Inflation and The Minimal Supersymmetric Standard Model

Rouzbeh Allahverdi
Department of Physics & Astronomy, University of New Mexico, Albuquerque, NM 87131, USA

There is strong evidence from cosmological data that the universe underwent an epoch of superluminal expansion called inflation. A satisfactory embedding of inflation in fundamental physics has been an outstanding problem at the interface of cosmology and high energy physics. We show how inflation can be realized within the Minimal Supersymmetric Standard Model (MSSM). The inflaton candidates are two specific combinations of supersymmetric partners of quarks and leptons. MSSM inflation occurs at a low scale and generates perturbations in the range experimentally allowed by the latest data from Wilkinson Microwave Anisotropy Probe (WMAP). The parameter space for inflation is compatible with supersymmetric dark matter, and the Large Hadron Collider (LHC) is capable of discovering the inflaton candidates in the allowed regions of parameter space.

I. INTRODUCTION

Inflation is the dominant paradigm of the early universe cosmology. It solves the flatness and isotropy problems of the hot big bang cosmology and generates the seeds for structure formation [1]. To date, experiments confirm simplest predictions of inflation: a flat universe and nearly scale-invariant adiabatic fluctuations with a gaussian spectrum, which are imprinted in temperature anisotropy of the Cosmic Microwave Background (CMB) [2]. In spite of its impressive success, a satisfactory explanation of the microscopic origin of inflation has been lacking over the past 27 years. In almost all proposed models [3] the inflaton is added as an absolute gauge singlet which has no natural place in particle physics 1. The inflaton mass and its couplings to other fields are not generally tied to any fundamental theory and instead are set by hand in order to match observations. In particular, since the inflaton couplings to the Standard Model (SM) fields are arbitrary, it is not clear how most of the inflaton energy eventually goes into the observable sector, as required by the primordial abundance of light elements made during Big Bang Nucleosynthesis (BBN) [7].

The Minimal Supersymmetric Standard Model (MSSM) is a well motivated extension of the SM [8]. It introduces scalar partners for the quarks and leptons (called squarks and sleptons respectively), and fermionic partners for the gauge and Higgs fields (called gauginos and Higgsinos respectively). The supersymmetric partners have the same quantum numbers as ordinary fields, and their mass is supposed to be in the $\sim 100 \text{ GeV} - 1 \text{ TeV}$ range in order to address the hierarchy problem of the SM. There exist a large number of flat directions in the field space, consisting of the squark and slepton fields, along which the scalar potential identically vanishes in the limit of exact supersymmetry [4]. These flat directions have important cosmological consequences [10, 11]. In particular, it has been recently shown that two specific flat directions can lead to a successful inflation [12, 13]. This is the first example of the inflaton finding a natural place in a well motivated extension of the SM. In particular, since the inflaton is related to the squark and slepton fields, the inflation sector can be probed in colliders, most notably the Large Hadron Collider (LHC).

The aim of this article is to present a brief review of MSSM inflation and its various achievements. In Section 2 we show how inflation can happen along MSSM flat directions. We then discuss properties of MSSM inflation and its predictions in light of the 5-year Wilkinson Microwave Anisotropy (WMAP) data in Section 3. Next we turn to the compatibility between MSSM inflation and supersymmetric dark matter and present a combined analysis of parameter space in the case of minimal supergravity in Section 4. In Section 5 we briefly discuss the reheating of the universe after MSSM inflation. Section 6 contains conclusions and some discussions.

II. INFLATION IN MSSM

Flat directions in the scalar potential of MSSM are classified by gauge-invariant monomials made of the squark, slepton and Higgs fields [9]. In the limit of unbroken supersymmetry, the potential identically vanishes along these directions. Supersymmetry breaking terms and non-renormalizable superpotential terms lift the flat directions [14, 15].

1 For the only exceptions known to the author, see Refs. 4, 5, 6.
Denoting the flat direction by $\Phi$, the superpotential terms have the form

$$W \supset \lambda \Phi^n n M^{n-3},$$

where $n > 3$ and $M$ is the scale of the new physics that induces these terms. Planck-suppressed non-renormalizable terms are generically expected to be induced by quantum gravity or string theory, for which $M = M_P \equiv 2.4 \times 10^{18}$ GeV and $\lambda_n \sim O(1)$. This is the case that we consider henceforth. Within the MSSM all flat directions are lifted by non-renormalizable operators with $4 \leq n \leq 9 \ [9]$. The scalar potential along the flat direction reads $\ [12]$

$$V = \frac{1}{2} m_\phi^2 \phi^2 + A \cos(n\theta + \theta_A) \frac{\lambda \phi^n}{n M_P^{n-3}} + \lambda^2 \phi^{2(n-1)} M_P^{2(n-3)} ,$$

where the first and last terms on the right-hand side are the soft mass term and the $A$-term respectively. Here $\phi$ and $\theta$ denote 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). For weak scale supersymmetry and gravity mediation of supersymmetry breaking to the observable sector we have $m_\phi \sim A \sim O$(TeV). After minimizing the potential along the angular direction we find

$$V(\phi) = \frac{1}{2} m_\phi^2 \phi^2 - A \frac{\lambda \phi^n}{n M_P^{n-3}} + \lambda^2 \phi^{2(n-1)} M_P^{2(n-3)}.$$

If

$$\frac{A^2}{8(n-1)m_\phi^2} \equiv 1 + \left(\frac{n-2}{2}\right)^2 \alpha^2 ,$$

where $|\alpha^2| \ll 1$, the potential has a point of inflection at $\ [2]$

$$\phi_0 = \left(\frac{m_\phi M_P^{n-3}}{\lambda \sqrt{2n-2}}\right)^{1/(n-2)},$$

at which $\ [13]$

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

$$V'(\phi_0) = \left(\frac{n-2}{2}\right)^2 \alpha^2 m_\phi^2 \phi_0 ,$$

$$V''(\phi_0) = 0 ,$$

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

We have kept terms that are leading order in $\alpha^2$.

There is a plateau in the vicinity of the point of inflection $\ [12]$, $\ [13]$,

$$|\phi - \phi_0| \sim \frac{\phi_0^3}{2(n(n-1)) M_P^2} ,$$

within which the slow roll parameters $\epsilon \equiv (M_P^2/2)(V'/V)^2$ and $\eta \equiv M_P^2(V''/V)$ are smaller than 1. If the field lies in this plateau and has a sufficiently small velocity, inflation occurs. The Hubble expansion rate during inflation is

---

2 The parameters $m_\phi$, $A$, $\lambda$ are all affected by radiative corrections. However they do not remove the point of inflection nor shift it to unreasonable values. All that happens is that the condition to have a point of inflection $\ [13]$ and the Vacuum Expectation Value (VEV) of the inflection point $\ [13]$ will be slightly modified $\ [13]$. Desirable initial conditions for inflation can be naturally set within prior phase(s) of false vacuum inflation either by the help of quantum fluctuations $\ [13]$, or as a result of attractor behavior of the inflection point $\ [13]$. Many models of high energy physics possess metastable vacua, and hence can lead to false vacuum inflation at some stage during the evolution of the early universe.
given by
\[ H_{\text{inf}} \simeq \frac{n-2}{\sqrt{6n(n-1)}} \frac{m_\phi \phi_0}{M_P}, \]  
(11)
and the total number of e-foldings in the slow-roll regime (where \(|\epsilon|, |\eta| < 1\)) can be as large as \(10^3\), which is more than enough to solve the flatness and isotropy problems.

The number of e-foldings between the time that observationally relevant perturbations exit the horizon and the end of inflation follows [13]:
\[ N_{\text{COBE}} \simeq 66.9 + \frac{1}{4} \ln \left[ \frac{V(\phi_0)}{M_P^4} \right]. \]  
(12)
The amplitude of perturbations thus produced \(\delta_H\) and the scalar spectral index \(n_s\) are given by [13, 16]
\[ \delta_H = \frac{1}{5\pi} \sqrt{\frac{2}{3}} n(n-1)(n-2) \frac{m_\phi M_P}{\phi_0^2} \frac{1}{\Delta^2} \sin^2[N_{\text{COBE}} \sqrt{\Delta^2}], \]  
(13)
and
\[ n_s = 1 - 4\sqrt{\Delta^2} \cot[N_{\text{COBE}} \sqrt{\Delta^2}], \]  
(14)
where
\[ \Delta^2 \equiv n^2(n-1)^2 \alpha^2 N_{\text{COBE}}^{-2} \left( \frac{M_P}{\phi_0} \right)^4. \]  
(15)

III. PROPERTIES OF MSSM INFLATION

Experimental data on density perturbations can be used to specify the properties of MSSM inflation. The amplitude of perturbations is [19] \(\delta_H \approx 1.91 \times 10^{-5}\). Then, for weak scale supersymmetry \(m_\phi \sim 100 \text{ eV} - 1 \text{ TeV}\), Eqs. (11,13) require that \(n = 6\). This singles out inflaton candidates as there are only two flat directions which are lifted by \(n = 6\) superpotential terms [9]. One is the \(u\) direction in which case the inflaton is
\[ \phi = \frac{u_1^\alpha + d_j^\beta + d_k^\gamma}{\sqrt{3}}, \]  
(16)
Here \(u\) and \(d\) are the right-handed up- and down-type squarks respectively. The superscripts \(1 \leq \alpha, \beta, \gamma \leq 3\) are color indices, and the subscripts \(1 \leq i, j, k \leq 3\) denote the quark families. The flatness constraints require that \(\alpha \neq \beta \neq \gamma\) and \(j \neq k\). The other direction is \(LLe\) for which
\[ \phi = \frac{L_a^\alpha + L_b^\beta + e_k^\gamma}{\sqrt{3}}, \]  
(17)
where \(L\) and \(e\) are the left-handed and right-handed sleptons respectively. The superscripts \(1 \leq a, b \leq 2\) are the weak isospin indices and the subscripts \(1 \leq i, j, k \leq 3\) denote the lepton families. The flatness constraints require that \(a \neq b\) and \(i \neq j \neq k\).

Eqs. (11,13,14), combined with the \(2\sigma\) allowed region for the spectral index from 5-year WMAP data \(0.934 \leq n_s \leq 0.988\) [2], result in
\[ H_{\text{inf}} \sim 100 \text{ MeV} - 1 \text{ GeV}, \]  
(18)
\[ 2 \times 10^{-6} \leq \Delta^2 \leq 5.2 \times 10^{-6}, \]  
(19)

\(^4\) There also exists an inflaton candidate within a minimal extension of MSSM which is represented by the \(NH_uL\) flat direction [20] (\(N\) and \(L\) denote the right-handed sneutrino and left-handed slepton respectively). The inflaton in this case is given by \(\phi = (N + H_u + L)/\sqrt{3}\). The simplest extension of the SM gauge group that allows such a flat direction includes \(U(1)_{B-L}\), where \(B\) and \(L\) denote the baryon and lepton number respectively. If neutrinos are Dirac in nature, density perturbations of the correct size will be obtained for neutrino masses of \(O(0.1 \text{ eV})\) [20], which is the mass indicated by atmospheric neutrino oscillations detected by Super–Kamiokande experiment [21].
and the VEV of the inflection point \( \phi_0 \) turns out to be

\[ \phi_0 \sim 10^{14} \text{ GeV}. \]  

(20)

In figure 1, we show \( \delta_H \) as a function of \( n_s \) for different values of \( m_\phi \). The horizontal blue band shows the 2\( \sigma \) allowed region for \( \delta_H \). The vertical green shaded region is the 2\( \sigma \) allowed band for \( n_s \), corresponding to Eq. (19), and the region enclosed by solid lines shows the 1\( \sigma \) allowed region. This figure is drawn for \( \lambda \approx 1 \), see Eq. (1), which is natural in the context of effective field theory. It is seen that \( m_\phi \) within the 60 – 440 GeV range is compatible with the experimentally allowed ranges of \( n_s \) and \( \delta_H \).

A remarkable property of MSSM inflation, which is related to inflation occurring near a point of inflection, is that it can give rise to a wide range for scalar spectral index. Indeed it can yield a spectral index within the whole 2\( \sigma \) range allowed by 5-year WMAP data \( 0.934 \leq n_s \leq 0.988 \). This stands in contrast with other models (for example, chaotic inflation, hybrid inflation, natural inflation, etc. [16]) and makes the model very robust.

The low scale of inflation \( 18 \) and the sub-Planckian VEV of the inflection point \( 20 \) have important consequences:

- Gravitational waves that are produced during MSSM inflation are negligible and cannot be detected in future CMB experiments.
- The model is free from the cosmological moduli problem \([24, 25]\). The moduli obtain a mass \( \sim \mathcal{O}(\text{TeV}) \) from supersymmetry breaking. However, since in our case \( H_{\text{inf}} \sim \mathcal{O}(\text{GeV}) \), 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. This ensures the absence of the cosmological moduli problem in MSSM inflation.
- Supergravity corrections to the inflaton mass are negligible. These corrections typically induce a term \( \sim H_{\text{inf}}^2 \phi^2 \) for the inflaton potential \([25, 26, 27, 28]\). In our case this is subdominant to the mass term in Eq. (3) since \( H_{\text{inf}} \ll m_\phi \).
- The fact that \( \phi_0 \) is sub-Planckian guarantees that the inflationary potential is free from the uncertainties about physics at super-Planckian VEVs. Moreover, the smallness of \( H_{\text{inf}} \) also precludes any large trans-Planckian

---

5 Smaller values of \( \lambda \) will lead to an increase in \( m_\phi \) \([31]\).

6 Inflection point inflation and sensitivity of its predictions have also been discussed in models of D-brane inflation in string theory \([22, 23]\).
correction that would generically go as \((H_{\text{inf}}/M_*)^2\), where \(M_*\) is the scale at which one would expect these effects to show up \cite{25, 30}.

IV. MSSM INFLATION AND DARK MATTER

The inflaton mass \(m_\phi\) is related to the mass of squarks and sleptons according to

\[
m_\phi^2 = \frac{m_{u_1}^2 + m_{d_1}^2 + m_{s_1}^2}{3} \quad (\text{udd inflaton}),
\]

\[
m_\phi^2 = \frac{m_{L_1}^2 + m_{L_2}^2 + m_{e_1}^2}{3} \quad (\text{LLe inflaton}).
\]

The bound on \(m_\phi\) from density perturbations will then be translated into the bounds on the scalar masses. These bounds apply at the scales \(\sim \phi_0\), see Eq. (3), around which inflation occurs. To make a connection with sparticle masses at the weak scale, which will be probed at colliders, one should use appropriate Renormalization Group Equations (RGEs). The one-loop RGEs are given by \cite{13, 31}

\[
\frac{dm_\phi^2}{d\mu} = -\frac{1}{6\pi^2} \left( 4M_3^2g_3^2 + \frac{2}{5}M_1^2g_1^2 \right) \quad (\text{udd inflaton}),
\]

\[
\frac{dm_\phi^2}{d\mu} = -\frac{1}{6\pi^2} \left( \frac{3}{2}M_2^2g_2^2 + \frac{9}{10}M_2^2g_2^2 \right) \quad (\text{LLe inflaton}).
\] (21)

Here \(g_1, g_2, g_3\) and \(M_1, M_2, M_3\) are gauge couplings and gaugino masses of \(U(1), SU(2), SU(3)\) respectively. Explicit calculations can be done in the case of minimal supergravity (mSUGRA), which is motivated by unification of gauge couplings at the Grand Unified Theory (GUT) scale \(M_G \approx 2 \times 10^{16}\) GeV. The models of mSUGRA depend only on four parameters and one sign: \(m_0\) (the universal scalar mass at \(M_G\)); \(m_{1/2}\) (the universal gaugino mass at \(M_G\)); \(A_0\) (the universal trilinear \(A\)-term at \(M_G\)); \(\tan \beta = (H_u)/(H_d)\) at the electroweak scale (where \(H_u\) and \(H_d\) give masses to up-type and down-type quarks respectively); and the sign of \(\mu\), the Higgs mixing parameter in the superpotential \((W_\mu = \mu H_u H_d)\).

The model parameters are already significantly constrained by different experimental results \cite{32, 33, 34, 35}. A further constraint arises from the requirement that the Lightest Supersymmetric Particle (LSP), which is stable in models with conserved \(R\)-parity, has the right relic abundance to be the cold dark matter (CDM) in the universe. The \(1\sigma\) bound from latest cosmological data \cite{2} gives a relic density bound of \(0.22 < \Omega_{\text{CDM}} < 0.246\) for CDM. In mSUGRA the lightest neutralino is the dark matter candidate. The allowed parameter space, at present, has mostly three distinct regions selected out by dark matter constraints \cite{36, 37, 38, 39, 40, 41, 42, 43}: (i) the stau-neutralino coannihilation region, (ii) the focus point region, and (iii) the scalar Higgs annihilation funnel region.

We show the mSUGRA parameter space in figure 2 for \(\tan \beta = 10, 40\) with the \(udd\) flat direction as the inflaton using \(\lambda \approx 1\) (the figure for the \(LLe\) flat direction is similar) \cite{7}. The contours correspond to different values of \(n_s\) within the \(2\sigma\) range allowed by 5-year WMAP data, for \(\delta_H = 1.91 \times 10^{-5}\). The constraints on the parameter space arising from inflation are compatible with those from dark matter and other experimental results. It is seen that \(\tan \beta\) needs to be smaller to allow for smaller values of \(n_s\). It is also interesting to note that the allowed region of \(m_\phi\) lies in the stau-neutralino coannihilation region which requires smaller values of the supersymmetric particle masses. The supersymmetric particles in this parameter space are, therefore, within the reach of the LHC very quickly. The detection of this region at the LHC has been considered in Ref. \cite{45}.

V. REHEATING AFTER MSSM INFLATION

After the end of inflation, the inflaton \cite{10, 17} rolls towards the global minimum of its potential at the origin. At this stage the dominant term in the scalar potential will be \(m_\phi \phi^2/2\), see Eq. (3), which results in oscillations with a frequency of \(m_\phi\). Since \(m_\phi \approx 10^3 H_{\text{inf}}\), the inflaton oscillates about the origin a large number of times within the

\footnote{Inflation and dark matter can be unified in the case that the \(N H_u L\) flat direction is the inflaton \cite{44}. The inflaton \(\phi = (N + H_u + L)/\sqrt{3}\) has a right-handed sneutrino component which can be the dark matter candidate. Sneutrino dark matter can be seen in the upcoming direct detection experiments and also be produced at the LHC. For more details, see Ref. \cite{44}.}
first Hubble time after the end of inflation. Hence the effect of expansion is negligible. Since the inflaton is a linear combination of squarks \(16\) or sleptons \(17\), it has gauge couplings to the gauge/gaugino fields and Yukawa couplings to the Higgs/Higgsino fields. To elucidate the physics, we consider the case when the \(Le\) flat direction is the inflaton.

The situation for the \(udd\) flat direction as the inflaton is similar. The VEV of the inflaton spontaneously breaks \(SU(2) \times U(1)\) symmetry, and therefore, induces a supersymmetry conserving mass to the electroweak gauge/gaugino fields (similar to what happens in electroweak symmetry breaking fashion via the Higgs mechanism). When the flat direction goes to its minimum this mass vanishes, and the gauge symmetry is restored. However, the mass undergoes a non-adiabatic time variation every time that the inflaton crosses the origin. This results in an efficient creation of gauge and gaugino quanta with a physical momentum \(k \lesssim (gm_\phi \phi_0)^{1/2}\) within a short interval \(\Delta t \sim (gm_\phi \phi_0)^{-1/2}\), where \(\phi_0\) is given by Eq. (5) and \(g\) is the corresponding gauge coupling. The number density of the gauge/gaugino quanta thus produced is given by \(46, 47\)

\[
n_g \approx \frac{(gm_\phi \phi_0)^{3/2}}{8\pi^3}.
\]  

(22)

As the inflaton VEV is rolling back to its maximum value \(\phi_0\), the mass of the produced quanta 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 total decay rate of the gauge/gaugino fields is then given by \(\Gamma = C (g^2/48\pi) g\phi\), where \(C \sim \mathcal{O}(10)\) is a numerical factor counting the multiplicity of final states.

The decay of gauge and gaugino quanta happens very quickly and converts a fraction \(f\) of the inflaton energy density (for details, see \(13\)) to relativistic (s)quarks where

\[
f \sim 10^{-2} g.
\]  

(23)

This is the so-called instant preheating mechanism \(48\). The rapid conversion of fraction \(f\) of the inflaton energy density into relativistic particles happens twice in each oscillation. Note that there will be hundreds of oscillations within the first few Hubble times after the end of inflation since \(m_\phi \sim 10^3 H_{\text{inf}}\). Reheating is therefore quite efficient.
in this model as almost all the energy density in the inflaton will decay into radiation within a couple of Hubble times. The resulting reheat temperature of the universe will be

\[ T_R \sim 10^7 \text{ GeV}, \]  

(24)

which is sufficiently low to avoid thermal overproduction of gravitinos (for more details, see [13]).

VI. DISCUSSION AND CONCLUSION

The existence of a point of inflection in the scalar potential of two MSSM flat directions provides all the necessary ingredients for a realistic and successful model of inflation. The exceptional feature of the model, which sets it apart from conventional singlet field inflation models, is that here the inflaton is not added as an ad hoc field whose sole purpose is to drive inflation. Instead it is a combination of the squark and slepton fields, and hence its couplings to the matter and gauge fields are known. This not only gives the inflaton a natural place within particle physics, but also makes it possible to address reheating after inflation in an unambiguous way.

MSSM inflation occurs at a low scale, corresponding to a Hubble expansion rate \( H_{\text{inf}} \sim \mathcal{O}(\text{GeV}) \) and sub-Planckian field values. This implies negligible supergravity and trans-Planckian corrections and solves the cosmological moduli problem. The model is robust as it can give rise to density perturbations of the correct size with a scalar spectral index in the entire range allowed by the 5-year WMAP data. The parameter space for inflation is compatible with that of supersymmetric dark matter. \( \lambda \sim \mathcal{O}(1) \) (as expected in an effective field theory approach) can be explained. In the context of mSUGRA the stau-neutralino coannihilation region is most preferred to satisfy the dark matter content of the universe. The masses of supersymmetric particles in this region are mostly within the reach of the LHC.

Reheating after MSSM inflation is very efficient. Non-perturbative production of gauge and gaugino fields and their subsequent decay to relativistic particles result in a reheat temperature \( T_R \sim 10^7 \text{ GeV} \) despite the small value of \( H_{\text{inf}} \). This is sufficiently high for various mechanisms of generating the baryon asymmetry of the universe and produce thermal dark matter while avoiding overproduction of gravitinos.

The existence of the inflection point requires a fine-tuning of the ratio of the soft supersymmetry breaking parameters \( m_{\phi} \) and \( A \). Radiative corrections change \( A/m_{\phi} \) only slightly so that inflection point inflation can always be achieved for some value of this ratio. A more detailed investigation is required to address the fine-tuning issue, but it is warranted by the success of MSSM inflation, which 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.

VII. ACKNOWLEDGMENTS

The author is indebted to Bhaskar Dutta, Kari Enqvist, Andrew Frey, Juan Garcia-Bellido, Asko Jokinen, Alex Kusenko and Anupam Mazumdar for collaboration and numerous discussions on various aspects of the physics presented here. He wishes to thank Robert Brandenberger, Cliff Burgess, Manuel Drees, Gordy Kane, Justin Khoury, David Lyth, Guy Moore, Scott Thomas and Maxim Pospelov for valuable discussions on this subject. Katie Richardson-McDaniel is acknowledged for careful reading of this manuscript and useful suggestions. This research was supported in part by Perimeter Institute for Theoretical Physics.

[1] For a review, see: A. D. Linde, hep-th/0503203.
[2] E. Komatsu et al. [WMAP Collaboration], arXiv:0803.0547 [astro-ph].
[3] For a review, see: D. H. Lyth and A. Riotto, Phys. Rept. 314, 1 (1999).
[4] H. Murayama, H. Suzuki, T. Yanagida and J. Yokoyama, Phys. Rev. Lett. 70, 1912 (1993).
[5] G. Lazarides and Q. Shafi, Phys. Lett. B 308, 17 (1993).
[6] S. Kasuya, T. Moroi and F. Takahashi, Phys. Lett. B 593, 33 (2004).
[7] K. A. Olive, G. Steigman and T. P. Walker, Phys. Rept. 333, 389 (2000).
[8] For a review on supersymmetry, see: H. P. Nilles, Phys. Rept. 110, 1 (1984).
[9] T. Gherghetta, C. F. Kolda and S. P. Martin, Nucl. Phys. B 468, 37 (1996).
[10] K. Enqvist and A. Mazumdar, Phys. Rept. 380, 99 (2003).
[11] M. Dine and A. Kusenko, Rev. Mod. Phys. 76, 1 (2004).
[12] R. Allahverdi, K. Enqvist, J. Garcia-Bellido and A. Mazumdar, Phys. Rev. Lett. 97, 191304 (2006).
[13] R. Allahverdi, K. Enqvist, J. Garcia-Bellido, A. Jokinen and A. Mazumdar, JCAP 0706, 019 (2007).
[14] M. Dine, L. Randall and S. Thomas, Phys. Rev. Lett. 75, 398 (1995).
[15] M. Dine, L. Randall and S. Thomas, Nucl. Phys. B 458, 291 (1996).
[16] J. C. Bueno Sanchez, K. Dimopoulos and D. H. Lyth, JCAP 0701, 015 (2007).
[17] R. Allahverdi, A. R. Frey and A. Mazumdar, Phys. Rev. D 76, 026001 (2007).
[18] R. Allahverdi, B. Dutta and A. Mazumdar, [arXiv:0806.4557] [hep-ph] (to appear in Phys. Rev. D).
[19] A. R. Liddle, D. Parkinson, S. M. Leach and P. Mukherjee, Phys. Rev. D 74, 083512 (2006).
[20] R. Allahverdi, A. Kusenko and A. Mazumdar, JCAP 0707, 018 (2007).
[21] Y. Fukuda et al. [Super-Kamiokande Collaboration], Phys. Rev. Lett. 81, 1562 (1998).
[22] D. Baumann, A. Dymarsky, I. R. Klebanov, L. McAllister and P. J. Steinhardt, Phys. Rev. Lett. 99, 141601 (2007).
[23] D. Baumann, A. Dymarsky, I. R. Klebanov and L. McAllister, JCAP 0801, 024 (2008).
[24] G.D. Coughlan, W. Fischler, E. W. Kolb, S. Raby and G. G. Ross, Phys. Lett. B 131, 59 (1983).
[25] M. Dine, W. Fischler and D. Nemeschansky, Phys. Lett. B 136, 169 (1984).
[26] G. D. Coughlan, R. Holman, P. Ramond and G. G. Ross, Phys. Lett. B 140, 44 (1984).
[27] A. S. Goncharov, A. D. Linde and M. I. Vysotsky, Phys. Lett. B 147, 279 (1984).
[28] O. Bertolami and G. G. Ross, Phys. Lett. B 183, 163 (1987).
[29] N. Kaloper, M. Kleban, A. E. Lawrence and S. Shenker, Phys. Rev. D 66, 123510 (2002).
[30] C. P. Burgess, J. Cline and R. Holman, JCAP 0310, 004 (2003).
[31] R. Allahverdi, B. Dutta and A. Mazumdar, Phys. Rev. D 75, 075018 (2007).
[32] C. Amsler et al. [Particle Data group], Phys. Lett. B 667, 1 (2008).
[33] G. Bennett et al. [Muon g − 2 Collaboration], Phys. Rev. Lett. 92, 161802 (2004).
[34] G. Abbiendi et al. [The LEP Working Group for Higgs Boson Searches], Phys. Lett. B 565, 61 (2003).
[35] M. Alam et al. [CLEO Collaboration], Phys. Rev. Lett. 74, 2885 (1995).
[36] J. Ellis, K. A. Olive, Y. Santoso, and V. Spanos, Phys. Lett. B 565, 176 (2003).
[37] R. Arnowitt, B. Dutta and B. Hu, arXiv:hep-ph/0310103.
[38] H. Baer, C. Balazs, A. Belyaev, T. Krupovnickas and X. Tata, JHEP 0306, 054 (2003).
[39] B. Lahanas and D. V. Nanopoulos, Phys. Lett. B 568, 55 (2003).
[40] U. Chattopadhyay, A. Corsetti and P. Nath, Phys. Rev. Lett 68, 035005 (2003).
[41] E. Baltz and P. Gondolo, JHEP 0410, 052 (2004).
[42] A. Djouadi, M. Drees and J. L. Kneur, JHEP 0603, 033 (2006).
[43] J. L. Feng, K. T. Matchev and F. Wilczek, Phys. Lett. B 482, 388 (2000).
[44] R. Allahverdi, B. Dutta and A. Mazumdar, Phys. Rev. Lett. 99, 261301 (2007).
[45] R. Arnowitt, B. Dutta, A. Gurrola, T. Kamon, A. Krislock and D. Toback, arXiv:0802.2968 [hep-ph].
[46] L. Kofman, A. D. Linde and A. A. Starobinsky, Phys. Rev. Lett. 73, 3195 (1994).
[47] L. Kofman, A. D. Linde and A. A. Starobinsky, Phys. Rev. D 56, 3258 (1997).
[48] G. N. Felder, L. Kofman and A. D. Linde, Phys. Rev. D 59, 123523 (1999).