$U(1)_R$ inspired inflation model in no-scale supergravityShaaban Khalil,¹ Ahmad Moursy,² Abhijit Kumar Saha,³ and Arunansu Sil³¹*Center for Fundamental Physics, Zewail City of Science and Technology,
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We consider a cosmological inflation scenario based on a no-scale supergravity sector with $U(1)_R$ symmetry. It is shown that a tree level $U(1)_R$ symmetric superpotential alone does not lead to a slowly rolling scalar potential. A deformation of this tree level superpotential by including an explicit R symmetry breaking term beyond the renormalizable level is proposed. The resulting potential is found to be similar (but not exactly the same) to the one in the Starobinsky inflation model. We emphasize that for successful inflation, with the scalar spectral index $n_s \sim 0.96$ and the tensor-to-scalar ratio $r < 0.08$, a correlation between the mass parameters in the superpotential and the vacuum expectation value of the modulus field T in the Kähler potential must be adopted.

DOI: [10.1103/PhysRevD.99.095022](https://doi.org/10.1103/PhysRevD.99.095022)**I. INTRODUCTION**

Planck satellite's four year data of the cosmic microwave background radiation and the large structure in the Universe support the predictions of cosmological inflation. The recent data confirmed that spectral index (scalar density fluctuations) is given by $n_s = 0.96 \pm 0.007$, and that the upper bound on the tensor-to-scalar ratio is $r < 0.08$ [1,2]. These results imposed severe challenges on several inflationary models. For example, the simple chaotic and hybrid inflationary models [3] are now ruled out. On the other hand, some other models of inflation with compatible cosmological fluctuation predictions have been receiving growing interest. One of these models is Starobinsky inflation [4], which is based on modified gravity.

A supergravity (SUGRA) realization of Starobinsky inflation was studied in Ref. [5] by considering a no-scale Kähler potential involving a modulus field T . It is well known that the no-scale SUGRA framework is free of the so-called η problem due to the involvement of logarithmic form in the Kähler potential. In Ref. [5], the no-scale Kähler potential of the T field is combined with a Wess-Zumino superpotential consists of a quadratic and a cubic term of the inflaton superfield S : $W = \mu S^2 - \kappa S^3$, with μ as a parameter of mass dimension and κ as a dimensionless parameter. It turns out that, at a specific point of the

parameter space, this construction becomes conformal equivalent to $R + \alpha R^2$ modified gravity models similar to the Starobinsky inflation model. Adding a term linear in field S to this renormalizable superpotential, Romao and King [6] have shown that it is also possible to realize supersymmetry breaking at the end of inflation. It is further indicated that a successful inflation consistent with correct n_s and r values may indicate an upper bound on gravitino mass once the Starobinsky limit is implemented. Few other studies having different kinds of motivation involving a Starobinsky type inflation model can be found in Refs. [7–19].

While the constructions in Refs. [5,6] are certainly elegant and minimal from their own perspectives, we notice that it is not possible to define an R charge for the superfield S so that superpotential W can have an R charge of two units. Hence no $U(1)_R$ symmetry is prevailing in this construction. Now it is well known that R symmetry plays important roles in many supersymmetric constructions. One such example is related to supersymmetry breaking. According to the Nelson-Seiberg theorem [20], existence of an R symmetry is a necessary condition in order to realize supersymmetry breaking. However, an exact R symmetry forbids gauginos and Higgsinos from having mass. Hence it must be broken (spontaneously or explicitly). It is customary to break R symmetry spontaneously, as is done in many dynamical supersymmetry breaking models leading to R axions [21].

In this paper, we start with a $U(1)_R$ global symmetry. We assume that the inflaton superfield S has an R charge unity. Thus, the tree level superpotential is given by $W = \frac{\mu}{2} S^2$, with μ as a mass scale. As we will show below, this tree level superpotential does not lead to a slowly rolling scalar

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potential. We propose a deformation of this tree level superpotential (having R charge 2) by including an explicit R symmetry breaking term beyond the renormalizable level. This new term is naturally expected to be suppressed by the cutoff scale M_* , and hence W can be expressed as

$$W = \frac{1}{2}\mu S^2 - \frac{1}{4}\lambda \frac{S^4}{M_*}, \quad (1)$$

where λ is a dimensionless coupling and M_* is a mass scale. Since a global symmetry is expected to be broken by gravity effects, a natural choice of M_* would be the Planck scale (M_P), $M_* \sim M_P$. A similar M_P suppressed R symmetry breaking term has been considered in a supersymmetric hybrid inflation scenario [22] with minimal Kähler potential. It was emphasized there that in order to get n_s within the preferred range, λ must be small enough ($\lesssim 10^{-7}$). Below we study the above superpotential in Eq. (1) and discuss how inflation can be realized in this framework.

The paper is organized as follows. In Sec. II, we study the associated inflation model originating from an interplay between a tree level $U(1)_R$ symmetric superpotential and a higher order explicit R symmetry breaking term along with no-scale Kähler potential. In Sec. III, inflationary predictions are discussed, in particular, the correlation between the spectral index n_s and the ratio r . Finally, our concluding remarks are given in Sec. IV.

II. THE MODEL

In addition to the superpotential W , we consider the Kähler potential (as standard in no-scale supergravity)

$$K = -3M_P^2 \ln \left[\frac{T + T^*}{M_P} - \frac{|S|^2}{3M_P^2} \right], \quad (2)$$

where T is the modulus field. The Kähler potential remains invariant under $U(1)_R$ symmetry with vanishing R charge for the modulus field. The supergravity potential can be obtained using

$$V_F = e^{\frac{K}{M_P^2}} \left[(D_j W^*) (K^{-1})^j_i (D^i W) - \frac{3|W|^2}{M_P^2} \right], \quad (3)$$

where

$$D^i W = \frac{\partial W}{\partial \phi_i} + \frac{W}{M_P^2} \frac{\partial K}{\partial \phi_i}, \quad D_j W^* = \frac{\partial W^*}{\partial \phi^{*j}} + \frac{W^*}{M_P^2} \frac{\partial K}{\partial \phi^{*j}},$$

$$K_i^j = \frac{\partial^2 K}{\partial \phi^{*i} \partial \phi_j},$$

where i, j refer to the modulus T and inflaton S . Now using the superpotential in Eq. (1) and the Kähler potential in Eq. (2), V_F can be obtained as

$$V_F = \frac{1}{\left(\frac{T+T^*}{M_P} - \frac{|S|^2}{3M_P^2} \right)^2} \left| \frac{\partial W}{\partial S} \right|^2. \quad (4)$$

This is a feature of no-scale supergravity that leaves the potential V_F independent of T (apart from the dependence through the prefactor e^{K/M_P^2}) and positive definite. Therefore, it can be an appropriate framework for inflationary scenarios.

Following Ref. [5], we assume here the modulus field T is stabilized at a fixed scale such that $\langle T + T^* \rangle = c$. This stabilization requires a nonperturbative effect at a high scale [23,24]. With this assumption, the effective Lagrangian turns out to be

$$\mathcal{L}_{\text{eff}} = \mathcal{L}_{\text{K.E}} - V_F = \frac{c/M_P}{\left(\frac{c}{M_P} - \frac{|S|^2}{3M_P^2} \right)^2} (\partial_\mu S^*)(\partial^\mu S) - \frac{1}{\left(\frac{c}{M_P} - \frac{|S|^2}{3M_P^2} \right)^2} \left| \mu S - \lambda \frac{S^3}{M_*} \right|^2. \quad (5)$$

In order to have the kinetic term for the complex scalar field S as a canonically normalized one, following the prescription of Ref. [5], we first redefine the S field in terms of χ ,

$$S = \sqrt{3cM_P} \tanh \left(\frac{\chi}{M_P \sqrt{3}} \right). \quad (6)$$

With the above definition of S and considering $\chi = (\chi_1 + i\chi_2)/\sqrt{2}$, the kinetic term ($\mathcal{L}_{\text{K.E.}}$) becomes

$$\mathcal{L}_{\text{K.E.}} = \sec^2 \left(\frac{2\chi_2}{\sqrt{3}M_P} \right) (\partial_\mu \chi^*)(\partial^\mu \chi), \quad (7)$$

and the F-term scalar potential responsible for inflation will have the form

$$V_F = 3M_P^4 \left(\frac{\mu^2}{M_P^2} \right) \left(\frac{M_P}{c} \right) \left[1 - \left| \tanh \left(\frac{\chi_1 + i\chi_2}{\sqrt{6}M_P} \right) \right|^2 \right]^{-2} \times \left| \tanh \left(\frac{\chi_1 + i\chi_2}{\sqrt{6}M_P} \right) - 3 \frac{c\lambda M_P}{\mu M_*} \tanh^3 \left(\frac{\chi_1 + i\chi_2}{\sqrt{6}M_P} \right) \right|^2. \quad (8)$$

From this potential, one can show that the field dependent mass squared of the imaginary component of χ , obtained by the second derivative of V_F with respect to χ_2 at the minimum $\chi_2 = 0$, is much larger than the Hubble scale squared during inflation (a numerical estimate will be provided in the next section). Therefore, imaginary part χ_2 will be stabilized at zero during the inflation. Hence we set χ_2 to be zero from now on and identify the F-term potential with the inflation potential, V_{Inf} . Note that with

FIG. 1. Schematic plots of the potential V_{Inf}/A against χ_1 for different values of B (ξ) around $B = 1$ ($\xi = 0$).

this choice, the kinetic term of the Lagrangian [see Eq. (7)] becomes canonical.

In this case, the inflation potential takes the form

$$V_{\text{Inf}} = A \cosh^4 \left(\frac{\chi_1}{\sqrt{6}} \right) \tanh^2 \left(\frac{\chi_1}{\sqrt{6}} \right) \left[1 - B \tanh^2 \left(\frac{\chi_1}{\sqrt{6}} \right) \right]^2, \quad (9)$$

where $A = 3 \frac{\mu^2}{c M_P}$ and $B = \lambda \frac{3 c M_P}{\mu M_s}$ are two dimensionless constants. In the last expression of V_{Inf} , we have set $M_P = 1$ unit. Unless otherwise stated, we will use this unit for the rest of our discussion. Note that when $B = 1$, this potential simplifies to the form

$$V_{\text{Inf}}^{(B=1)} = A \tanh^2 \left(\frac{\chi_1}{\sqrt{6}} \right). \quad (10)$$

We have shown the form of this potential (normalized by A) in Fig. 1 for different choices of B . With $B = 1$, the shape of $V_{\text{Inf}}^{(B=1)}$ (denoted by the brown line in Fig. 1) turns out to be similar to the standard Starobinsky potential [4]. If we reduce the value of B from 1 by a tiny amount, the potential starts to become steep. On the other hand, if we enhance B from 1, another distant minimum appears (other than at $\chi_1 = 0$) at some very large value of the field χ_1 .

The inference of the above discussion is that (i) the field χ_1 can now be identified as the inflaton in the limit $\chi_2 \rightarrow 0$, and (ii) B is very close to 1 as the required flatness for inflation is obtainable from the associated potential V_{Inf} of Eq. (9). In order to show the importance of the R symmetry breaking term λ , we include a plot of the potential against χ_1 , with $\lambda = 0$ (i.e., $B = 0$) in Fig. 1 denoted by the red curve. It is evident that such a potential can not provide sufficient inflation. Hence inclusion of an explicit R symmetry breaking term becomes instrumental in realizing inflation, and that too by a restricted amount. From the

nature of the plots, it is expected that, for any large deviation of B from unity, the slow roll of the inflaton might be spoiled.

In order to have better control over different values of B , we parametrize the deviation of B from 1 by ξ , $B = 1 - \xi$. Then the inflation potential in Eq. (9) can be expanded for small ξ as

$$V_{\text{Inf}} \simeq A \tanh \left(\frac{\chi_1}{\sqrt{6}} \right)^2 \left\{ 1 + 2\xi \sinh \left(\frac{\chi_1}{\sqrt{6}} \right)^2 \right\}. \quad (11)$$

In $M_P = 1$ unit, the slow-roll parameters are given by

$$\epsilon = \frac{1}{2} \left(\frac{V'_{\text{Inf}}}{V_{\text{Inf}}} \right)^2, \quad \eta = \frac{V''_{\text{Inf}}}{V_{\text{Inf}}}. \quad (12)$$

The number of e -folds is written as

$$N_e = \int_{x_e}^{x^*} \frac{1}{\sqrt{2\epsilon}} dx, \quad (13)$$

where x_e is the inflaton field value at which inflation ends, and x^* corresponds to the crossing horizon value of the inflaton. The three inflationary observables, tensor-to-scalar ratio (r), spectral index (n_s), and power spectrum (P_s), are provided by

$$r \simeq 16\epsilon, \quad (14)$$

$$n_s \simeq 1 - 6\epsilon + 2\eta, \quad (15)$$

$$P_s = \frac{V_{\text{Inf}}}{24\pi^2 \epsilon}. \quad (16)$$

These observables are to be determined at $x = x^*$.

III. INFLATIONARY PREDICTIONS

Let us now proceed to determine the inflationary observables in this scenario. The inflationary potential in Eq. (11) contains two free parameters, A and ξ . Between them, ξ takes part in determining r and n_s . The other parameter, A , will be fixed by the observed value of scalar perturbation spectrum $P_s = 2.2 \times 10^{-9}$. In Fig. 2, we show the logarithmic plot of the spectral index n_s vs the tensor-to-scalar ratio r , as predicted by our model. We also use the Planck limits for comparison purpose. The brown, black, and green curves represent the predictions for n_s and r corresponding to values of $B = 1, 0.99995$ and 0.9999 , respectively. Similarly the blue and orange lines are for $B = 1.00005$ and 1.0001 , respectively. The color codes are in accordance with Fig. 1. In this plot, a single colored line segment represents the variation of the number of e -foldings N_e from 50 to 60, where the prediction for $N_e = 55$ is denoted by a black dot over the respective line. Note that, due to the presence of another minimum at a

FIG. 2. A logarithmic plot (for vertical axis only) for the $n_s - r$ predictions of the inflationary potential through Eq. (9). Here scanning over $N_e = 50$ – 60 for a fixed B value results in one segment, on top of which a black dot is present denoting the prediction for $N_e = 55$ only. The color codes are in accordance with Fig. 1, representing different values of parameter $B(\xi)$. *Planck* observation regions, the dark and light blue regions, correspond to 1σ and 2σ contours of the *Planck* 2018 (TT, TE, EE + LowE + Lensing + BK14) data [1].

large field value for χ_1 in case with $B > 1$ (see Fig. 1), an initial condition on the field value of the inflaton has to be set. For such choices of $B > 1$, we assume that near the onset of inflation, the inflaton starts with a not so large field value; rather it was close enough to the flat part (i.e., near maximum) of the potential $\lesssim 8$ (in M_P unit). Then it can slowly roll toward the minimum at origin, and inflation can be realized.

Here we tabulate a few reference points which provide correct values of n_s and r within the allowed range of the *Planck* limit, considering $N_e = 55$. Values of A are fixed from the value of the required power spectrum. Note that the parameters A and ξ in Table I are simply combinations of the original variables: c , μ , λ , and M_* . All of these variables serve significant importance from the model point of view. Therefore, we should also estimate their magnitude in the setup. For this purpose, we consider $M_* = M_P$, argued as the natural choice in the Introduction.

Corresponding to the reference points (I–V) in Table I, in Table II we provide values of c and μ in $M_P = 1$ units for different values of λ . Note that, with λ being the parameter

TABLE I. Inflationary predictions ($n_s - r$) for five reference points in our setup, considering $N_e = 55$. These points are also indicated by black dots in Fig. 2.

Sl no.	A	ξ	x^*	n_s	r
I	1.3×10^{-10}	0	6.13	0.963 87	0.0038
II	1.55×10^{-10}	5×10^{-5}	6.20	0.970 66	0.0046
III	1.85×10^{-10}	1×10^{-4}	6.25	0.977 95	0.0056
IV	1.07×10^{-10}	-5×10^{-5}	6.07	0.957 35	0.0032
V	0.9×10^{-10}	-1×10^{-4}	6.02	0.951 56	0.0026

TABLE II. Values of c and μ (in M_P units) with $M_* = M_P$ for five reference points from Table I.

Sl. no.	λ	c	μ
I	10^{-4}	4.81481×10^{-4}	1.44444×10^{-7}
	(10^{-6})	(4.81481)	(1.44444×10^{-5})
II	10^{-4}	5.74017×10^{-4}	1.72214×10^{-7}
	(10^{-6})	(5.74017)	(1.72214×10^{-5})
III	10^{-4}	6.85048×10^{-4}	2.05535×10^{-7}
	(10^{-6})	(6.85048)	(2.05535×10^{-5})
IV	10^{-4}	3.96336×10^{-4}	1.18895×10^{-7}
	(10^{-6})	(3.96336)	(1.18895×10^{-5})
V	10^{-4}	3.3340×10^{-4}	1.0001×10^{-7}
	(10^{-6})	(3.3340)	(1.0001×10^{-5})

associated with the explicit R -symmetry breaking term, it is expected to be small. Hence in obtaining c and μ values, we have kept $\lambda \ll 1$. It can be noted from Table II that there exists a correlation between the two mass parameters c and μ for different values of λ . For example, in reference point I of Table II with $\lambda \sim 10^{-4}$, the values of c and μ are found to be $\sim 4.81 \times 10^{-4}$ and $\sim 1.44 \times 10^{-7}$, respectively. For a comparatively smaller value of $\lambda \sim 10^{-6}$, magnitudes of c and μ become $\mathcal{O}(1)$ and $\mathcal{O}(10^{-5})$, respectively. This can be interpreted by looking at the expressions of A and B which involve all of the parameters $c, \mu, \lambda, M_* = M_P$, and keeping in mind that in order to achieve successful inflation, we have to have a $B = 3\lambda c/\mu$ value very close to unity. Therefore, with a fixed choice of λ , the ratio c/μ is uniquely fixed. Then the parameter $A = 3\mu^2/c$ will fix the value of μ from the requirement that the power spectrum $P_s \sim \mathcal{O}(10^{-9})$.

It is also important to discuss the value of the modulus field vacuum expectation value (VEV) ($c = \langle T + T^* \rangle$) in terms of high scale dynamics. In fact the Kachru-Kallosh-Linde-Trivedi (KKLT) scenario assumes [25,26] that the volume modulus field can be stabilized by nonperturbative corrections to the superpotential that arise from instanton effects or gaugino condensation. Considering a single modulus as in the KKLT model [25], the nonperturbative part of the superpotential $W_{np} \sim e^{-aT}$, where a is a positive constant. It is also emphasized in Ref. [25] that the condition $aT \gg 1$ should be ensured in order to have control over supergravity approximation. Now depending on the magnitude of a , the VEV of the modulus field T could be bigger or smaller than 1 (in $M_P = 1$ units). In case $\langle T \rangle$ is sub-Planckian, a must be greater than 1 (to maintain the condition $aT \gg 1$). Such a case is discussed in Ref. [27], where it is shown that $a > 1$ can be realized through the choice of a hidden sector gauge group in order to perform the gaugino condensation. On the other hand, to establish $\langle T \rangle \gtrsim 1$, a could be either bigger or smaller than 1. In our scenario, both of the cases regarding the modulus VEV ($c < 1$ or > 1) can be accommodated as displayed in

Table II, depending on the magnitude of λ . For example, $\lambda \sim 10^{-4}$ can make the modulus VEV $c \sim \mathcal{O}(10^{-4})$, while for $\lambda \sim 10^{-6}$, c could be $\mathcal{O}(1)$ or more.

For all of the reference points mentioned in Table I, numerically it is found that the mass of the χ_2 field during inflation is significantly higher than the Hubble scale (H) during inflation given by $H^2 = V_{\text{Inf}}/3$. In particular, one finds $m_{\chi_2}^2/H^2 \simeq 4$ at the minimum of χ_2 ($\chi_2 = 0$). Hence χ_2 would be stabilized at origin during inflation. Furthermore we have also found numerically that the slope along the χ_2 direction is much steeper than the one for χ_1 . Hence χ_2 will move faster and reaches the minimum much earlier than χ_1 . This justifies our assumption that $\chi_2 = 0$ during inflation. Inflaton mass (m_{χ_1}) at its minimum for the abovementioned points is $\mathcal{O}(10^{13})$ GeV, as expected for this type of inflation scenario. We end this section by observing that, even if χ_1 is super-Planckian as required by the slow-roll condition, the field S remains sub-Planckian, as seen in Eq. (6). Now S being sub-Planckian, higher order $U(1)_R$ breaking terms in W are accordingly less important.

IV. CONCLUSION

In this paper, we propose a global R symmetry motivated inflation model within no-scale SUGRA. We find that the minimal $U(1)_R$ symmetric superpotential (quadratic in inflaton superfield) is unable to provide a successful inflation, as the associated scalar potential turns out to be extremely steep. Then we introduce an explicit R symmetry breaking term in the superpotential at a

nonrenormalizable level, which provides the required flatness for inflation. The introduction of such a $U(1)_R$ breaking term is motivated by the fact the any global symmetry will be broken by the gravity effect. For this reason, we associate the cutoff scale of this nonrenormalizable term with M_P . The effective inflation potential resulting from our proposed setup carries similarity with Starobinsky-like inflation models in the limit, one combination of parameters of the superpotential and no-scale Kahler potential as $B = 1$. Varying B from unity by a tiny amount leads to the predictions for the spectral index and tensor-to-scalar ratio. In order to keep these predictions within the limit allowed by *Planck* data, we evaluate the magnitudes of the relevant mass parameters of the model. Such a construction involving an explicit R symmetry breaking term may also have some interesting consequences, while supersymmetry breaking will also be involved. Since any dynamical supersymmetry breaking model requires that R symmetry should spontaneously be broken, leading to the presence of an R axion; such an explicit breaking term, connected with inflation, in our setup can be helpful in providing the mass of it.

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