimum but the potential may be too steep for slow roll inflation.

Obtaining sufficient inflation, i.e. such that the number of e-foldings is  $\geq \mathcal{N}_{\text{COBE}}$ , with observationally acceptable properties results in a following constraint on  $\alpha$ ,

$$\alpha \leq \frac{1}{30\mathcal{N}_{\text{COBE}}} \left( \frac{\phi_0}{M_P} \right)^2, \quad (27)$$

in both cases with  $\delta > 1$  and  $\delta < 1$  [3]. For typical values of  $\phi_0 \lesssim 10^{15} \text{ GeV}$ , the saddle point condition in Eq. (3) requires tuning at a level;

$$\alpha \sim 10^{-9}. \quad (28)$$

Since radiative corrections modify  $\alpha$ , we need to finetune the potential to a few (but not all) orders

in perturbation theory. Although not disastrous, this can hardly be considered a satisfactory situation. However, it is conceivable that the mechanism of supersymmetry breaking could remove the fine-tuning in some natural, dynamical way. For instance,  $A/m$  could turn out to be a renormalization group (RG) fixed point so that once the ratio is fixed, it would remain fixed under quantum corrections (for example, see [17]). This is an interesting possibility which requires a detailed investigation.

## V. THE INFLATON AND LHC

Let us recall that the constraint on the mass of the  $n = 6$  flat direction inflaton in Eq. (23). As mentioned earlier, this is the bound on the mass of the flat direction during inflation, determined at the scale  $\phi = \phi_0$ . Since the inflaton mass runs from  $\phi_0$  down to the LHC energy scales, it will also get scaled.

For **LLe** the flat direction mass only gets larger due to the gaugino loops. As a result, the lower bound on the slepton masses is larger by  $\leq 50\%$  compared to that in Eq. (23) [3]). The situation would be similar for **udd** without the top squark. For the **u<sub>3</sub>dd** direction it is possible that the inflaton mass gets even smaller at the weak scale. We do not claim that LHC can discover MSSM inflation, but it can certainly rule out the possibility. If LHC does not find low energy supersymmetry within  $\sim \text{TeV}$ , then MSSM inflation is effectively ruled out.

Unlike  $m_\phi$ , there is no prospect of measuring the  $A$  term, because it is related to the non-renormalizable interactions which are suppressed by  $M_P$ . However, a knowledge of supersymmetry breaking sector and its communication with the observable sector may help to link the non-renormalizable  $A$ -term under consideration to the renormalizable ones.

To elucidate this, let us consider the Polonyi model where a general  $A$ -term at a tree level is given by

$$m_{3/2}[(a-3)W + \phi(dW/d\phi)],$$

with  $a = 3 - \sqrt{3}$  [16]. One then finds a relationship between  $A$ -terms corresponding to  $n = 6$  and  $n = 3$  superpotential terms, denoted by  $A_6$  and  $A_3$  respectively, at high scales:

$$A_6 = \frac{3 - \sqrt{3}}{6 - \sqrt{3}} A_3. \quad (29)$$

One can then use relevant RG equations to relate  $A_6$  which is relevant for inflation, to  $A_3$  at the weak scale, which can be constrained and/or measured. In principle this can also be done in general, provided that we have sufficient information about the supersymmetry breaking sector and its communication with the MSSM sector.

## VI. END OF MSSM INFLATION

After the end of inflation, the flat direction starts rolling towards its global minimum. At this stage the dominant term in the scalar potential will be:  $m_\phi \phi^2/2$ . Since the frequency of oscillations is  $\omega \sim m_\phi \sim 10^3 H_{\text{inf}}$ , the flat direction oscillates a large number of times within the first Hubble time after the end of inflation. Hence the effect of expansion is negligible.

We recall that the curvature of the potential along the angular direction is much larger than  $H_{\text{inf}}^2$ . Therefore, the flat direction has settled at one of the minima along the angular direction during inflation from which it cannot be displaced by quantum fluctuations. This implies that no torque will be exerted, and hence the flat direction motion will be one dimensional, i.e. along the radial direction.

Flat direction oscillations excite those MSSM degrees of freedom which are coupled to it. The inflaton, either **LLe** or **udd** flat direction, is a linear combination of slepton or squark fields. Therefore it has gauge couplings to the gauge/gaugino fields and Yukawa couplings to the Higgs/Higgsino fields. Let us elucidate the physics, by considering the case when **LLe** flat direction is the inflaton.

An efficient bout of particle creation occurs when the inflaton crosses the origin, which happens twice in every oscillation. The reason is that fields which are coupled to the inflaton are massless near the point of enhanced symmetry. Mainly electroweak gauge fields and gauginos are then created as they have the largest coupling to the flat direction. The production takes place in a short interval,  $\Delta t \sim (gm_\phi \phi_0)^{-1/2}$ , where  $\phi_0 \sim 10^{14}$  GeV is the initial amplitude of the inflaton oscillation, during which quanta with a physical momentum  $k \lesssim (gm_\phi \phi_0)^{1/2}$  are produced. As the inflaton VEV is rolling back to its maximum value  $\phi_0$ , the mass of the produced quanta  $g\langle\phi(t)\rangle$  increases. The gauge and gaugino fields can (perturbatively) decay to the fields which are not coupled to the inflaton, for instance to (s)quarks. Note that (s)quarks are not coupled to the flat direction, hence they remain massless throughout the oscillations.

The decay is very quick compared with the frequency of inflaton oscillations. The ratio of energy density in relativistic particles thus produced  $\rho_{\text{rel}}$  with respect to the total energy density  $\rho_0$  is

$$\frac{\rho_{\text{rel}}}{\rho_0} \sim 10^{-2} g. \quad (30)$$

This implies that a fraction  $\sim \mathcal{O}(10^{-2})$  of the inflaton energy density is transferred into relativistic (s)quarks every time that the inflaton passes through the origin. This is so-called instant preheating mechanism [18]. It is quite an efficient mechanism in our model as it can convert almost all of the energy density in the inflaton into radiation within a Hubble time

The maximum temperature attained by the plasma would be given by:

$$T_{max} \sim (m_\phi \phi_0)^{1/2} \geq 10^9 \text{ GeV}. \quad (31)$$

This temperature may be too high and could lead to thermal overproduction of gravitinos [19, 20]. However the dominant source of gravitino production in a thermal bath is scatterings which include an on-shell gluon or gluino leg.

In order to suppress thermal gravitino production it is therefore sufficient to make gluon and gluino fields heavy enough such that they are not kinematically accessible to the reheated plasma. This suggests a natural solution to the thermal gravitino problem in the case of our model. Consider another flat direction with a non-zero VEV, denoted by  $\varphi$ , which spontaneously breaks the  $SU(3)_C$  group. For example, if **LLe** is the inflaton, then **udd** provides a unique candidate which can simultaneously develop VEV.

So long as  $g\varphi \gg T$ , the gluon/gluino fields will be too heavy and not kinematically accessible to the reheated plasma. Once  $g\langle\varphi\rangle \simeq T$ , gluon/gluino fields come into equilibrium with the thermal bath. As pointed out in Refs. [7, 8], if the initial VEV of **udd** is  $\varphi_0 > 10^{10}$  GeV, then the temperature at which gluon/gluino become kinematically accessible, i.e.  $g\langle\varphi\rangle \simeq T$ , is given by [8]

$$T_R \leq 10^7 \text{ GeV}. \quad (32)$$

This is the final reheat temperature at which gluons and gluinos are all in thermal equilibrium with the other degrees of freedom. The standard calculation of thermal gravitino production via scatterings can then be used for  $T \leq T_R$ . Note however that  $T_R$  is sufficiently low to avoid thermal overproduction of gravitinos.

Finally, we also make a comment on the cosmological moduli problem. The moduli are generically displaced from their true minimum if their mass is less than the expansion rate during inflation. The moduli obtain a mass  $\sim \mathcal{O}(\text{TeV})$  from supersymmetry breaking. They start oscillating with a large amplitude, possibly as big as  $M_P$ , when the Hubble parameter drops below their mass. Since moduli are only gravitationally coupled to other fields, their oscillations dominate the Universe while they decay very late. The resulting reheat temperature is below MeV, and is too low to yield a successful primordial nucleosynthesis.

However, in our case  $H_{\text{inf}} \ll \text{TeV}$ . This implies that quantum fluctuations cannot displace the moduli from their true minima during the inflationary epoch driven by MSSM flat directions. Moreover, any oscillations of the moduli will be exponentially damped during the inflationary epoch. Therefore our model is free from the infamous moduli problem.

## VII. CONCLUSION

The existence of a saddle point in the scalar potential of the **udd** or **LLe** MSSM flat directions appears, perhaps surprisingly, to provide all the necessary ingredients for an observationally realistic model of inflation. The exceptional feature of the model, which sets it apart from conventional singlet field inflation models, is the fact that here the inflaton is a gauge invariant combination of the squark or slepton fields. As a consequence, the mass of the inflaton is not a free parameter but is related to the masses of e.g. sleptons, should the **LLe** direction be the inflaton. Hence LHC can indeed put a constraint on the model: it may not be able to verify it, but it certainly can rule it out. Moreover, the couplings of the inflaton to the MSSM matter and gauge fields are known. This makes it possible to address the questions of reheating in an unambiguous way. As we saw, the model is free from the gravitino and moduli problems.

To summarize, MSSM flat direction inflation is unique in being both a successful model of inflation and at the same time having a concrete and real connection to physics that can be observed in earth bound laboratories.

## Acknowledgments

I am indebted to my collaborators Kari Enqvist, Juan Garcia-Bellido, Asko Jokinen, Alexander Kusenko and Anupam Mazumdar. I also thank the organizers of IPM-LHP06 for their invitation and kind hospitality. This work was supported by the National Sciences and Engineering Research Council of Canada (NSERC) and by Perimeter Institute for Theoretical Physics. Research at Perimeter Institute is supported in part by the province of Ontario through MEDT.

[1] D.N. Spergel, et.al., astro-ph/0603449.

[2] R. Allahverdi, K. Enqvist, J. Garcia-Bellido and A. Mazumdar, Phys. Rev. Lett. (to appear), arXiv:hep-ph/0605035.

[3] R. Allahverdi, K. Enqvist, J. Garcia-Bellido, A. Joki-

nen and A. Mazumdar, arXiv:hep-ph/0610134.

[4] K. Enqvist and A. Mazumdar, Phys. Rept. **380**, 99 (2003) [arXiv:hep-ph/0209244]; M. Dine and A. Kusenko, Rev. Mod. Phys. **76**, 1 (2004) [arXiv:hep-ph/0303065].

- [5] R. Allahverdi, A. Kusenko and A. Mazumdar, arXiv:hep-ph/0608138.
- [6] D. H. Lyth and A. Riotto, Phys. Rept. **314**, 1 (1999) [arXiv:hep-ph/9807278].
- [7] R. Allahverdi and A. Mazumdar, arXiv:hep-ph/0505050.
- [8] R. Allahverdi and A. Mazumdar, JCAP **10**, 008 (2006) [arXiv:hep-ph/0512227].
- [9] R. Allahverdi and A. Mazumdar, arXiv:hep-ph/0603244.
- [10] R. Allahverdi and A. Mazumdar, arXiv:hep-ph/0608296.
- [11] T. Gherghetta, C. F. Kolda and S. P. Martin, Nucl. Phys. B **468**, 37 (1996) [arXiv:hep-ph/9510370].
- [12] M. Dine, L. Randall and S. Thomas, Phys. Rev. Lett. **75**, 398 (1995) [arXiv:hep-ph/9503303]; M. Dine, L. Randall and S. Thomas, Nucl. Phys. B **458**, 291 (1996) [arXiv:hep-ph/9507453].
- [13] R. Allahverdi, K. Enqvist, A. Jokinen and A. Mazumdar, JCAP **10** 007 (2006) [arXiv:hep-ph/0603255].
- [14] C. P. Burgess, R. Easther, A. Mazumdar, D. F. Mota and T. Multamaki, JHEP **0505**, 067 (2005) [arXiv:hep-th/0501125].
- [15] R. Allahverdi and A. Mazumdar, arXiv:hep-ph/0610069.
- [16] H. P. Nilles, Phys. Rept. **110**, 1 (1984).
- [17] M. Lanzagorta and G. G. Ross, Phys. Lett. B **349**, 319 (1995) [arXiv:hep-ph/9501394]; M. Lanzagorta and G. G. Ross, Phys. Lett. B **364**, 163 (1995) [arXiv:hep-ph/9507366].
- [18] G. N. Felder, L. Kofman and A. D. Linde, Phys. Rev. D **59**, 123523 (1999) [arXiv:hep-ph/9812289].
- [19] J. R. Ellis, J. E. Kim and D. V. Nanopoulos, Phys. Lett. B **145**, 181 (1984).
- [20] M. Bolz, A. Brandenburg and W. Buchmüller, Nucl. Phys. B **606**, 518 (2001) [arXiv:hep-ph/0012052].
- [21] In a recent similar model with small Dirac neutrino masses, the observed microwave background anisotropy and the tilted power spectrum are related to the neutrino masses [5]. The model relies solely on renormalizable couplings.
- [22] The importance of the A-term has also been highlighted in a successful MSSM curvaton model [13].
- [23] Obtaining  $n_s > 0.92$  requires deviation from the saddle point condition in Eq. (3). Recently we have illustrated that within MSSM inflation the spectral tilt varies from  $0.92 \leq n_s \leq 1$  [15]. Inflation occurring exactly at the saddle point yields the lower limit, while the upper limit is obtained when the total number of e-foldings is saturated by the number of e-foldings required for the COBE normalization, i.e.  $\mathcal{N}_{\text{tot}} \approx \mathcal{N}_{\text{COBE}} \sim 50$ .