Higgs inflation from new Kähler potential

Sayantan Choudhury∗, Trina Chakraborty† and Supratik Pal ‡

Physics and Applied Mathematics Unit, Indian Statistical Institute, 203 B.T. Road, Kolkata 700 108, India
(Dated: February 18, 2018)

We introduce a new class of models of Higgs inflation using the superconformal approach to supergravity by modifying the Kähler geometry. Using the above-mentioned mechanism, we construct a phenomenological functional form of a new Kähler potential followed by construction of various types of models which are characterized by a superconformal symmetry breaking parameter χ. Depending on the numerical values of χ we classify the proposed models into three categories. Models with minimal coupling are identified by χ = ± 2 3 branch which are made up of shift symmetry preserving flat directions. We also propose various other models by introducing a non-minimal coupling of the inflaton field to gravity described by χ ≠ ± 2 3 branch. We employ all these proposed models to study the inflationary paradigm by estimating the major cosmological observables and confront them with recent observational data from WMAP9 along with other complementary data sets, as well as independently with PLANCK. We also mention an allowed range of non-minimal couplings and the Yukawa type of couplings appearing in the proposed models used for cosmological parameter estimation.

Contents

I. Introduction

II. Superconformal mechanism in Kähler geometry

III. Inflationary model building for different values of the non-minimal coupling (χ)
   A. Models with χ = 2 3
   B. Models with χ = − 2 3
   C. Models with χ ≠ ± 2 3

IV. Summary and outlook

Acknowledgments

References

I. INTRODUCTION

Cosmological inflation has been a paradigm in which the pathological problems of the Standard Big Bang Cosmology are addressed in a sophisticated way. The inflaton field yields scale-dependent nearly Gaussian spectrum of density fluctuations. Moreover, during the inflationary epoch, cosmological perturbation via quantum fluctuation provides seed for the large-scale structure formation as we perceive today. Inflation is governed by a flat potential which has a proper field theoretic origin [1–5]. In this context, supersymmetry (SUSY) or its local extension (i.e. supergravity (SUGRA)) is the most successful candidate, which imposes certain constraints on the non-supersymmetric models of particle physics and cosmology [6–9]. A well known example of such restrictions is the fact that the supersymmetric version of the Standard Model (SM) of particle physics requires at least two Higgs superfields [4, 10]. On the other hand, the SUSY embedding of the Higgs model in inflation requires SUGRA [11–18]. Thus, it is interesting to see how SUSY may affect various inflationary models, where the gravity sector is minimally or non-minimally coupled to scalar fields [11, 19].

∗ Electronic address: sayanphysicsisi@gmail.com
† Electronic address: trinachakraborty.27@gmail.com
‡ Electronic address: supratik@isical.ac.in
A competent idea is to exercise the SM Higgs doublet as the inflaton [20, 21] through the well-known Higgs inflation [22] within the SUGRA domain [12–14, 16]. In this framework, inflation is realized via a large non-minimal coupling of Higgs doublet to Einstein gravity instead of a tiny Higgs quartic coupling, as it contradicts the observed Higgs mass bound at Large Hadron Collider (LHC) [23]. Earlier it has been shown in various works [24–27] that by applying power counting formalism, Hubble scale during inflation approaches the unitarity limit on the new scale in conjunction with the breakdown of the semi-classical approximation in the effective field theory of inflation in four dimension below the Ultra-Violet (UV) cut-off. However, a customary notion is prevalent amongst physicists for the study of effective field theory of inflation in which a singlet field with non-minimal coupling can act as an inflaton for a small singlet self interaction motivated quartic coupling. In such cases Hubble scale can be smaller than the unitarity bound. The well posed hierarchy problem in the context of SM has been resolved by implementing the well-known weak scale SUSY [10, 28], which is one of the most important topic of research in particle physics collider phenomenology. In the framework of Minimal Supersymmetric Standard Model (MSSM) [3, 10, 28, 29], there is an existence of two Higgs doublets and the equivalent self coupling can be expressed in terms of the electroweak (EW) gauge couplings. Setting apart the unitarity problem in the context of MSSM Higgs inflation, the implementation of Higgs inflation without fine tuning [30] is inconceivable due to the appearance of instability in the ratio of two Higgs VEVs. An interesting situation may emerge when the superpotential term provides the vacuum energy via the introduction of an additional self interacting coupling required for inflation governed by Next-to-Minimal Supersymmetric Standard Model (NMSSM) [1, 12, 13, 15, 31]. As this new self coupling can be made small without any violation of the recently observed LHC bound on the Higgs mass, there might be another physical possibility appearing where the Higgs inflation can be performed within the semi-classical limit of effective field theory.

However the supergravity theory has a dark side in the context of Higgs inflation. The main problem was rooted in the functional form of the Kähler potential which involves typical contributions proportional to quadratic combination of the superfields in the canonical version. One elegant way to overcome such problem is to search for shift symmetry [14, 32–34] protected flat directions in supergravity which can take part in inflation. The flatness of the potential is broken only by introducing a superconformal symmetry breaking parameter in the supergravity Kähler potential [12, 13]. Such terms are directly connected with non-minimal interactions of the inflaton field to the Einstein gravity sector. This class of non-minimal models of Kähler potential has many interesting features, which were explored in the context of superconformal approach to supergravity. Specifically, in the context of canonical superconformal supergravity (CSS) models [12, 13, 15], kinetic terms in the preferred frame of reference (Jordan frame) are canonical and the corresponding potential is exactly same as that appearing in global supersymmetry. For this purpose, in this article we propose a phenomenological model of a new Kähler potential with two singlet chiral superfields ($H, S$) which successfully address the problems of supergravity inflation with non-minimal coupling ($\xi_1, \xi_2$). Here one singlet field plays the role of inflaton and the other one is the background which will trigger preheating [35, 36] reheating [37–39] depending on the branching ratios of different decay channels of the inflaton. Our result can be applied directly to the Higgs inflation by satisfying D-flat constraints. In this article, our primary target is to do a thorough survey of inflationary models from Kähler potential using superconformal transformation followed by confrontation with latest observational data from WMAP9 [40] and other complementary datasets. The results have also been confronted independently with PLANCK data [41].

The paper is organized as follows. We first explain a general framework for $\mathcal{N}=1, D=4$ Jordan frame supergravity where superconformal symmetry breaking parameters for scalar fields are suitably implemented. Then we introduce a new phenomenological model of Kähler potential with two singlet chiral superfields. Next we discuss the implication of the Higgs inflation from various types of inflationary potentials derived from the Jordan frame Kähler potential for four distinct physical branches of the symmetry breaking parameter ($\chi$). Next imposing the constraints from LHC we employ these models for cosmological parameter estimation by using a numerical code CAMB [42]. Finally, we confront the cosmological observables with the latest available datasets.

II. SUPERCONFORMAL MECHANISM IN KähLER GEOMETRY

In this section we start our discussion with $\mathcal{N}=1, D=4$ SUGRA action in the Jordan frame with generalized frame function $\Phi (z, \bar{z})$ in the Planckian unit described by [12, 13]

$$S_{\Phi} = \int d^4x \sqrt{-g_\mathcal{J}} \left[ R(4) - 2\Lambda (4) + e^{-1}_{(4)} S_{\text{SUGRA}} \right]$$

(2.1)

where

$$e^{-1}(4) S_{\text{SUGRA}} := -\frac{\Phi(z, \bar{z})}{6} \left[ R(4) - \bar{\Psi}_\mu R^\mu \right] - \frac{1}{6} (\partial_\mu \Phi) \left( \bar{\Psi}^\alpha \gamma_\mu \Psi^\alpha \right) + L_0 + L_\Phi + L_1 + L_m + L_{mix} + L_d + L_{4f} - V_\mathcal{J}.$$  

(2.2)
In equation (2.2) the notations used are: $\Psi_{\mu} \Rightarrow$ gravitino field, $R^\mu \Rightarrow$ gravitino kinetic term, $L_0 \Rightarrow$ scalar d.o.f., $L_4 \Rightarrow$ fermion d.o.f., $L_1 \Rightarrow$ vector d.o.f., $L_m \Rightarrow$ fermion mass term, $L_{mix} \Rightarrow$ mixing term, $L_d \Rightarrow$ kinetic D term, $L_{dif} \Rightarrow$ four fermion term and the SUGRA potential in Jordan frame is given by \cite{12, 13}

$$V_J = \frac{\Phi^2(z, \bar{z})}{9} \left[ e^{K(z, \bar{z})} \left\{ (\nabla_\alpha W(z)) G^{\alpha \bar{\beta}} (\nabla_{\bar{\beta}} \bar{W}(z)) - 3|W(z)|^2 \right\} + \frac{1}{2} (\text{Re } f(z))^{-1} A B P_A P_B \right]$$ \quad (2.3)

where $\alpha = 1, 2, \ldots, n$ represents the number of complex scalars in the SUGRA chiral multiplet, $K(z, \bar{z})$ is the Kähler potential, $W(z)$ is the holomorphic superpotential, $f_{AB}(z)$ is the holomorphic kinetic gauge matrix field and the Killing potential or momentum map is denoted by $P_A$ which includes all the Yang-Mills transformation of the scalars through which Fayet-Iliopoulos terms are also taken care of. In equation (2.1) the supergravity verbien (inverse of fünfbinien) is characterized by the transformation rule \cite{43}

$$g^J_{\mu
u} := \eta_{AB} \left( V_{\mu}^A \otimes V_{\nu}^B \right)$$ \quad (2.4)

with

$$\text{Det}(V) = \sqrt{-g_J} = e_{(4)}.$$ \quad (2.5)

Here we use the following definition of covariant derivative:

$$\nabla_\alpha W := W_\alpha + K_\alpha W$$ \quad (2.6)

where the subscript $\alpha$ denotes differentiation with respect to complex field $z^\alpha$. By setting $\Phi = -3$, the SUGRA action in the Jordan frame reduces to the well known action in the Einstein frame. Consequently the potential stated in equation (2.3) can be related to its Einstein frame counterpart as

$$V_J = \frac{\Phi^2(z, \bar{z})}{9} V_E,$$ \quad (2.7)

where the subscripts $J$ and $E$ are used to denote Jordan and Einstein frame. Here both the frames are connected via the superconformal transformation defined in terms of the metric as

$$g^J_{\mu\nu} = \Omega^2(z, \bar{z}) g^E_{\mu\nu}$$ \quad (2.8)

where we identify the conformal factor with

$$\Omega^2(z, \bar{z}) = -\frac{\Phi(z, \bar{z})}{3} = e^{-\frac{K(z, \bar{z})}{3}}$$ \quad (2.9)

which yields a purely bosonic action in $N=1, D=4$ SUGRA in a specific Jordan frame triggering the superHiggs mechanism. The SUGRA action includes $SU(2, 2|1)$ superconformal symmetry, local dilation, special conformal symmetry, special SUSY and local $U(1)_R$ symmetry and other local symmetries of $N=1, D=4$ SUGRA. Such a superconformal mechanism is very useful to embed a class of scale invariant Global Supersymmetric (GSUSY) models into SUGRA theory. By “embedding”, here we actually point towards the fact that the $N=1, D=4$ self-interacting SUGRA multiplets has a local Poincare SUSY which can be obtained by the breakdown of above mentioned superconformal symmetry. Consequently the pure SUGRA sector in the action stated by equation (2.1) breaks superconformal symmetry and the matter part remains superconformal after gauge fixing. The non-canonical nature of the kinetic term is generally guaranteed by the following choice of superconformal factor \cite{12-14}:

$$\Omega^2(z, \bar{z}) = 1 - \frac{1}{3} \left\{ \delta_{\alpha\bar{\beta}} z^\alpha \bar{z}^{\bar{\beta}} + J(z) + \bar{J}(\bar{z}) \right\}$$ \quad (2.10)

where $J(z)$ and $\bar{J}(\bar{z})$ are the phenomenological holomorphic functions considered in the Kähler gauge. It is important to mention here that the dilatation symmetry implies $\Omega^2(z, \bar{z})$ to be homogeneous of first degree in both $z$ and $\bar{z}$, $W(z)$ to be homogeneous of third degree in $z$. Additionally local $U(1)_R$ symmetry implies $\Omega^2(z, \bar{z})$ is neutral and $W(z)$ has chiral weight three (which has been taken care of in equation (2.12)). We also assume that the resultant potential is obtained only from the supergravity F-term as the kinetic sector is gauge fixed by imposing the D-flat constraints.

Now using equation (2.9) and equation (2.10) one can find out the explicit expressions for SUGRA frame function and Kähler potential in this context. Using these results we obtain the following expression for Kähler metric:

$$G^{\alpha \bar{\beta}} = \left( \frac{\partial^2 \Omega^2(z, \bar{z})}{\partial z_{\alpha} \partial \bar{z}_{\bar{\beta}}} \right) = \left\{ 1 - \frac{1}{3} \left\{ \delta_{\alpha \bar{\beta}} z^{\alpha} \bar{z}_{\bar{\beta}} + J(z) + \bar{J}(\bar{z}) \right\} \right\} \left[ \delta^{\alpha \bar{\beta}} - \frac{1}{3} \left( z^{\alpha} \bar{z}_{\bar{\beta}} + \delta^{\alpha \bar{\beta}} \left( J(z) + \bar{J}(\bar{z}) \right) \right) \right]$$ \quad (2.11)
Assuming non-canonical structure of the superconformal factor stated in equation (2.10) let us prove the equivalence of F-term SUGRA potential in superconformal Jordan frame and in GSUSY. We start with a renormalizable \( N = 1, D = 4 \) SUGRA where the most generalized expression of the superpotential is constrained to the following cubic form:

\[
W(z) = \frac{1}{3} d_{\alpha \beta \gamma} z^\alpha z^\beta z^\gamma \tag{2.12}
\]

where \( d_{\alpha \beta \gamma} \)'s are the trilinear couplings in SUGRA theory. Equation (2.12) breaks the \( SU(1, n) \) symmetry. Now considering the fact that the SUGRA superpotential is homogeneous of the third degree in \( z^\alpha \)'s we get:

\[
W_{\alpha} z^\alpha = 3W,
\]

\[
\frac{\partial W}{\partial z^\alpha} = 3W_{\alpha} \tag{2.13}
\]

Considering all the above facts the Jordan frame D-flat potential turns out to be

\[
V^F_{\text{J}} = \left( 1 - \frac{1}{3} (J(z) + \bar{J}(\bar{z})) \right) [V^F_{\text{GSUSY}}(z) + W(\partial_{z^\alpha} J(z)) + W(\partial_{\bar{z}^\alpha} \bar{J}(\bar{z}))] + |W|^2 \left\{ \delta_{\alpha \beta} z^\alpha z^\beta + J(z) + \bar{J}(\bar{z}) \right. \\
- \frac{1}{3} z^\alpha z^\beta \left[ \delta_{\alpha \beta} \delta_{\alpha \gamma} z^\alpha z^\gamma + \delta_{\alpha \beta} \delta_{\beta \gamma} z^\alpha z^\gamma + \delta_{\alpha \gamma} z^\gamma + \delta_{\beta \gamma} z^\beta \right] \\
- \left. \frac{1}{3} \right\} \left( \frac{1}{2} |W|^2 \delta_{\alpha \beta} + W \delta_{\alpha \beta} \partial_{\bar{z}^\alpha} \partial_{\bar{z}^\beta} \right) \tag{2.14}
\]

where GSUSY potential \( V^F_{\text{GSUSY}}(z) = \delta_{\alpha \beta} W_{\alpha} \bar{W}_{\beta} \). Here the superscript \( F \) denotes F-term potential. Here it is important to mention that when superconformal symmetry is gauge fixed, the matter multiplets are preserved, which implies \( J(z) = 0 \) and \( \bar{J}(\bar{z}) = 0 \). Consequently equation (2.14) reduces to the following D-flat form of the effective potential:

\[
V^F_{\text{J}} = V^F_{\text{GSUSY}}(z) - \frac{1}{3} \delta_{\alpha \gamma} \delta_{\beta \alpha} z^\alpha z^\beta \bar{z}^\gamma \bar{z}^\beta |W|^2 \tag{2.15}
\]

where in the last non-renormalizable term of the above expansion the superpotential is highly suppressed by the UV cut-off scale (\( \Lambda_{\text{UV}} \)) of the effective theory of gravity in presence of \( O(1/L^2_{\text{UV}}) \) order term. Here \( \Lambda_{\text{UV}} \) is fixed at the value of reduced Planck scale \( M_{\text{Pl}} \sim 2.43 \times 10^{18} \text{GeV} \) in the Planckian unit system beyond which the theory becomes unprotective from UV end and the effective field theory prescription doesn’t hold good in our proposed setup. The contribution from the last term of Eq (2.15) originates from the quadratically Planck scale suppressed higher dimensional Kähler operators in \( N = 1 \) SUGRA theory. Most importantly, in four dimension, such Kähler corrections don’t contribute to the leading order of effective theory. Consequently below such high scale UV cut-off, renormalizability of the effective potential is automatically demanded within the effective field theory prescription and finally we have:

\[
V^F_{\text{J}} \simeq V^F_{\text{GSUSY}}(z \leq \Lambda_{\text{UV}} = M_{\text{Pl}}) \tag{2.16}
\]

leading to the equivalence of F-term potentials as claimed above. Next we will concentrate on a specific situation where the superconformal symmetry is broken via the non-minimal coupling parameter \( \chi \) with gravity. Consequently the frame function stated in equation (2.10) is modified as [12, 13]:

\[
\Omega^2(z, \bar{z}) = -|z^0|^2 + |z^\alpha|^2 - \chi \left( \Theta_{\alpha \beta} \frac{z^\alpha z^\beta \bar{z}^0}{z^0} + \Theta_{\alpha \beta} \frac{\bar{z}^\alpha \bar{z}^\beta z^0}{z^0} \right) \tag{2.17}
\]

which characterizes the non-flat moduli space geometry in SUGRA. Now gauge fixing criteria demands that in Planckian Unit system the compensator fields satisfy \( z^0 = 0 = \bar{z}^0 = \sqrt{3} \). This implies a subsequent modification in the matter part of the inverse Kähler metric of the enlarged space which can be expressed as:

\[
G^{\alpha \beta} = \delta^{\alpha \beta} - \frac{4\chi^2 \delta^{\alpha \lambda} \delta^{\beta \rho} \Theta_{\lambda \zeta} \bar{\Theta}_{\zeta} \bar{z}^\lambda \bar{z}^\rho}{[3 - \chi (\Theta_{\gamma \eta} z^\gamma z^\eta + \Theta_{\gamma \eta} \bar{z}^\gamma \bar{z}^\eta)} + 4\chi^2 \delta^{\gamma \eta} \Theta_{\gamma \zeta} \Theta_{\eta \bar{\rho} z^\zeta \bar{z}^\rho}],
\]

\[
G^{\alpha \bar{\beta}} = \frac{-3\chi (\Theta_{\gamma \eta} z^\gamma z^\eta + \Theta_{\gamma \eta} \bar{z}^\gamma \bar{z}^\eta)}{2\sqrt{3} \chi^2 \delta^{\alpha \lambda} \Theta_{\lambda \zeta} \bar{z}^\lambda} \tag{2.18}
\]
subject to the orthonormalization condition
\[ G^{\alpha\beta} G_{\alpha\beta} + G^{\alpha\beta} G_{\alpha\beta} = \delta^2. \]  
(2.19)

This will directly modify the Jordan frame potential stated in equation (2.3). In the next two sections we will discuss elaborately the cosmological consequences of such non-minimal coupling parameter in the context of superHiggs theory.

III. INFLATIONARY MODEL BUILDING FOR DIFFERENT VALUES OF THE NON-MINIMAL COUPLING (χ)

In this section we will start our discussion with a simple gauge fixed version of frame function in the presence of a superconformal symmetry breaking term (χ) in the Planckian unit:

\[
\Phi(H, S, H, S) = -3 - \frac{1}{4} \left( 1 + \frac{3\chi}{2} \right) [(H - H)^2 + (S - S)^2] + \frac{1}{4} \left( 1 - \frac{3\chi}{2} \right) [(H + H)^2 + (S + S)^2].
\]  
(3.1)

Using equation (2.9), the conformal factor turns out to be:

\[
\Omega^2(H, S, H, S) = 1 + \frac{1}{12} \left( 1 + \frac{3\chi}{2} \right) [(H - H)^2 + (S - S)^2] - \frac{1}{12} \left( 1 - \frac{3\chi}{2} \right) [(H + H)^2 + (S + S)^2].
\]  
(3.2)

Here the superHiggs sector \( H = \frac{H_1 + iH_2}{\sqrt{2}} \) and \( S = \frac{S_1 + iS_2}{\sqrt{2}} \) are complex scalar fields in the SUGRA chiral multiplet. Depending on the numerical values of \( \chi \), shift symmetry of \( H \) and \( S \) fields are preserved, which is one of the necessary tools to resolve SUGRA η problem in the context of inflation. However, the model will suffer from the well known tachyonic mass problem \([12, 13, 15, 44]\) in superHiggs theory, which can be resolved by adding higher order non-minimal quartic correction terms \( \beta_1(HH)^2 \) or \( \beta_2(SS)^2 \) in the frame function as well as in the conformal factor stated in equation (3.1) and equation (3.2) respectively. Here \( \beta_1, \beta_2 \) are two dimensional non-minimal couplings which are highly suppressed by the UV cut-off scale of the effective theory by \( O(1/M_{PL}) \) order term in Planckian unit. Once we add such corrections to the proposed model, tachyonic mass problem is resolved immediately in the next to leading order of the effective theory. But this will explicitly break the shift symmetry, the result of which is reappearance of SUGRA η problem. However, in our prescribed effective field theory setup the tachyonic mass problem will not at all appear as the VEV of the Higgs field is too small compared to the UV cut-off of the effective theory \( \Lambda_{UV} = M_{PL} \) and the scale of superHiggs inflation \( (\sqrt{\Lambda_{inf}} \sim 4.11 \times 10^{-3} M_{GUT}) \). Here we fix the VEV of the Higgs field at \( v = 1.01 \times 10^{-15} M_{PL} \sim 246 \text{ GeV} \), which sets the Higgs mass at the observed value \( m_H = 5.14 \times 10^{-17} M_{PL} \sim 125 \text{ GeV} \) by LHC [23]. The behavior of the Higgs potential for various values of the non-minimal coupling is explicitly shown in figure (1). This shows that with the increasing strengths of non-minimal coupling, the corresponding potential becomes more and more flat. In the next subsections we will study the cosmological consequences of these models in detail.

A. Models with \( \chi = \frac{2}{3} \)

In this branch the conformal factor is given by:

\[
\Omega^2(H, S, H, S) = 1 + \frac{1}{6} [(H - H)^2 + (S - S)^2]
\]  
(3.3)

which is connected to the Kähler potential via equation (2.9). In this context the following transformations

\[
H \rightarrow H + C_H, \\
S \rightarrow S + C_S
\]  
(3.4)

lead to the shift symmetry of the Kähler potential with respect to \( (H - H) \) and \( (S - S) \), provided \( C_H \) and \( C_S \) are constant shifts along real axis of \( H \) and \( S \) complex plane.

In table (1) we have listed several classes of Jordan frame and Einstein frame potentials obtained from all possible physical combinations of \( H \) and \( S \) of the superconformal transformation mentioned in equation (3.3). In this article, potentials obtained from \( H \) and \( S \) in any branch are exactly similar. So we will restrict ourselves to the \( H \) dependent models for cosmological parameter estimation. In order to confront with the recently observed Higgs at LHC, here we fix the VEV, \( v_1 = 246 \text{ GeV} \) with mass 125 GeV.
FIG. 1: Family of Higgs potentials for different numerical values of non-minimal coupling $\xi$, starting from $\xi = 0$. Inflation occurs either when the field $\phi = (H, S)$ rolls down from its large numerical values or when it rolls down from $\phi = 0$.

$$\Omega^2 = 1 - \lambda_1 S \left(H \bar{H} - \frac{v^2}{2}\right)$$

$$W = \frac{\lambda_2}{4} \left(H^2 - v^2\right)^2$$

$$V_f = \frac{\lambda_3}{4} \left(H^2 - v^2\right)^2$$

$$V_{\Lambda} = \frac{\lambda_4}{4} \left(S^2 - v^2\right)^2$$

TABLE II: Cosmological parameter estimation for observationally allowed models obtained from $\chi = \frac{4}{3}$ branch.

B. Models with $\chi = -\frac{4}{3}$

In this branch the conformal factor in equation (3.2) reduces to the following:

$$\Omega^2(H, S, \bar{H}, \bar{S}) = 1 - \frac{1}{6} \left[(H + \bar{H})^2 + (S + \bar{S})^2\right]$$

which is connected to the Kähler potential via equation (2.9). In this context the following transformations

$$H \rightarrow H + \tilde{C}_H,$$

$$S \rightarrow S + \tilde{C}_S$$

TABLE I: Jordan frame and Einstein frame potentials obtained from $\chi = \frac{2}{3}$ branch.
lead to the shift symmetry of the Kähler potential with respect to \((H + \bar{H})\) and \((S + \bar{S})\), provided \(\tilde{C}_H\) and \(\tilde{C}_S\) are constant shifts along imaginary axis of H and S complex plane.

$$\text{TABLE III: Jordan frame and Einstein frame potentials obtained from } \chi = -\frac{4}{3} \text{ branch.}$$

In table (III) we have mentioned the various classes of Jordan frame and Einstein frame potentials obtained from all possible physical combinations of H and S of the superconformal transformation mentioned in equation (3.5).

In table (II) and table (IV) we have mentioned all the cosmological parameters estimated from the observationally allowed potentials of \(\chi = \frac{2}{3}\) and \(\chi = -\frac{4}{3}\) branch respectively. From the numerical analysis, we have explicitly shown that almost all of these proposed models confront well with latest WMAP9 data combined with several complementary datasets of SPT, ACT, and h0 observations in the \(\Lambda\)CDM background and PLANCK data set as well. Next implementing the information obtained for each model from the cosmological code CAMB we estimate dark energy density \((\Omega_\Lambda)\), matter density \((\Omega_m)\) and its r.m.s. fluctuation \((\sigma_8)\) etc. Hence we plot the behavior of CMB angular power spectrum for \(TT\), \(TE\) and \(EE\) polarization obtained from \(\chi = \pm \frac{2}{3}\) branch as shown in figure (2(a)-2(c)) and figure (3(a)-3(c)) for scalar mode.

In this article our prime objective is to study the cosmological consequences of single field inflationary potentials. For such cases the fields other than inflaton (i.e. background fields) can trigger the two phenomenological scenarios: preheating and reheating. Further this will directly or indirectly affect the leptogenesis [37, 38, 45-48] and baryogenesis.
In this context, all the potentials are derived from SUGRA or from its superconformal extension. Consequently the energy scale of the potentials is around GUT scale. This directly satisfies the constraint on energy scale as $\mu_{\text{GUT}} < \Lambda_{\text{UV}}$, where $\mu_{\text{GUT}} \sim 10^{16}\text{GeV}$ is the corresponding energy scale of SUGRA and $\Lambda_{\text{UV}} = M_{\text{PL}}$ be the UV(Ultra-Violet) cut-off theory. Here all the Yukawa type couplings $(\lambda_1, \lambda_2)$ are energy scale dependent which will follow the Renormalization Group (RG) flow [51–53] via Callan-Symanzik equation. For the numerical estimation we fix the values of the Yukawa type couplings at GUT scale in the present context. Moreover, after applying RG flow from GUT to EWSB scale all of them becomes large ($\sim 2.065 \times 10^{-3}$) imposing the experimental constraints from LHC. It is a ray of hope for near future that proper bound on the self coupling is measurable in the next run of the LHC.

For further details on these aspects see [52, 53] where RG flow analysis has been discussed thoroughly. Most importantly, the very recent Higgs mass bound observed at LHC and latest observational data from WMAP9 and PLANCK have already ruled out the possibility of all the proposed inflationary potentials at the EWSB scale in the present context. Moreover, after applying RG flow the energy scale of the potentials is around GUT scale. This directly satisfies the constraint on energy scale as $\mu_{\text{GUT}} < \Lambda_{\text{UV}}$, where $\mu_{\text{GUT}} \sim 10^{16}\text{GeV}$ is the corresponding energy scale of SUGRA and $\Lambda_{\text{UV}} = M_{\text{PL}}$ be the UV(Ultra-Violet) cut-off theory. Here all the Yukawa type couplings $(\lambda_1, \lambda_2)$ are energy scale dependent which will follow the Renormalization Group (RG) flow [51–53] via Callan-Symanzik equation. For the numerical estimation we fix the values of the Yukawa type couplings at GUT scale in the present context. Moreover, after applying RG flow from GUT to EWSB scale all of them becomes large ($\sim 2.065 \times 10^{-3}$) imposing the experimental constraints from LHC. It is a ray of hope for near future that proper bound on the self coupling is measurable in the next run of the LHC.

For further details on these aspects see [52, 53] where RG flow analysis has been discussed thoroughly. Most importantly, the very recent Higgs mass bound observed at LHC and latest observational data from WMAP9 and PLANCK have already ruled out the possibility of all the proposed inflationary potentials at the EWSB scale in the absence of any symmetry breaking non-minimal coupling. In this article by thorough numerical analysis we explicitly show that without introducing any non-minimal coupling all the proposed inflationary potentials obtained from the $\chi = \pm \frac{2}{3}$ branches are observationally favored at the GUT scale. On the other hand such running in the Yukawa type of couplings induces the possibility of Primordial Black Hole (PBH) formation [4, 54, 55] depending on the running on the model dependent cosmological parameter $\alpha_{\phi}$. A very interesting fact for the inflationary model building is that the present observation from PLANCK (using WMAP9 data as a prior and the complementary data set (PLANCK lensing+CMB high $l$+BAO) [41] has predicted $\alpha_{\phi}$ and $\kappa_{S}$ to be $-0.013 \pm 0.009$ (although at 1.5$\sigma$) and $0.020^{+0.016}_{-0.015}$ respectively. Additionally for both $\chi = \pm \frac{2}{3}$ branches tensor to scalar ratio ($r$) are within the observational upper bound of PLANCK.

C. Models with $\chi \neq \pm \frac{2}{3}$

In this context the symmetry breaking parameter $\chi$ is connected with the non-minimal coupling $\xi$ present as $\frac{\xi}{2}\phi^2 R$ in the action. To explore more features from this sector we consider two physical situations given by:

$$\chi - \frac{2}{3} = 4\xi_1 \quad (3.7)$$

$$\chi + \frac{2}{3} = 4\xi_2 \quad (3.8)$$

where $\xi_1$ and $\xi_2$ are the two non-minimal couplings approaching from $\frac{2}{3}$ and $-\frac{2}{3}$ respectively. From equation (3.7) and equation (3.8) the superconformal factors can be expressed as:

$$\Omega_1^2(H, \bar{H}, S, \bar{S}) = 1 + \frac{1}{2} \left( \xi_1 + \frac{1}{3} \right) [(H - \bar{H})^2 + (S - \bar{S})^2] + \frac{\xi_1}{2} [(H + \bar{H})^2 + (S + \bar{S})^2] \quad (3.9)$$

FIG. 3: Variation of CMB angular power spectra vs multipoles $(l)$ for (a) TT, (b) TE and (c) EE mode from $\chi = -\frac{2}{3}$ branch. The statistical error bars are obtained from WMAP9 data.
\[ \Omega_2^2 (H, \bar{H}, S, \bar{S}) = 1 + \frac{\xi_2}{2} \left[ (H - \bar{H})^2 + (S - \bar{S})^2 \right] + \frac{1}{2} \left( \xi_2 - \frac{1}{3} \right) \left[ (H + \bar{H})^2 + (S + \bar{S})^2 \right] \]  

(3.10)

In table (V) and table (VII) we mention all types of inflationary potentials in Jordan frame and Einstein frame as obtained from the two possible physical branches of the superconformal transformations mentioned in equation (3.9) and equation (3.10) respectively.

| Class of models | \( \Omega_1^2 \) | \( \mathcal{V} \) | \( V_j \) | \( V_E \) |
|----------------|----------------|----------------|-----------|-----------|
| H real, S=0    | \((1 + \xi_1 H_1^2)\) | -\(\lambda_1 S \left( H \bar{H} - v_1^2 \right)\) | \( \frac{\lambda_1^2}{2} \left( H_1^2 - v_1^2 \right)^2 \) | \( \frac{\lambda_2^2}{2 (1 + \xi_1 H_1^2)} \) |
| H=0, S real    | \((1 + \xi_1 S_1^2)\) | -\(\lambda_2 H \left( S \bar{S} - v_2^2 \right)\) | \( \frac{\lambda_2^2}{2} \left( S_1^2 - v_2^2 \right)^2 \) | \( \frac{\lambda_3^2}{2 (1 + \xi_1 S_1^2)} \) |
| H complex, S=0 | \(1 - (\xi_1 + \frac{1}{3}) H_2^2 + \xi_1 H_1^2\) | -\(\lambda_1 S \left( H \bar{H} - v_2^2 \right)\) | \( \frac{\lambda_2^2}{2} \left( H_2^2 + H_1^2 - v_2^2 \right)^2 \) | \( \frac{\lambda_3^2}{2 (1 - (\xi_1 + \frac{1}{3}) H_2^2 + \xi_1 H_1^2)} \) |
| H=0, S complex | \(1 - (\xi_1 + \frac{1}{3}) S_2^2 + \xi_1 S_1^2\) | -\(\lambda_2 H \left( S \bar{S} - v_2^2 \right)\) | \( \frac{\lambda_2^2}{2} \left( S_2^2 + S_1^2 - v_2^2 \right)^2 \) | \( \frac{\lambda_3^2}{2 (1 - (\xi_1 + \frac{1}{3}) S_2^2 + \xi_1 S_1^2)} \) |

TABLE V: Jordan frame and Einstein frame potentials obtained from \( \chi - \frac{2}{3} = 4\xi_1 \) branch.

| Potential                  | Confronts with                  | Coup-lings | \( \xi_1 \) | \( P_R \) \((\times 10^{-9})\) | \( a_s \) \((\times 10^{-4})\) | \( r \) | \( \Omega_0 \) | \( \Omega_m \) | \( \sigma_8 \) | \( d_{H0} \) \(\text{Mpc}\) | \( d_{H0} \) \(\text{Mpc}\) |
|----------------------------|--------------------------------|------------|-------------|-----------------|-----------------|------|-------------|-------------|-------------|----------------|----------------|
| \( \frac{\lambda_1^2}{2} (H_1^2 - v_1^2)^2 \) | ACDM (WMAP9 + spt + \( h_0 \)) / PLANCK | 5.167      | 0.1        | 2.330          | 0.961           | -11.752 | 0.015       | 0.684       | 0.316       | 0.821           | 280.38          | 14184.8        |
| \( \frac{\lambda_2^2}{2} (H_2^2 + H_1^2 - v_2^2)^2 \) | ACDM (WMAP9 + spt + \( h_0 \)) / PLANCK | 6.789      | 0.1        | 2.310          | 0.960           | -9.94   | 0.013       | 0.684       | 0.316       | 0.816           | 280.38          | 14184.8        |
| \( \frac{\lambda_3^2}{2} (S_2^2 + S_1^2 - v_2^2)^2 \) | ACDM (WMAP9 + spt + \( h_0 \)) / PLANCK | 5.818      | -0.5       | 2.318          | 0.962           | -8.801  | 0.011       | 0.684       | 0.316       | 0.821           | 280.38          | 14184.8        |

TABLE VI: Cosmological parameter estimation from observationally allowed models obtained from \( \chi - \frac{2}{3} = 4\xi_1 \) branch.

![CMB Angular Power Spectrum](image1)

(a) TT

![CMB Angular Power Spectrum](image2)

(b) TE

![CMB Angular Power Spectrum](image3)

(c) EE

FIG. 4: Variation of CMB angular power spectra vs multipoles (l) (a) TT, (b) TE and (c) EE mode from \( \chi - \frac{2}{3} = 4\xi_1 \) branch. The statistical error bars are obtained from WMAP9 data.

Next we have mentioned all the cosmological parameters estimated from \( \chi \neq \frac{2}{3} \) (\( \chi - \frac{2}{3} = 4\xi_1 \) and \( \chi + \frac{2}{3} = 4\xi_2 \)) branches in table (VI) and table (VIII). This clearly shows non-minimal coupling \((\xi_1, \xi_2)\) dependent models confront
with latest data. We have also shown that if we allow the above mentioned non-minimal couplings along with very recent LHC Higgs mass bound and latest observational constraints, then almost all of the proposed inflationary potentials are favored starting from EWSB to GUT scale depending on the RG flow in Yukawa type coupling. Throughout the numerical analysis we have allowed both the signatures of the non-minimal coupling. We also avoided specific values of the non-minimal couplings for which divergences are appearing in the proposed potentials. During the analysis we have observed that only for $(\xi_1,\xi_2) > 0$ the first two models appearing in table(VI) and table(VIII) are in good agreement with latest observation. On the contrary for $(\xi_1,\xi_2) < 0$ only the third model fairs well with WMAP9 and PLANCK data set. Moreover, for the numerical estimations we consider only those values of the non-minimal couplings for which the proposed models are free from any poles. The behavior of tensor to scalar ratio ($r$) with respect to the scalar spectral index ($n_s$) for all class of proposed models of inflation are depicted in figure(6).

![CMB Angular Power Spectrum](image1.png)

(a) TT  
(b) TE  
(c) EE

FIG. 5: Variation of CMB angular power spectra vs multipoles (l) for (a) TT, (b) TE (c)EE mode from $\chi + \frac{2}{3} = 4\xi_2$ branch. The statistical error bars are obtained from WMAP9 data.
FIG. 6: Variation of tensor to scalar ratio \((r)\) vs scalar spectral index \((n_S)\) for the family of Higgs potentials for different numerical values of non-minimal coupling \(\xi\). The value of the non-minimal coupling increases as we go down towards the plot. This also shows \(\chi \pm \frac{2}{3} = 4\xi\) branches are more observationally favored compared to the \(\chi = \pm \frac{2}{3}\) branches.

IV. SUMMARY AND OUTLOOK

In this article we have proposed a class of supergravity motivated models to implement Higgs inflation, where the Higgs field is non-minimally coupled to gravity sector via symmetry breaking coupling \((\chi)\). We have followed the analysis by making use of superconformal techniques in the Kähler manifold. Using such tools we have introduced a phenomenological Kähler potential which preserves shift symmetry for two minimal coupling \(\chi = \pm \frac{2}{3}\) with gravity. This results in various classes of inflationary models which are made up of shift symmetry protected flat directions. We have elaborately discussed the consequences of superconformal techniques in the two preferred frame of references namely, Jordan and Einstein frames. Then we have explored the features of non-minimal coupling \((\xi_1, \xi_2)\) connected with shift symmetry breaking branch \(\chi \neq \frac{2}{3}\) in the context of Higgs inflation. Hence we have studied inflation from these proposed models by estimating the observable parameters which originates from primordial quantum fluctuation for scalar and tensor modes. We have further confronted our results with WMAP9 and various complementary dataset \((SPT, ACT, h_0)\) by using CAMB and as well as independently with PLANCK data set. Further we have compared the behavior of theoretical CMB polarization power spectra for \(TT, TE\) and \(EE\) mode obtained from all of these proposed models with observational power spectra. We have also commented on the allowed range for non-minimal couplings \((\xi_1, \xi_2)\) and phenomenological Yukawa type of couplings which are very crucial inputs in the context of inflationary model building. This, collectively, provides an exhaustive study of the class of Higgs inflation from Kähler potential and consequently, their pros and cons.

An interesting open issue in this context is to study the role of Hiesenber symmetry [56–58]. Other open issues is to study primordial black hole formation and its cosmological consequences from the running of the spectral index \((\alpha_S)\) and its running \((\kappa_S)\) as the very recent PLANCK data gives an estimation of the above mentioned indexes at \(1.5\sigma\) [41]. Moreover, the phenomenological consequences of all of these proposed models via reheating and leptogenesis are also a promising issue for future study.

Acknowledgments

SC thanks Council of Scientific and Industrial Research, India for financial support through Senior Research Fellowship (Grant No. 09/093(0132)/2010). SC also thanks Soumya Sadhukhan for useful discussions. TC thanks Council of Scientific and Industrial Research, India for financial support through Senior Research Fellowship (Grant
[1] D. H. Lyth and A. Riotto, Phys. Rept. 314 (1999) 1.
[2] L. Alabidi and D. H. Lyth, JCAP 0605(2006) 016.
[3] A. Mazumdar and J. Rocher, Phys. Rept. 497 (2011) 85.
[4] S. Choudhury and S. Pal, JCAP 04 (2012) 018.
[5] S. Choudhury, A. Mazumdar and S. Pal, JCAP 07 (2013) 041.
[6] F. L. Bezrukov and M. Shaposhnikov, Phys. Lett. B 659 (2008) 703.
[7] F. Bezrukov, D. Gorbunov and M. Shaposhnikov, JCAP 0906 (2009) 029.
[8] F. Bezrukov and M. Shaposhnikov, JHEP 0907 (2009) 089.
[9] A. De Simone, M. P. Hertzberg and F. Wilczek, Phys. Lett. B 678 (2009) 1.
[10] H. P. Nilles, Phys. Rept. (1984) 110.
[11] Martin B Einhorn and D R Timothy Jones, JHEP 1003 (2010) 026.
[12] S. Ferrara, R. Kallosh, A. Linde, A. Marrani and A. Van Proeyen, Phys. Rev. D 82 (2010) 045003.
[13] S. Ferrara, R. Kallosh, A. Linde, A. Marrani and A. Van Proeyen, Phys. Rev. D 83 (2011) 025008.
[14] R. Kallosh and A. Linde, JCAP 1011 (2010) 011.
[15] Hyun Min Lee, JCAP 1008 (2010) 003.
[16] K. Nakayama and F. Takahashi, JCAP 1102 (2011) 010.
[17] M. Arai, S. Kawai and N. Okada, Phys. Rev. D 84 (2011) 123515.
[18] M. Arai, S. Kawai and N. Okada, Phys. Rev. D 86 (2012) 063507.
[19] A. Linde, M. Noorbala and A. Westphal, JCAP 1103 (2011) 013.
[20] J.L. Cervantes-Cota and H. Dehnen, Nucl. Phys. B442 (1995) 391-412
[21] J.L. Cervantes-Cota and H. Dehnen, Phys.Rev. D51 (1995) 395-404
[22] Martin B. Einhorn and D. R. Timothy Jones, JCAP 1211 (2010) 049
[23] Large Hadron Collider collaboration, [http://lhc.web.cern.ch/lhc/](http://lhc.web.cern.ch/lhc/)
[24] C. P. Burgess, H. M. Lee and M. Trott, JHEP 1007 (2010) 007
[25] J. L. F. Barbon and J. R. Espinosa, Phys. Rev. D 79 (2009) 081302.
[26] Mark P. Hertzberg, JHEP 1011 (2010) 023.
[27] C. Germani and A. Kehagias, Phys. Rev. Lett. 105 (2010) 011302.
[28] S. P. Martin, arXiv: hep-ph/0709356.
[29] A. Mazumdar, arXiv: 1106.5408.
[30] A. Chatterjee and A. Mazumdar, JCAP 1109 (2011) 009.
[31] U. Ellwanger, C. Hugonie and A. M. Teixeira, Phys. Rept. 496 (2011) 1.
[32] S. Antusch, K. Dutta and P. M. Kostka, Phys. Lett. B 677 (2009) 221.
[33] R. Kallosh, A. Linde, K. A. Olive and T. Rube, Phys. Rev. D 84 (2011) 083519.
[34] I. Ben-Dayan and M. B. Einhorn, JCAP 1012 (2010) 002.
[35] Ling-Fei Wang, JCAP 1112 (2011) 018.
[36] R. Allahverdi, A. Ferrantelli, J. Garcia-Bellido and A. Mazumdar, Phys. Rev. D 83 (2011) 123507.
[37] S. Choudhury and S. Pal, Nucl. Phys. B 857 (2012) 85.
[38] S. Choudhury and S. Pal, Phys. Rev. D 85 (2012) 043529.
[39] R. Allahverdi, R. Brandenberger, Francis-Yan Cyr-Racine and A. Mazumdar, Ann. Rev. Nucl. Part. Sci. 60 (2010) 27.
[40] WMAP9 collaboration, [http://lambda.gsfc.nasa.gov/product/map/current](http://lambda.gsfc.nasa.gov/product/map/current).
[41] Planck Collaboration, arXiv:1303.5082, arXiv:1303.5076. For more up to date see: [http://www.esa.int/Our-Activities/Space-Science/Planck/](http://www.esa.int/Our-Activities/Space-Science/Planck/).
[42] CAMB, Online link: [http://camb.info/](http://camb.info/).
[43] S. Choudhury and S. SenGupta, JHEP 1302 (2013) 136.
[44] A. Sen, JHEP 0204 (2002) 048.
[45] C. S. Fong, E. Nardi and A. Riotto, Adv. High Energy Phys. 2012 (2012) 158303.
[46] A. Mazumdar and S. Morisi, Phys. Rev. D 86 (2012) 045031.
[47] R. Allahverdi and A. Mazumdar, JCAP 0610 (2006) 008.
[48] A. Mazumdar and A. Perez-Lorenzana, Phys. Rev. Lett. 92 (2004) 251301.
[49] R. Allahverdi, S. Hannestad, A. Jokinen, A. Mazumdar and S. Pascoli, arXiv: hep-ph/0504102.
[50] R. Allahverdi and A. Mazumdar, arXiv: hep-ph/0505050.
[51] L. H. Ryder, Quantum Field Theory, Cambridge University Press (2008).
[52] A.O. Barvinsky, A.Yu. Kamenshchik, C. Kiefer, A.A. Starobinsky and C. Steinwachs, JCAP 0912 (2009) 003
[53] A. O. Barvinsky, A. Yu. Kamenshchik, C. Kiefer, A. A. Starobinsky and C. F. Steinwachs, Eur.Phys.J. C72 (2012) 2219
[54] M. Drees and E. Erfani, JCAP 1201 (2012) 035.
[55] M. Drees and E. Erfani, JCAP 1104 (2011) 005.
[56] S. Antusch, M. Bastero-Gil, K. Dutta, S. F. King and P. M. Kostka, Phys. Lett. B 679 (2009) 428.
[57] S. Antusch, M. Bastero-Gil, K. Dutta, S. F. King, P. M. Kostka, JCAP 0901 (2009) 040.
[58] M. K. Gaillard, H. Murayama and K. A. Olive, Phys. Lett. B 355 (1995) 71.