PROCEEDINGS OF SCIENCE

Challenges in Particle Physics and Cosmology

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We discuss the possibility that inflation is driven by supersymmetry breaking with the superpartner of the goldstino (sgoldstino) playing the role of the inflaton and charged under a gauged $U(1)$ R-symmetry. Imposing a linear superpotential allows to satisfy easily the slow-roll conditions, avoiding the so-called η-problem, and leads to an interesting class of small field inflation models, characterised by an inflationary plateau around the maximum of the scalar potential near the origin, where R-symmetry is restored with the inflaton rolling down to a minimum describing the present phase of the Universe. Inflation can be driven by either an F- or a D-term, while the minimum has a positive tuneable vacuum energy. The models agree with cosmological observations and in the simplest case predict a rather small tensor-to-scalar ratio of primordial perturbations. Upon coupling the inflaton sector to the (supersymmetric) Standard Model, we examined decay modes of the inflaton, with the resulting reheating temperature around 10^8 GeV.

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1. Introduction

Several experiments have raised inflation to the level of a cornerstone of modern Cosmology, but many important questions, including the origin of inflation, the inflaton field and its associated symmetries protecting its mass, the issue of initial conditions for inflation and the nature of dark energy and dark matter, among other fundamental questions, remain in need of an answer. A first step towards the understanding of these puzzles is to study possible connections between the electroweak scale of the Standard Model or its possible extension (such as the supersymmetry breaking scale) with that of inflation. An additional constraint would be to impose at the electroweak vacuum the presence of a tiny tuneable cosmological constant in order to accommodate the observed dark energy, without necessarily trying to explain it.

Despite the absence of evidence of low energy supersymmetry at the Large Hadron Collider at CERN, a substantial part of the theoretical community believes that supersymmetry should play a role at some very fundamental level. However, inflationary models in supergravity suffer in general from several problems, such as fine-tuning to satisfy the slow-roll conditions, large field initial conditions that break the validity of the effective field theory and stabilisation of the (pseudo) scalar companion of the inflaton arising from the fact that the number of bosonic components of superfields are always even. A solution to all three problems, with the sgoldstino as inflaton, was recently proposed in [1] and [2]. The model has a gauged R-symmetry and generalises models of the so-called “minimal inflation". The superpotential is linear and the slow-roll conditions are automatically satisfied. Also, since the inflation arises at a plateau around the maximum of the scalar potential (hill-top) no large field initial conditions are needed, while the pseudo-scalar companion of the inflaton is absorbed into the R-gauge field that becomes massive, leading to the inflaton being present as a single scalar field in the low-energy spectrum. Moreover, this model allows the presence of a realistic minimum describing our present Universe with an infinitesimal positive vacuum energy arising due to a cancellation between F- and D-term contributions to the scalar potential. This proposal was studied using an effective field theory approach, in perturbation around the origin of the inflaton field potential where R-symmetry is restored. Both cases have been analysed in detail, corresponding to inflation dominated by F-term or D-term supersymmetry breaking. The second case is possible only in the presence of a new Fayet-Iliopoulos (FI) term constructed recently [3, 4].

The Outline of this chapter is the following. In Sect. 2, we briefly review the framework of inflation by supersymmetry breaking and explain the coupling of the supersymmetry breaking sector to the MSSM. We then estimate reheating temperature in Sect. 3. In Sect. 4, we review the new FI term analyze the consequences of the new term in the models of inflation driven by supersymmetry breaking.

	Q	$\bar{u}$	$\bar{d}$	L	$\bar{e}$	$\tilde{H}_u$	$\tilde{H}_d$	ζ	λ_R	λ_1	λ_2	λ_3
$U(1)_R$	$-\frac{1}{2}$	$-\frac{1}{2}$	$-\frac{1}{2}$	$-\frac{1}{2}$	$-\frac{1}{2}$	$-\frac{1}{2}$	$-\frac{1}{2}$	$\frac{1}{2}$	$\frac{1}{2}$	$\frac{1}{2}$	$\frac{1}{2}$	$\frac{1}{2}$
$U(1)_Y$	$\frac{1}{6}$	$-\frac{2}{3}$	$\frac{1}{3}$	$-\frac{1}{2}$	1	$\frac{1}{2}$	$-\frac{1}{2}$	0	0	0	0	0
$SU(2)_L$	**2**	**1**	**1**	**2**	**1**	**2**	**2**	**1**	**1**	**1**	**3**	**1**
$SU(3)_c$	**3**	$\bar{\mathbf{3}}$	$\bar{\mathbf{3}}$	**1**	**1**	**1**	**1**	**1**	**1**	**1**	**1**	**8**

Table 1: MSSM and $U(1)_R$ charges of the fermions. ζ is the inflatino, λ_R is the $U(1)_R$ gaugino, and $\lambda_{1,2,3}$ are bino, wino, and gluino, respectively. The gravitino has the same R-charge as λ_R.

2. Inflation by Supersymmetry Breaking

The starting point is a class of models with gauged $U(1)_R$ phase symmetry, defined by Kähler potential and superpotential,

$$\mathcal{K}(X,\overline{X},\phi,\bar{\phi}) = \sum \phi\bar{\phi} + J(X\overline{X}) \,, \tag{1}$$

$$\mathcal{W}(X,\phi) = \kappa[f\kappa^{-3} + \Omega(\phi)]X \,, \tag{2}$$

where X is the inflaton/sgoldstino superfield, ϕ collectively denotes matter superfields, and J is the inflaton Kähler potential. In the superpotential, f is a dimensionless real constant, while Ω describes the MSSM part

$$\Omega = \hat{y}_u\bar{u}QH_u - \hat{y}_d\bar{d}QH_d - \hat{y}_e\bar{e}LH_d + \hat{\mu}H_uH_d \,. \tag{3}$$

Here $\bar{u},\bar{d},\bar{e},Q,L,H_u,H_d$ are chiral superfields. As usual, we denote the corresponding SM matter fields (quarks, leptons, and Higgs fields) with the same character, while tildes will be used for their superpartners (squarks, sleptons, and Higgsinos). The un-normalized Yukawa couplings y and the μ-parameter are denoted by hats which will be removed after proper rescaling, once X settles at the minimum.

The total gauge group of the model is,

$$SU(3)_c \times SU(2)_L \times U(1)_Y \times U(1)_R \,. \tag{4}$$

Squarks, sleptons, and Higgs scalars are neutral under $U(1)_R$, while X carries the same R-charge as the superpotential. The R-charges of the MSSM fermions are fixed later as in Table 1.

In this class of models, the X-dependent part of the potential drives inflation, after which X and its auxiliary field F^X settle at non-zero vacuum expectation values (VEVs), spontaneously breaking both supersymmetry (SUSY) and $U(1)_R$. At the minimum of the potential, the gravitino mass and the the auxiliary fields of X and $U(1)_R$ are given by,

$$\begin{aligned} m_{3/2} &= f\langle e^{\kappa^2 J/2}|X|\rangle \,, \\ \langle F^X\rangle &= -f\langle e^{\kappa^2 J/2}J^{X\bar{X}}(\kappa^{-2} + J_X X)\rangle \,, \\ \langle \mathcal{D}_R\rangle &= g\langle \kappa^{-2} + J_X X\rangle \,, \end{aligned} \tag{5}$$

where we assume that matter fields ϕ vanish at the minimum.

The Yukawa couplings $\hat{y}$ and the parameter $\hat{\mu}$ in (3) are related to their properly normalized versions y and μ by

$$\{\hat{y},\hat{\mu}\} = \left\langle \frac{e^{-\kappa^2 J/2}}{\kappa|X|} \right\rangle \times \{y,\mu\} \ . \tag{6}$$

This is due to the overall factor of $e^{\mathcal{K}}$ in the F-term potential, as well as the coupling of Ω to X as shown in Eq. (2). At the minimum, X and $J(X,\overline{X})$ take non-vanishing VEVs, which leads to this rescaling.

The scalar potential can be written as $\mathcal{V} = \mathcal{V}_F + \mathcal{V}_D$ where [1]

$$\mathcal{V}_F = e^{\kappa^2\mathcal{K}}\left\{\mathcal{K}^{I\bar{J}} D_I\mathcal{W} D_{\bar{J}}\overline{\mathcal{W}} - 3\kappa^2|\mathcal{W}|^2\right\} , \tag{7}$$

$$\mathcal{V}_D = \frac{1}{2}\mathrm{Re}(\mathcal{F}^{AB})\mathcal{D}_A\mathcal{D}_B \ . \tag{8}$$

In our notation, the indices I,J run through all the chiral (super)fields, while A,B are the gauge group indices. The relevant part of supergravity Lagrangian that we use here can be found in the Appendix of Ref. [5] and its derivation in Ref. [6].

For the gauge kinetic matrix, we use $\mathcal{F}^{AB} \equiv \mathcal{F}^{-1}_{AB}$. Kähler covariant derivatives are defined as $D_I\mathcal{W} \equiv \mathcal{W}_I + \kappa^2\mathcal{K}_I\mathcal{W}$, where the indices denote the respective partial derivatives. The Killing potential and Killing vector are related by

$$\mathcal{D}_A = ik_A^I\left(\mathcal{K}_I + \kappa^{-2}\frac{\mathcal{W}_I}{\mathcal{W}}\right) , \tag{9}$$

where the gauge couplings and charges are included in the Killing vectors k_A^I. For example, if X transforms under $U(1)_R$ as $X \to Xe^{-igq\vartheta}$ (where ϑ is a transformation parameter and q is its R-charge), its Killing vector is $k_R^X = -igqX$. The gauge couplings of $U(1)_R$, $U(1)_Y$, $SU(2)_L$, and $SU(3)_c$ are g, g_1, g_2, and g_3, respectively.

We use a convention where the superpotential transforms under $U(1)_R$ with unit R-charge, $\mathcal{W} \to \mathcal{W}e^{-ig\vartheta}$, and the fermionic superspace coordinate transforms with half-unit R-charge, $\theta \to \theta e^{-ig\vartheta/2}$. Then X has unit R-charge, while its fermionic partner has half-unit R-charge. For a scalar field with R-charge q, its fermionic partner has R-charge $q-1/2$. With this convention, the fermion charges under the total gauge group of our model are summaised in Table 1.

Let us focus on the possibility of inflation in this model by ignoring matter fields so that $\mathcal{K} = J$ and $\mathcal{W} = \kappa^{-2}fX$. Here we would like to introduce a simpler choice of J with finite number of perturbative corrections, namely,

$$J = X\overline{X} + A\kappa^2(X\overline{X})^2 + B\kappa^4(X\overline{X})^3 , \tag{10}$$

where the parameters A and B are dimensionless. One should think of the above form as a perturbative expansion around the canonical kinetic terms with coefficients less than unity. Then

[1] The mass dimensions of Kähler potential, superpotential, gauge kinetic function, Killing potential, and Killing vector are $[\mathcal{K}] = M^2$, $[\mathcal{W}] = M^3$, $[\mathcal{F}_{AB}] = M^0$, $[\mathcal{D}_A] = M^2$, $[X_A^I] = M$, respectively while scalar fields have canonical mass dimension M.

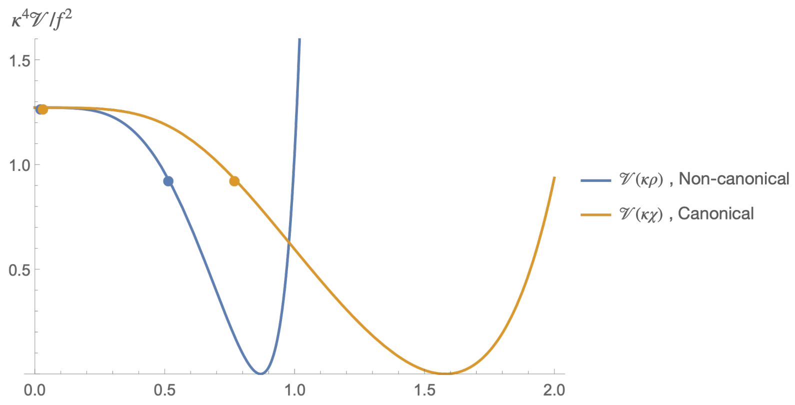

Figure 1: Scalar potential (11) for the parameter set (23). Both non-canonical (ρ) and canonical (χ) parametrizations are shown, where the latter is found numerically. The markers represent the start and end of 60 e-folds of inflation (the starting point of inflation almost coincides for the two curves).

the scalar potential reads,

$$\mathcal{V} = \frac{f^2}{\kappa^4} e^{\left(|\kappa X|^2 + A|\kappa X|^4 + B|\kappa X|^6\right)} \left\{ \frac{(1 + |\kappa X|^2 + 2A|\kappa X|^4 + 3B|\kappa X|^6)^2}{1 + 4A|\kappa X|^2 + 9B|\kappa X|^4} - 3|\kappa X|^2 \right\} + \frac{g^2}{2\kappa^4}\left(1 + |\kappa X|^2 + 2A|\kappa X|^4 + 3B|\kappa X|^6\right)^2 , \tag{11}$$

where we set gauge kinetic function $\mathcal{F} = 1$ for now. Note that we can write

$$X = \rho e^{i\theta}. \tag{12}$$

Thus, the scalar potential is only a function of the modulus ρ and we can indentify the field ρ as the inflaton

$$\mathcal{V} = \frac{f^2}{\kappa^4} \exp\left(\kappa^2\rho^2 + A\kappa^4\rho^4 + B\kappa^6\rho^6\right) \left\{ \frac{(1 + \kappa^2\rho^2 + 2A\kappa^4\rho^4 + 3B\kappa^6\rho^6)^2}{1 + 4A\kappa^2\rho^2 + 9B\kappa^4\rho^4} - 3\kappa^2\rho^2 \right\} + \frac{g^2}{2\kappa^4}\left(1 + \kappa^2\rho^2 + 2A\kappa^4\rho^4 + 3B\kappa^6\rho^6\right)^2 . \tag{13}$$

The phase θ get absorbed by the $\mathrm{U}(1)_\mathrm{R}$ gauge field in the standard Brout-Englert-Higgs mechanism. However, in order to calculate the slow-roll parameters, we introduce the canonically normalised field χ satisfying

$$\frac{d\chi}{d\rho} = \sqrt{2\mathcal{K}_{X\bar{X}}}. \tag{14}$$

The slow-roll parameters can be defined in terms of the canonical field χ as:

$$\epsilon = \frac{1}{2\kappa^2}\left(\frac{d\mathcal{V}/d\chi}{\mathcal{V}}\right)^2, \quad \eta = \frac{1}{\kappa^2}\frac{d^2\mathcal{V}/d\chi^2}{\mathcal{V}}. \tag{15}$$

The number of e-folds N during inflation is determined by

$$N = \kappa^2 \int_{\chi_*}^{\chi_\mathrm{end}} \frac{\mathcal{V}}{\partial_\chi \mathcal{V}} d\chi = \kappa^2 \int_{\rho_*}^{\rho_\mathrm{end}} \frac{\mathcal{V}}{\partial_\rho \mathcal{V}} \left(\frac{d\chi}{d\rho}\right)^2 d\rho, \tag{16}$$

where we choose $|\eta(\chi_{\text{end}})| = 1$ for this section. The scalar potential as a function of non-canonical (ρ) and canonical (χ) parametrizations are shown in Fig. 1. Notice that our model is classified as hilltop inflation which may encounter the so-called overshoot problems. It was shown in [7] that initial condition for models of this type can be determined quantum mechanically such that the inflaton is driven toward the slow-roll attractor solution exponentially fast and the overshoot problem is partially evade.

Since inflation arises near the maximum $\kappa\rho = 0$, we expand

$$\epsilon = 4\left(\frac{-4A + y^2}{2 + y^2}\right)^2 (\kappa\rho)^2 + O(\rho^4), \tag{17}$$

$$\eta = 2\left(\frac{-4A + y^2}{2 + y^2}\right) + O(\rho^2), \tag{18}$$

where we defined $y = g/f$. The above equation implies $\epsilon \simeq \eta^2(\kappa\rho)^2 \ll \eta$. For simplicity, we focus on the special case $y \to 0$ where F-term contribution to the scalar potential is dominant. By considering the behaviour near the origin, we can put some constraints on the coefficient A of the quadratic term of the Kähler potential defined in (10). We can easily show that $A > 0$ is required for having a local maximum of the scalar potential at $\rho = 0$. Furthermore, the slow-roll condition $|\eta| \ll 1$ sets an upper bound $A \ll 0.25$. Taking these requirements into account, the constraint on A is

$$0 < A \ll 0.25. \tag{19}$$

We can choose $A \sim 0.005$ to obtain $\eta \sim -0.02$ which is in agreement with CMB observational data. [2] In the following, we will compare our theoretical predictions to the CMB observational data. The amplitude of density fluctuations A_s, the spectral index n_s and the tensor-to-scalar ratio r can be written in terms of the slow-roll parameters:

$$\mathcal{A}_s = \frac{\kappa^4 \mathcal{V}_*}{24\pi^2\epsilon_*}, \tag{20}$$

$$n_s = 1 + 2\eta_* - 6\epsilon_* \simeq 1 + 2\eta_*, \tag{21}$$

$$r = 16\epsilon_*, \tag{22}$$

evaluated at the horizon exit.

As a concrete example we choose the following parameter values

$$A = 0.139\,, \quad B = 0.6\,, \quad y = 0.7371\,, \quad f = 2.05 \times 10^{-7}\,, \tag{23}$$

which leads to the inflationary parameters

$$\mathcal{A}_s = 2.1 \times 10^{-9}\,, \quad n_s = 0.9543\,, \quad r = 1.72 \times 10^{-6}\,, \tag{24}$$

for 60 e-fold. The Hubble parameter during inflation is

$$H_{\text{inf}} = \kappa\sqrt{\mathcal{V}_*/3} = 3.25 \times 10^{11}\ \text{GeV}\,. \tag{25}$$

[2] In [8], a generalisation version of the Fayet-Iliopoulos (FI) model [9] was introduced as an example of the microscopic origin for the effective field theory of this class of inflation models.

the scalar potential depicted in Figure 1 where the non-canonical scalar ρ is shown in blue, while the canonically normalized scalar χ, found numerically, is shown in orange. The corresponding inflaton VEV is $\langle\kappa\rho\rangle = 0.89$.

3. Reheating after inflation by supersymmetry breaking

In our model defined by (1) and (2) (with general J) soft scalar masses are universal,

$$m_Q^2 = m_u^2 = m_d^2 = m_L^2 = m_e^2 = m_{H_u}^2 = m_{H_d}^2 = m_0^2\,, \tag{26}$$

where m_0^2 is given by

$$m_0^2 = \kappa^2\langle J_{X\bar{X}}F^X\overline{F}^X\rangle - 2m_{3/2}^2\,, \tag{27}$$

and for the MSSM μ-parameter we assume $|\mu| \ll |m_0|$ to avoid extreme fine-tuning of the Higgs boson mass (since m_0 is close to the inflationary scale). From the requirement of Minkowski minimum, we have the relation,

$$\langle V\rangle = \langle J_{X\bar{X}}F^X\overline{F}^X\rangle - 3\kappa^{-2}m_{3/2}^2 + \tfrac{1}{2}\langle\mathcal{D}_R\rangle^2 = 0\,. \tag{28}$$

The relation (28) allows us to rewrite m_0^2 in terms of the D-term contribution,

$$m_0^2 = m_{3/2}^2 - \tfrac{\kappa^2}{2}\langle\mathcal{D}_R\rangle^2\,, \tag{29}$$

and this leads to the requirement $m_{3/2} > \kappa\langle\mathcal{D}_R\rangle/\sqrt{2}$, in order to avoid tachyonic instabilities in the MSSM sector.

For the bilinear H_uH_d coupling we have

$$e^{-1}\mathcal{L} \supset -B_0\mu H_uH_d + \text{h.c.}\,, \tag{30}$$

where

$$B_0 = \frac{\kappa^2\langle J_{X\bar{X}}F^X\overline{F}^X\rangle - m_{3/2}^2}{m_{3/2}}\,. \tag{31}$$

The MSSM gaugino masses are generated at one loop via the Green–Schwarz mechanism of anomaly cancellation, where the gauge anomalies due to triangle diagrams involving the fermions (all the fermions of the model carry non-zero R-charges) are cancelled by appropriate $U(1)_R$ transformations of the following terms depending on the imaginary part of the gauge kinetic matrix:

$$e^{-1}\mathcal{L} \supset \tfrac{1}{8}\mathrm{Im}(\mathcal{F}_{AB})\epsilon^{mnkl}F^A_{mn}F^B_{kl}\,. \tag{32}$$

The gauge kinetic matrix takes the form,

$$\mathcal{F}_{AB} = \begin{pmatrix} \mathcal{F}_R & & & \\ & \mathcal{F}_1 & & \\ & & \mathcal{F}_2 & \\ & & & \mathcal{F}_3 \end{pmatrix}, \tag{33}$$

where $\mathcal{F}_{R,1,2,3}$ are gauge kinetic functions for $U(1)_R$, $U(1)_Y$, $SU(2)_L$, and $SU(3)_c$, respectively. To cancel the anomalies we fix these kinetic functions as,

$$\mathcal{F}_R = 1 + \beta_R \log(\kappa\rho)\,, \tag{34}$$
$$\mathcal{F}_a = 1 + \beta_a \log(\kappa\rho)\,, \tag{35}$$

where $a = 1, 2, 3$ stands for the Standard Model gauge groups. Here β are constants which we determine by using the methods described in Refs. [10–12]. The result is,

$$\beta_R = -\frac{g^2}{3\pi^2}\,, \quad \beta_1 = -\frac{11 g_1^2}{8\pi^2}\,, \quad \beta_2 = -\frac{5 g_2^2}{8\pi^2}\,, \quad \beta_3 = -\frac{3 g_3^2}{8\pi^2}\,, \tag{36}$$

where β_R is found from the cancellation of $U(1)_R^3$ anomaly, β_1 from $U(1)_R \times U(1)_Y^2$ anomaly, β_2 from $U(1)_R \times [SU(2)_L]^2$ anomaly, and β_3 from $U(1)_R \times [SU(3)_c]^2$ anomaly.

The values of β_a are the same as in the model of Ref. [13], because the MSSM fermions in the two models have the same R-charges, while β_R is different due to the difference in the hidden sector fermion (inflatino) R-charges. Since g/κ in our models is not far from the Hubble scale (e.g. the parameter choice (23) leads to $g \sim 10^{-7}$), we have

$$\mathcal{F}_R = 1 + \beta_R \log(\kappa\rho) \approx 1\,, \tag{37}$$

if $\kappa\rho$ is around unity. Gauged $U(1)_R$ also leads to a gravitational anomaly which can be cancelled in a similar fashion (see for example Refs. [10–12]).

This brings us to the MSSM gaugino masses,

$$m_{ab} = \frac{1}{2}\left|\langle F^X \partial_X \mathcal{F}_{ab}\rangle\right| = \frac{f}{2}\left|\left\langle e^{\kappa^2 J/2} J^{X\bar{X}} (\kappa^{-2} + J_{\bar{X}} \bar{X}) \partial_X \mathcal{F}_{ab}\right\rangle\right|\,. \tag{38}$$

Using Eq. (35) we get,

$$m_a = \left|\frac{\langle \kappa F^X\rangle \beta_a}{2\langle\kappa X\rangle}\right|\,, \tag{39}$$

where we denote $m_a \equiv m_{aa}$. Finally, m_a should be rescaled after taking into account non-canonical kinetic terms of the gaugini,

$$e^{-1}\mathcal{L} \supset -\tfrac{i}{2}\langle \mathrm{Re}\mathcal{F}_a\rangle \lambda^a \sigma^m D_m \bar{\lambda}^a + \text{h.c.} = -\tfrac{i}{2}(1 + \beta_a \log\langle\kappa X\rangle)\lambda^a \sigma^m D_m \bar{\lambda}^a + \text{h.c.} \tag{40}$$

However, if $|\beta_a| \log\langle\kappa X\rangle \ll 1$, as in the models that we consider here, the rescaling of the gaugini can be neglected. The trilinear couplings between the MSSM scalars are

$$e^{-1}\mathcal{L} \supset -A_0(y_u \bar{\tilde{u}} \tilde{Q} H_u - y_d \bar{\tilde{d}} \tilde{Q} H_d - y_e \bar{\tilde{e}} \tilde{L} H_d) - \mu(y_u \bar{\tilde{u}} \tilde{Q} \overline{H}_d - y_d \bar{\tilde{d}} \tilde{Q} \overline{H}_u - y_e \bar{\tilde{e}} \tilde{L} \overline{H}_u) + \text{h.c.}\,, \tag{41}$$

where for A_0 we have

$$A_0 = \frac{\kappa^2 \langle J_{X\bar{X}} F^X \overline{F}^X\rangle}{m_{3/2}}\,, \tag{42}$$

which is related to B_0 from Eq. (31) as $A_0 = B_0 + m_{3/2}$. Here we show explicit values of the MSSM soft parameters for the parameter set (23), as well as the mass spectrum of the model. The results

m_z	m_ζ	$m_{3/2}$	m_0	m_1	m_2	m_3
1.25×10^{12}	6.15×10^{11}	7.51×10^{11}	2.68×10^{11}	1.03×10^{10}	6.54×10^{9}	5.84×10^{9}

Table 2: Masses (in GeV) of inflaton, inflatino, gravitino, and MSSM sparticles derived from our model with parameter set (23).

are summarized in Table 2, where we take one-loop values of the Standard Model gauge couplings [3] at the reheating temperature 10^8 GeV (estimated below),

$$g_1 = 0.5\,, \quad g_2 = 0.59\,, \quad g_3 = 0.72\,. \tag{43}$$

As for the $U(1)_R$ gauge boson, its mass generated by the Higgs mechanism is 9.61×10^{11} GeV, close to the inflaton mass.

The parameters A_0 and B_0 are estimated as

$$A_0 = 1.6\times10^{12}\ \text{GeV}\,, \quad B_0 = 8.46\times10^{11}\ \text{GeV}\,. \tag{44}$$

For the model we consider, the inflaton ρ can perturbatively decay into the MSSM scalars, gaugini, and inflatino since their masses are smaller than $m_\rho/2$. However, the inflaton mass is smaller than two times the gravitino mass, $m_\rho < 2m_{3/2}$, which prohibits the perturbative decay of the inflaton into gravitini. As shown in [5], the total decay rate is

$$\Gamma_{\text{tot}} = 6.54\times10^{-3}\ \text{GeV}\,, \tag{45}$$

and we can estimate the reheating temperature as,

$$T_{\text{reh}} \simeq \sqrt{M_P\Gamma_{\text{tot}}} = 1.26\times10^{8}\ \text{GeV}\,. \tag{46}$$

4. Inflation by supersymetry breaking with the new Fayet-Iliopoulos term

In the previous sections, we discuss a class of minimal inflation models in supergravity that identify the inflaton with the goldstino superpartner in the presence of a gauged R-symmetry. We notice that the D-term has a constant Fayet–Iliopoulos (FI) contribution but plays no role during the inflation and can be neglected, while the pseudoscalar partner of the inflaton is absorbed by the $U(1)_R$ gauge field that becomes massive away from the origin.

In this section, we discuss the consequences of a new gauge invariant FI term proposed recently to the class of inflation models mentioned above. It turns out that the resulting D-term scalar potential provides an alternative realisation of inflation from supersymmetry breaking, driven by a D- instead of an F-term. The inflaton is still a superpartner of the goldstino which is now a gaugino within a massive vector multiplet, where again the pseudoscalar partner is absorbed by the gauge field away from the origin. In the following, we use the notation in [14] and κ is set to 1 for simplicity.

[3] We choose non-SUSY running of the couplings because SUSY breaking scale is very high in our models.

4.1 Component action of new FI term in superconformal tensor calculus

A new constant FI term was proposed recently in [3] (see also in [4]) of the form $\mathcal{L}_{\rm FI} = \xi_2\, D+$ fermions, that can be coupled to supergravity without gauging the R-symmetry. This new term contain inverse powers of some auxiliary field. It is non-singular when the D-auxiliary filed has a non vanishing vacuum expectation value (VEV). The corresponding supergravity Lagrangian can be written as:

$$\mathcal{L}_{\rm FI} = \xi_2 \left[S_0 \bar{S}_0 \frac{w^2 \bar{w}^2}{\bar{T}(w^2) T(\bar{w}^2)} (V)_D \right]_D , \tag{47}$$

where ξ_2 is a constant parameter. In the superconformal formalism, the chiral compensator field S_0, with Weyl and chiral weights $(\delta, w') = (1,1)$, has components $S_0 = (s_0, P_L \Omega_0, F_0)$. The vector multiplet $V = \left(v, \zeta, \mathcal{H}, v_\mu, \lambda, D\right)$ has weights $(0,0)$. We will use the Wess-Zumino gauge in which the first components $v = \zeta = \mathcal{H} = 0$. The multiplet w^2 with weights $(1,1)$ is given by

$$w^2 = \frac{\bar{\lambda} P_L \lambda}{S_0^2}, \qquad \bar{w}^2 = \frac{\lambda P_R \bar{\lambda}}{\bar{S}_0^2}, \tag{48}$$

where we have (in the components form)

$$\bar{\lambda} P_L \lambda = \left(\bar{\lambda} P_L \lambda \ ; \ \sqrt{2} P_L \big(-\frac{1}{2} \gamma \cdot \hat{F} + i D \big) \lambda \ ; \ 2 \bar{\lambda} P_L \not{\mathcal{D}} \lambda + \hat{F}^- \cdot \hat{F}^- - D^2 \right). \tag{49}$$

The corresponding kinetic terms in supergravity Lagrangian for the gauge multiplet are

$$\mathcal{L}_{\rm kin} = -\frac{1}{4} \left[\bar{\lambda} P_L \lambda \right]_F + \text{h.c.} \ . \tag{50}$$

The operator T ($\bar{T}$) in (47) is defined in [15] and [16], and can be used to define a chiral (antichiral) multiplet. For example, the chiral multiplet $T(\bar{w}^2)$ has weights $(2,2)$. This corresponds to the usual chiral projection operator $\bar{D}^2$ in the case of global supersymmetry. Note that we will drop the notation of h.c. and implicitly assume its presence for every $[\ \ ]_F$ term in the Lagrangian. Finally, the multiplet $(V)_D$ is a $(2,0)$ linear multiplet. Its components are given by

$$(V)_D = \left(D, \not{\mathcal{D}} \lambda, 0, \mathcal{D}^b \hat{F}_{ab}, -\not{\mathcal{D}} \not{\mathcal{D}} \lambda, -\Box^C D \right). \tag{51}$$

The component $\not{\mathcal{D}}\lambda$ and the covariant field strength $\hat{F}_{ab}$ are defined in eq. (17.1) of [14]. In our case, we have

$$\begin{aligned} \hat{F}_{ab} &= e_a^{\ \mu} e_b^{\ \nu} \left(2 \partial_{[\mu} A_{\nu]} + \bar{\psi}_{[\mu} \gamma_{\nu]} \lambda \right) , \\ \mathcal{D}_\mu \lambda &= \left(\partial_\mu - \frac{3}{2} b_\mu + \frac{1}{4} w_\mu^{ab} \gamma_{ab} - \frac{3}{2} i \gamma_* \mathcal{A}_\mu \right) \lambda - \left(\frac{1}{4} \gamma^{ab} \hat{F}_{ab} + \frac{1}{2} i \gamma_* D \right) \psi_\mu , \end{aligned} \tag{52}$$

where $e_a^{\ \mu}$ is the vierbein, with frame indices a, b and coordinate indices μ, ν. The gauge fields w_μ^{ab}, b_μ, and $\mathcal{A}_\mu$ correspond to Lorentz transformations, dilatations, and T_R symmetry of the conformal algebra respectively, while ψ_μ denotes the gravitino. The conformal d'Alembertian operator is defined by $\Box^C \equiv \eta^{ab} \mathcal{D}_a \mathcal{D}_b$.

Let us consider first the case of pure supergravity coupled to a $U(1)$ gauge multiplet with the FI term in (47). The supergravity Lagrangian can be written as

$$\mathcal{L} = -3\left[S_0\bar{S}_0\right]_D + \left[S_0^3 W_0\right]_F - \frac{1}{4}\left[\bar{\lambda}P_L\lambda\right]_F + \mathcal{L}_{\text{FI}}. \tag{53}$$

Supersymmetry is broken via a non-vanishing VEV of the D-auxiliary component of the vector multiplet driven by the linear term in D, with the Goldstino being the $U(1)$ gaugino. By fixing the compensator $S_0 = 1$, integrating out the auxiliary fields, and choosing the unitary gauge where the Goldstino vanishes, the Lagarangian in component form is

$$e^{-1}\mathcal{L} = \frac{1}{2}\left(R - \bar{\psi}_\mu\gamma^{\mu\nu\rho}D_\nu\psi_\rho + m_{3/2}\bar{\psi}_\mu\gamma^{\mu\nu}\psi_\nu\right) - \frac{1}{4}F^{\mu\nu}F_{\mu\nu} - \left(-3m_{3/2}^2 + \frac{1}{2}\xi_2^2\right), \tag{54}$$

with a constant superpotential $m_{3/2} = W_0$. In the absence of matter, any non-vanishing value of ξ_2 breaks supersymmetry and uplifts the vacuum energy by a constant term $\mathcal{V}_{FI} = \xi_2^2/2$. It is also important to note that the FI term in eq. (47) breaks the Kähler invariance and does not require the gauging of an R-symmetry.

Let us now couple the FI-term given by eq. (47) to additional matter fields charged under the $U(1)$. For simplicity, we focus on a single chiral multiplet X. The Lagrangian is given by

$$\mathcal{L} = -3\left[S_0\bar{S}_0 e^{-\frac{1}{3}\mathcal{K}(X,\bar{X})}\right]_D + \left[S_0^3\mathcal{W}(X)\right]_F - \frac{1}{4}\left[\mathcal{F}(X)\bar{\lambda}P_L\lambda\right]_F + \mathcal{L}_{\text{FI}}. \tag{55}$$

Here $\mathcal{K}(X,\bar{X})$, $\mathcal{W}(X)$ and $\mathcal{F}(X)$ are a Kähler potential, a superpotential and a gauge kinetic function respectively. The first three terms in eq. (55) are the usual supergravity Lagrangian [14]. Assuming that the multiplet X transforms under the $U(1)$ as

$$\begin{aligned} V &\longrightarrow V + i\Lambda - i\bar{\Lambda}, \\ X &\longrightarrow X e^{-iq\Lambda}, \end{aligned} \tag{56}$$

where Λ is a gauge multiplet parameter. In the case we consider, the superpotential does not transform under the gauge symmetry therefore the $U(1)$ is not an R-symmetry. For a model with a single chiral multiplet, the superpotential must be constant

$$\mathcal{W}(X) = F. \tag{57}$$

To ensure gauge invariance of the supergravity action, the Kähler potential must be a function of $Xe^{qV}\bar{X}$. However, for notational simplicity, in the following we drop the e^{qV} factors.

Indeed, in this case we can consistently add the FI-term $\mathcal{L}_{\text{FI}}$ to the theory, similar to [3], and the resulting D-term potential acquires an extra term proportional to ξ_2

$$\mathcal{V}_D = \frac{1}{2}\text{Re}\left(\mathcal{F}(X)\right)^{-1}\left(ik_X\partial_X\mathcal{K} + \xi_2 e^{\frac{1}{3}\mathcal{K}}\right)^2, \tag{58}$$

where the Killing vector is $k_X = -iqX$. For a constant superpotential (57), the F-term potential reduces to

$$\mathcal{V}_F = |F|^2 e^{\mathcal{K}(X,\bar{X})}\left(-3 + g^{X\bar{X}}\partial_X\mathcal{K}\partial_{\bar{X}}\mathcal{K}\right). \tag{59}$$

From eq. (58) it is easy to see that if the Kähler potential has a term proportional to $\xi_1 \log(X\bar{X})$, the D-term contribution to the scalar potential obtains another constant contribution. For example, if

$$\mathcal{K}(X, \bar{X}) = X\bar{X} + \xi_1 \ln(X\bar{X}), \tag{60}$$

the D-term potential becomes

$$\mathcal{V}_D = \frac{1}{2}\mathrm{Re}\,(\mathcal{F}(X))^{-1} \left(qX\bar{X} + q\xi_1 + \xi_2 e^{\frac{1}{3}\mathcal{K}}\right)^2 . \tag{61}$$

The term proportional to ξ_1 is the usual FI term in a non R-symmetric Kähler frame. It can be consistently added to the model with the new FI term proportional to ξ_2.

In the absence of the extra term, a Kähler transformation

$$\begin{aligned} \mathcal{K}(X, \bar{X}) &\to \mathcal{K}(X, \bar{X}) + \mathcal{J}(X) + \bar{\mathcal{J}}(\bar{X}), \\ \mathcal{W}(X) &\to \mathcal{W}(X) e^{-\mathcal{J}(X)}, \end{aligned} \tag{62}$$

with $\mathcal{J}(X) = -\xi_1 \ln X$ allows us to recast the model in the form

$$\begin{aligned} \mathcal{K}(X, \bar{X}) &= X\bar{X}, \\ \mathcal{W}(X) &= m_{3/2} X, \end{aligned} \tag{63}$$

where $m_{3/2} = F$. The two models result in the same Lagrangian, at least classically[4]. However, in the Kähler frame of eqs. (63) the superpotential transforms nontrivially under the gauge symmetry. As a consequence, the gauge symmetry becomes an R-symmetry.

Note that the extra term (47) violates the Kähler invariance of the theory, and the two models related by a Kähler transformation are no longer equivalent. The model written in the Kähler frame where the gauge symmetry becomes an R-symmetry in eqs. (63) can not be consistently coupled to $\mathcal{L}_{\mathrm{FI}}$. A generalized Kähler invariant FI term has been built in [17] and [19].

4.2 The scalar potential in a Non R-symmetry frame

In this section, we work in the Kähler frame where the superpotential does not transform, and take into account the two types of FI terms which were discussed in the last section. For convenience, we repeat here the Kähler potential in eq. (60) and restore the inverse reduced Planck mass κ:

$$\mathcal{K} = X\bar{X} + \kappa^{-2}\xi_1 \ln \kappa^2 X\bar{X}. \tag{64}$$

The superpotential and the gauge kinetic function are set to be constant [5]:

$$\mathcal{W} = \kappa^{-3} F, \quad \mathcal{F}(X) = 1. \tag{65}$$

[4]At the quantum level, a Kähler transformation also introduces a change in the gauge kinetic function f, see for example [18].

[5]In order to cancel the chiral anomalies [2], the gauge kinetic function gets a field-dependent correction $\propto q^2 \ln \rho$. However, the correction turns out to be very small and can be neglected below, since q is chosen to be of order of 10^{-5} or smaller.

After performing a change of the field variable $X = \rho e^{i\theta}$, the F-term contribution to the scalar potential is given by

$$\mathcal{V}_F = \frac{1}{\kappa^4} F^2 e^{\kappa^2\rho^2} (\kappa\rho)^{2b} \left[\frac{\left(b + \kappa^2\rho^2\right)^2}{(\kappa\rho)^2} - 3 \right], \tag{66}$$

and the D-term contribution is

$$\mathcal{V}_D = \frac{q^2}{2\kappa^4} \left(b + (\kappa\rho)^2 + \xi (\kappa\rho)^{\frac{2b}{3}} e^{\frac{1}{3}\kappa^2\rho^2} \right)^2. \tag{67}$$

Note that we set $b = \xi_1$ and rescaled the second FI parameter by $\xi = \xi_2/q$. We are interested in the role of the new FI-term in inflationary models driven by supersymmetry breaking.

For $F = 0$, one finds that for $\xi < -1$ and $b = 3$ the potential has a maximum at the origin, and a supersymmetric minimum. Since we set the superpotential to zero, the SUSY breaking is measured by the D-term order parameter, i.e. the Killing potential associated with the gauged U(1), which is given by

$$\mathcal{D} = i\kappa^{-2} \frac{-iqX}{W} \left(\frac{\partial W}{\partial X} + \kappa^2 \frac{\partial \mathcal{K}}{\partial X} W \right) + \kappa^{-2} q \xi (\kappa\rho)^2 e^{(\kappa\rho)^2/3}. \tag{68}$$

This enters the scalar potential as $\mathcal{V}_D = \mathcal{D}^2/2$. So, at the local maximum and during inflation $\mathcal{D}$ is of order q and supersymmetry is broken. On the other hand, at the global minimum, supersymmetry is preserved and the potential vanishes. Strictly speaking, the supersymetric minimum is not valid because the new FI term becomes singular since the D-auxiliary vanishes. Therefore a small F is required in any case.

For $F \neq 0$, the potential has still a local maximum at $\rho = 0$ for $b = 3$ and $\xi < -1$. For this choice, the derivatives of the potential have the following properties,

$$\mathcal{V}'(0) = 0, \quad \mathcal{V}''(0) = 6\kappa^{-4} q^2 (\xi + 1). \tag{69}$$

For $\xi < -1$, the extremum is a local maximum, as desired.

Let us comment on the global minimum after turning on the F-term contribution. As long as $F^2/q^2 \ll 1$, the change in the global minimum $\kappa\rho_v$ is very small, of order $O(F^2/q^2)$, The plot of this change is shown in Fig. 2.

Let us comment on super symmetry breaking in the present case $F \neq 0$, the order parameters are both the Killing potential $\mathcal{D}$ and the F-term contribution $\mathcal{F}_X$, which read

$$\mathcal{D} \propto q[3 + (\kappa\rho)^2(1 + \xi e^{(\kappa\rho)^2/3})], \quad \mathcal{F}_X \propto F(\kappa\rho)^2(3 + (\kappa\rho)^2) e^{(\kappa\rho)^2/2}, \tag{70}$$

where the F-term order parameter $\mathcal{F}_X$ is defined by

$$\mathcal{F}_X = -\frac{1}{\sqrt{2}} e^{\kappa^2 \mathcal{K}/2} \left(\frac{\partial^2 \mathcal{K}}{\partial X \partial \bar{X}} \right)^{-1/2} \left(\frac{\partial \bar{W}}{\partial \bar{X}} + \kappa^2 \frac{\partial \mathcal{K}}{\partial \bar{X}} \bar{W} \right). \tag{71}$$

Therefore, near the local maximum, $\mathcal{F}_X/\mathcal{D} \sim \frac{F}{q}\rho^2$. On the other hand, at the global minimum, both $\mathcal{D}$ and $\mathcal{F}_X$ are of the same order i.e. $\mathcal{F}_X/\mathcal{D} \sim \frac{F}{q}$, assuming that ρ at the minimum is of order $O(1)$, which is true in our models below. This makes tuning of the vacuum energy between the F- and D-contribution in principle possible.

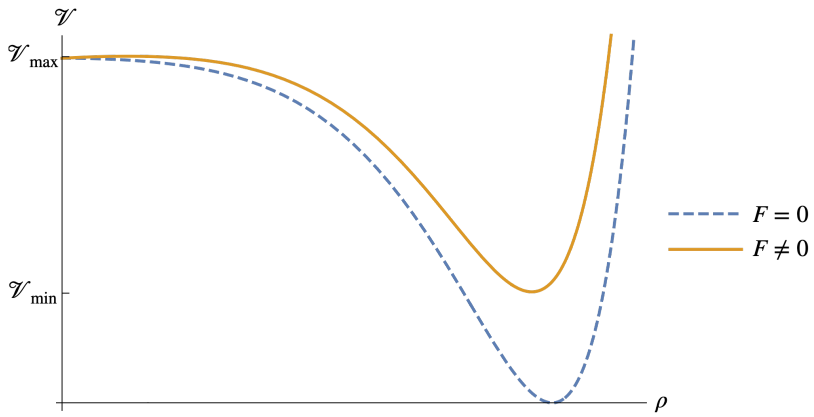

Figure 2: This plot presents the scalar potentials for $F = 0$ and $F \neq 0$ cases. For $F = 0$, we have a local maximum at $\rho = 0$ and the global minimum has zero cosmological constant. For $F \neq 0$, the origin $\rho = 0$ is still the maximum but the global minimum now has a positive cosmological constant.

Let us make a comment on the case $b = 0$ where only the new FI parameter ξ contributes to the potential. In this case, the condition for the local maximum of the scalar potential at $\rho = 0$ can be satisfied for $-3 < \xi < 0$. In the case where F is set to zero, the scalar potential (67) has a minimum at $\kappa\rho^2_{\text{min}} = 3\ln\left(-\frac{3}{\xi}\right)$. In order to have $\mathcal{V}_{\text{min}} = 0$, we can choose $\xi = -\frac{3}{e}$. However, we find that this choice of parameter ξ does not allow slow-roll inflation near the maximum of the scalar potential. Similar to the previous section, it may be possible to achieve both the scalar potential satisfying slow-roll conditions and a small cosmological constant at the minimum by adding correction terms to the Kähler potential and turning on a parameter F. However, in the next section, we will focus on $b = 3$ case where less parameters are required to satisfy the observational constraints.

4.3 An example for D-term inflation model

Let us focus on the $b = 3$ case and assume that the scalar potential is D-term dominated by fixing $F = 0$, the model has only two free parameters, namely q and ξ. The first parameter controls the overall scale of the potential and it will be fixed by the amplitude A_s of the CMB data. The only free-parameter left over is the second parameter ξ. We derive the condition that leads to slow-roll inflation scenarios, where the start of inflation (or, horizon crossing) is near the maximum of the potential at $\rho = 0$.

Since we assume inflation to start near the origin $\rho = 0$, the expansion of slow-roll parameters for small $\kappa\rho$ can be written as

$$\begin{aligned} \epsilon &= \frac{4}{9}(\xi + 1)^2(\kappa\rho)^2 + O(\rho^3), \\ \eta &= \frac{2(1+\xi)}{3} + O(\rho^2). \end{aligned} \tag{72}$$

Note also that η is negative when $\xi < -1$. We can therefore tune the parameter ξ to avoid the η-problem. The observation is that at $\xi = -1$, the effective charge of X vanishes and thus the ρ-dependence in the D-term contribution (67) becomes of quartic order.

Note that we obtain the same relation between ϵ and η as in the model of inflation from supersymmetry breaking driven by an F-term from a linear superpotential and $b = 1$ (see eq. (17) and (18)). Thus, there is a possibility to have flat plateau near the maximum that satisfies the slow-roll condition and at the same time a small cosmological constant at the minimum nearby.

The number of e-folds N during inflation is determined by using eq. (16) where we choose $|\epsilon(\chi_{\text{end}})| = 1$. Notice that the slow-roll parameters for small ρ^2 satisfy the simple relation $\epsilon = \eta(0)^2\rho^2 + O(\rho^4)$ by eq. (72). Therefore, the number of e-folds between $\rho = \rho_1$ and ρ_2 ($\rho_1 < \rho_2$) takes the following simple approximate form as in [2],

$$N \simeq \frac{1}{|\eta(0)|} \ln\left(\frac{\rho_2}{\rho_1}\right) = \frac{3}{2|\xi + 1|} \ln\left(\frac{\rho_2}{\rho_1}\right). \tag{73}$$

as long as the expansions in (72) are valid in the region $\rho_1 \leq \rho \leq \rho_2$. Note that we used the approximation $\eta(0) \simeq \eta_*$, which holds in this case.

We are now comparing the theoretical predictions of this model to the observational data via the power spectrum of scalar perturbations of the CMB such as the amplitude A_s, tilt n_s and the tensor-to-scalar ratio r. From the relation of the spectral index above, one should have $\eta_* \simeq -0.02$, and thus eq. (73) gives approximately the desired number of e-folds when the logarithm is of order one. Actually, using this formula, we can estimate the upper bound of the tensor-to-scalar ratio r and the Hubble scale H_* following the same argument given in [2]; the upper bounds are given by computing the parameters r, H_* assuming that the expansions (72) hold until the end of inflation. We then get the bound

$$r \lesssim 16(|\eta_*|\kappa\rho_{\text{end}} e^{-|\eta_*|N})^2 \simeq 10^{-4}, \quad H_* \lesssim 10^{12}\,\text{GeV}, \tag{74}$$

where we used $\eta_* = -0.02$, $N \simeq 50 - 60$ and $\rho_{\text{end}} \lesssim 0.5$, which are consistent with our models.

4.4 A small field inflation model from supergravity with observable tensor-to-scalar ratio

Supergravity models with higher r are of particular interest. In this section we show that our model can get large r at the price of introducing some additional terms in the Kähler potential. Let us consider the previous model with additional quadratic and cubic terms in $X\bar{X}$:

$$K = X\bar{X} + A\kappa^2(X\bar{X})^2 + B\kappa^4(X\bar{X})^3 + \kappa^{-2} b \ln(\kappa^2 X\bar{X}), \tag{75}$$

while the superpotential and the gauge kinetic function remain as in eq. (65). We assume that inflation is driven by the D-term by setting the parameter $F = 0$. The scalar potential in terms of the field variable ρ can be written as:

$$\mathcal{V} = \frac{q^2}{\kappa^4}\left(b + (\kappa\rho)^2 + 2A(\kappa\rho)^4 + 3B(\kappa\rho)^6 + \xi(\kappa\rho)^{\frac{2b}{3}} e^{\frac{1}{3}\left(A(\kappa\rho)^4 + B(\kappa\rho)^6 + (\kappa\rho)^2\right)}\right)^2. \tag{76}$$

We now have two additional parameters A and B. These parameters do not affect our previous discussions on the choices of the parameter b because they appear in higher orders in ρ in the scalar potential. Therefore, we can continue with the $b = 3$ case. The formula (73) for the number of e-folds also holds for small $(\kappa\rho)^2$ even when A, B are not zero because the new parameters appear at order ρ^4 and higher. However these two parameters can increase the value of the tensor-to-scalar ratio r. To obtain $r \approx 0.01$, we can choose for example

$$q = 8.68 \times 10^{-6}, \quad \xi = -1.101, \quad A = 0.176, \quad B = 0.091. \tag{77}$$

By choosing the initial condition $\kappa\rho_* = 0.445$ and $\kappa\rho_{\text{end}} = 1.155$, we get the results $N = 58$, $n_s = 0.96$, $r = 0.01$ and $\mathcal{A}_s = 2.2 \times 10^{-9}$, which is in agreement with the CMB data. Note that an application of the new FI term in no-scale supergravity model for inflation can be found for example in [19–21].

5. Conclusions

In this chapter we discussed the possibility that inflation is driven by supersymmetry breaking with the scalar component of the goldstino superfield playing the role of the inflaton. Imposing a gauged R-symmetry allows to satisfy easily the slow-roll conditions, leading to an interesting class of small field inflation models, characterised by an inflationary plateau around the maximum of the scalar potential near the origin, where R-symmetry is restored with the inflaton rolling down to a minimum with an infinitesimal tuneable positive vacuum energy. Inflation can be driven by either an F- or a new FI D-term. The above models are in agreement with cosmological observations and in the simplest case predict a rather small tensor-to-scalar ratio of primordial perturbations.

We also described the MSSM-inflaton couplings, and estimated the reheating temperature, generally $T_{\text{reh}} \sim 10^8$ GeV, from perturbative decay channels of the inflaton. In our model the inflaton can decay into all the MSSM sparticles. The inflaton mass is smaller than two times the gravitino mass, which prohibits perturbative decay of the former into two gravitini. The full picture of reheating, however, requires further investigation after taking into account non-perturbative effects such as Bose condensation and possible resonant production of fermions. Finally, as explained in [5], our minimal models do not allow for thermal LSP dark matter, but superheavy LSP dark matter (e.g. neutralino) is possible depending on the parameter choice.

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