Dynamically induced topological inflation

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(Received 25 April 2022; accepted 19 January 2023; published 6 February 2023)

We propose an inflation model in which the inflationary era is driven by the strong dynamics of $Sp(2)$ gauge theory. The quark condensation in the confined phase of $Sp(2)$ gauge theory generates the inflaton potential comparable to the energy of the thermal bath at the time of phase transition. Afterwards, with a super-Planckian field value at the global minimum, the inflation commences at a false vacuum region lying between true vacuum regions and hence the name "topological inflation." Featured by the huge separation between the scale of the false vacuum ($V(0)^{1/4} \sim 10^{15}$ GeV) and the field value at the global minimum ($\phi_{\rm min} \sim M_P$), the model can be consistent with cosmic microwave background (CMB) observables without suffering from the initial condition problems. Crucially, this is achieved only with some mild tuning of parameters in $V(\phi)$. In addition to $Sp(2)$, this model is based on an anomaly free Z_{6R} discrete R symmetry. Remarkably, while all parameters are fixed by CMB observations, the model predicts a hierarchy of energy scales including the inflation scale, supersymmetry-breaking scale, R-symmetry breaking scale, Higgsino mass and the right-handed neutrino mass given in terms of the dynamical scale of $Sp(2)$.

DOI: 10.1103/PhysRevD.107.036005

I. INTRODUCTION

The cosmic inflation has been the firmly established solution to the main cosmological issues including the flatness problem and the horizon problem [1–4]. The requirement for the de Sitter period and its end leads to the slow-roll inflation models [3,4]. In these models, starting from the origin in the field space, the inflaton field (ϕ) goes through the slow-rolling and eventually arrives at the field value at the global minimum, which results in the end of inflation.

The "new inflation" is an important type of slow-roll inflation with a Mexican-hat-like potential and a small initial inflaton field value near the origin [3,4]. It is the first inflation scenario that solves the "graceful exit" problem in the original version of inflation. However, there are problems bothering the new inflation type models. Soon after they were proposed, it was pointed out that the potential height ($V(0)^{1/4}$) needs to be tuned to be much smaller than the global minimum ($\phi_{\rm min}$) in order to obtain a small enough quantum fluctuation to be consistent with observations [5–8]. The cosmic microwave background (CMB) observations prefer a plateaulike inflaton potential with a small potential height (in Planck units) [9]. The width the plateau is required to large enough to avoid the problems posed by the initial kinetic and gradient energies of the inflaton field [9]. These require a new inflation scenario with a noticeable hierarchy between the scale of the false vacuum ($V(0)^{1/4} \ll M_{\rm Pl}$) and the field value at the global minimum ($\phi_{\rm min} \sim M_{\rm Pl}$). This hierarchy also concerns the initial location of the inflaton ($\phi_{\rm ini}$). Taking a Coleman-Weinberg form of the inflaton potential, a small coupling constant is required to achieve a small potential height. But, a weakly coupled inflaton also means it was not in thermal equilibrium with the thermal bath at early times. The lack of "supercooling" then puts in question the initial condition $\psi \sim 0$ for the new inflation; see [10] for a

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brief review. Resolving these problems usually requires extreme fine-tunings of the model parameters and/or the initial conditions.

In this work, given the aforementioned problems of inflation models with a small initial field value,[1] we propose an inflation model which establishes the separation of scales of $V(0)^{1/4}$ and $\phi_{\rm min}$ with all $\mathcal{O}(1)$ dimensionless parameters. To this end, we consider the supersymmetric $Sp(2)$ gauge theory with $N_F = 6$ flavors of quarks Q_i transforming as the fundamental representation. With an integer R-charge of Q_i, Z_{6R} can be easily shown to be free of the mixed anomaly $Z_{6R} - [Sp(2)]^2$ and thus we choose Z_{6R} as the R-symmetry of the theory. On top of this, because the discrete Z_6 is also free of the mixed anomaly $Z_6 - [Sp(2)]^2$ within $Sp(2)$ gauge theory, we introduce Z_6 as a gauge symmetry of the theory. We note that $SP(n)$ is needed to suitable effective superpotential. $n = 2$ is chosen to have Z_6R which is the unique anomaly-free discrete R symmetry in the SUSY standard model. Once the $Sp(2)$ gauge theory enters the confinement phase at the dynamical scale Λ_*, the quark condensate $\langle QQ\rangle$ forms to produce $V(\phi)$ [12]. As such, the scale of $\langle QQ\rangle \sim \Lambda_*^2$ determines the inflation scale. We shall show that how the presence of Z_6 resolves the extreme fine-tuning of parameters in $V(\phi)$ which the similar setup without Z_6 suffers from [13,14]. In this way, we trade the extreme parameter fine-tuning problem with the additional symmetry, which is nonetheless nontrivial.

Quark fields Q_i being charged under Z_{6R}, its condensation in the nonperturbative regime of $Sp(2)$ drives the spontaneous breaking of Z_{6R} to Z_{2R}. However, in our model, there is no domain wall problem associated with the discrete R-symmetry since the domain walls are diluted away by the inflation. The model gains more attractiveness when we specify the role of quark condensate more than generating the inflaton potential: as a spurion field arising from the spontaneous breaking of Z_{6R}, the quark condensate explains the dimensionful parameters in the Minimal Supersymmetric Standard Model (MSSM). This interesting attribute of the model unifies the origin of various energy scales including inflation scale ($H_{\rm inf}$), SUSY-breaking scale ($\sqrt{F_Z}$), R-symmetry breaking scale, Higgsino mass (μ_H) and the right-handed (RH) neutrino mass (m_N) and explains those based on a single energy scale Λ_* inferred from the cosmic microwave background (CMB) observable.

In this paper, we invoke the Planck unit in which the reduced Planck scale is set to the unity, i.e., $M_P = (8\pi G)^{-1/2} = 1$ unless the unit is specified.

[1]By inflation models with a small initial field value, we mean the models where $\phi_{\rm ini}$ starts the slow-roll near the false vacuum of $V(\phi)$ which is taken to be located at $\phi = 0$. Usually a PT from a symmetry breaking precedes the inflationary era, generating $V(\phi)$ for inflation. The "new inflation" [3,4] and the "natural inflation" [11] are the well-known examples of the small field inflation scenario.

II. MODEL

On top of the MSSM gauge group, we introduce

$$G = Sp(2) \otimes Z_{6R} \otimes Z_6. \tag{1}$$

as the additional symmetry group. The matter contents we assume are shown in Table. I. In addition to the MSSM particle contents[2] and the right-handed neutrino (N), $Sp(2)$ quark chiral multiplet Q_i, the singlet antisymmetric field $S_{ij} = -S_{ji}$, the inflaton chiral multiplet Φ, and the SUSY-breaking field Z are newly introduced. The indices of Q_i and S_{ij} run from 1 to $N_F = 6$, which makes the mixed anomaly of $Z_{6R} - [Sp(2)]^2$ vanish with account taken of the contribution from the $Sp(2)$ gaugino [15–17]. Thanks to this, the symmetry group $Sp(2) \otimes Z_{6R}$ is the gauged one.[3] With Q_i assigned the same Z_6 charge as the Z_{6R} charge, it is clear that Z_6 is free of anomaly with respect to $Sp(2)$ within $Sp(2)$ sector. But since Z_6 is anomalous with respect to $SU(2)_L$, for arguing the gauged Z_6, we need additional fields contributing to $Z_{6R} - [SU(2)_L]^2$. We will get back to this point later.

For a high enough energy scale where $Sp(2)$ is in its perturbative regime, the symmetry group in Eq. (1) allows for the superpotential

$$\begin{aligned} W \supset &-\lambda_{ij} S_{ij} Q_i Q_j + \lambda_{ij} g S_{ij} Q_i Q_j \Phi^2 \\ &+ \lambda_{H,ijk\ell} Q_i Q_j Q_k Q_\ell H_u H_d + \lambda_{N,ij} Q_i Q_j N N \\ &+ \lambda_{Z,ijk\ell} Q_i Q_j Q_k Q_\ell Z. \end{aligned} \tag{2}$$

where all λ's and g are $\mathcal{O}(1)$ dimensionless coupling constants.[4] The superpotential in the first line of Eq. (2) is responsible for the inflationary dynamics, the second line for the Higgsino mass term (a.k.a μ-term), the third line for the right-handed neutrino mass, and the final line for the supersymmetry (SUSY)-breaking. As we shall see shortly, in the nonperturbative regime of $Sp(2)$, the quark condensate from the quantum moduli constraint $\mathrm{Pf}[QQ] = \Lambda_*^6$ [12] and its powers generate dimensionful parameters, explaining various energy scales in the theory at the tree-level.

As the $Sp(2)$ gauge theory becomes strongly coupled for the energy scale below Λ_*, it is described by 15 composite meson fields $\mathcal{M}_{ij} \equiv \langle Q_i Q_j\rangle/(4\pi\Lambda_*)$ with the deformed moduli constraint $\mathrm{Pf}[\mathcal{M}_{ij}] = \Lambda_*^3/(4\pi)^3$. For simplicity, we

[2]H_u and H_d are the MSSM up-type and down-type Higgs $SU(2)_L$ doublets respectively. We denote other matter contents in the MSSM ($\mathbf{5}^*$ and $\mathbf{10}$) by using the representations of $SU(5)_{\rm GUT}$.

[3]One can also check that the charge assignment of the MSSM particle contents under Z_{6R} makes Z_{6R} anomaly free with respect to the MSSM gauge group [18].

[4]Although in principle the higher dimension operators $(c_{2n}/(2n)!)QQS(\Phi)^{2n}$ are allowed, for perturbative c_{2n}s they are negligible. So it suffices for us to consider the second term in the rhs of Eq. (2).

TABLE I. Quantum numbers of the matter contents of the model under the additional symmetry group in Eq. (1). All the MSSM nongauge interactions are consistent with the charge assignment.

	Q_i	S_{ij}	Φ	$\mathbf{5}^*$	$\mathbf{10}$	H_u	H_d	N	Z
$Sp(2)$	$\square$	⋯	⋯	⋯	⋯	⋯	⋯	⋯	⋯
Z_{6R}	1	0	3	0	0	2	2	0	4
Z_6	1	4	3	4	0	0	2	2	2

can make a choice of vacuum expectation values (VEV) of the quark fields (Q_i) such that the only nonvanishing meson fields are

$$\langle Q_i Q_{i+1} \rangle = \Lambda_*^2 \quad \text{for } i = 1, 3, 5, \tag{3}$$

where Λ_* is the dynamical scale of $Sp(2)$. The quark condensation induces the spontaneous breaking of Z_{6R} to Z_{2R}. This physics will be taken as the fundamental origin of the various energy scales in our scenario.

Given Eq. (3), in the nonperturbative regime of $Sp(2)$, the superpotential in Eq. (2) transforms to

$$\begin{aligned} W_{\text{eff}} \supset & -\lambda\Lambda_*^2 S(1 + g\Phi^2) \\ & + \lambda_H \Lambda_*^4 H_u H_d + \lambda_N \Lambda_*^2 NN + \lambda_Z \Lambda_*^4 Z \\ & = W_{\text{inf}} + W_H + W_N + W_{\cancel{\text{SUSY}}}, \end{aligned} \tag{4}$$

where all λs without the flavor indices are $\mathcal{O}(1)$ dimensionless coupling constants obtained after rescaling λs in Eq. (2). Here S is a linear combination of S_{ii+1} with $i = 1, 3, 5$.[5]

III. INFLATION (W_{inf})

As mentioned above, the inflation of the universe in our scenario is attributed to the first term of Eq. (4). Along with the Kähler potential of the form[6]

$$K(\Phi, S) \supset |S|^2 + |\Phi|^2 + c|S|^2|\Phi|^2 + \cdots, \tag{5}$$

the F-term contribution of S to the scalar potential of the model yields the inflaton potential below [14]

$$\begin{aligned} V &= e^K \left[\sum_{m,n} \left(\frac{\partial^2 K}{\partial X_m \partial X_n^*} \right)^{-1} D_{X_m} W D_{X_n^*} W^* - 3|W|^2 \right] \\ &\simeq \Lambda_*^4 e^{\frac{\phi^2}{2}} \left(1 - g\frac{\phi^2}{2} \right)^2 \left(1 + c\frac{\phi^2}{2} \right)^{-1}, \end{aligned} \tag{6}$$

where $X_m = S, \Phi$, and $D_{X_m} W = (\partial W/\partial X_m) + W(\partial K/\partial X_m)$, $\phi/\sqrt{2}$ is the real part of the scalar component of Φ and we have taken $|S|$ and the imaginary part of Φ to be their origin as we shall explain later. We note that we assume all couplings of Φ are suppressed by the Planck scale and hence loops diagrams are suppressed by more than Planck scale by the loop effects. However, we consider that there are tree-level operators like $c|S|^2|\Phi|^2$ in the Kähler potential which might dominate over such loop-induced operators. We assume the coefficient c is an unknown free parameter.

In order for the slow-roll of ϕ from somewhere near $\phi = 0$ to the field value at the global minimum to be responsible for the early inflationary era of the universe, the power spectrum of the curvature perturbation ($P_\zeta(k) = A_s(k/k_\star)^{n_s - 1}$) needs to satisfy $A_s = 2.1 \times 10^{-9}$, $n_s = 0.9649 \pm 0.0042$ (68% C.L., Planck TT, TE, EE + lowE + lensing) [21] and $r < 0.036$ (95% C.L., BICEP/Keck) [22] should hold at the CMB pivot scale $k_\star = 0.05\ \text{Mpc}^{-1}$. According to the scanning of the parameter space (Λ_*, c, g) performed in [14] with $0 < \phi_\star < 1$, $\Lambda_* \sim 10^{-3}$, $c \gtrsim 0.4$, $g \sim 0.3$ with $\phi_\star \sim 0.6$ were found to provide a good fit to the CMB observables. The model predicts a rather low tensor-to-scalar ratio $r = \mathcal{O}(10^{-4})$. The details were given in our previous work [14]. We briefly summarize the process of parameter inference. We note that our model has only three parameters. The observed spectral index n_s and the requirement of a consistent reheating history determine the coefficients in the superpotential and the Kähler potential, i.e., g and c. Then, with $A_s = V^3/12\pi^2 V'^2$, we can determine the potential height Λ_*^4. We note that it is remarkable for the inflation model in Eq. (6) to accomplish the successful fit to CMB observables with $\mathcal{O}(1)$ dimensionless parameters. Our previous work [14] pointed out that the spectrum of the gravitational wave induced by the short-lived cosmic string present during the reheating era can help test our scenario.

It is worth emphasizing a crucial difference of our model from our previous work [14] and other inflation models in the literature. There are two kinds of parameters, dimensionless parameters and dimensionful parameters; see Table II. Dimensionless parameters are naturally of $\mathcal{O}(1)$, while dimensionful parameters, such as the $SU(2)$ dynamical scale Λ_* in our Eq. (3), are naturally much smaller than the Planck scale. In the new inflation scenarios, the initial condition needs to be finely tuned to near the potential top [3,4]. This problem can be avoided in the large-scale topological inflation [23,24], but a dimensionless self-coupling parameter needs to be

[5] All other components of S_{ii+1} besides S have masses of order of the dynamical scale Λ_* together with composite bound states $Q_i Q_j$ and they are stabilized at the origins during the inflation [19,20].

[6] We comment on the role of the $|S|^4$ term later, and we assume all the other higher order terms are suppressed. However, we consider that those terms are not important for the inflation dynamics, since the inflation ends before the inflaton ϕ reaches its expectation value $\langle \phi \rangle = \mathcal{O}(M_p)$.

TABLE II. Dimensionful and dimensionless parameters.

Dimensionful parameters	Λ_*		
Dimensionless parameters	c	g	λ_{ij}

imposed to have a small potential height. Our previous work [14] can dynamically generate a topological inflation potential of the same form as Eq. (6), but is still plagued by the fact that it contains an unnaturally small dimensionless parameter in the UV completion. In our model, on the contrary, all the dimensionless parameters are of $\mathcal{O}(1)$, resulting in a vacuum expectation value (VEV) of ϕ around the Planck scale. Meanwhile, the height of the inflaton potential is given by the $SP(2)$ strong dynamics, which is naturally much smaller than the Planck scale (M_p^4, to be precise). This disparity of two scales is a key for generating the flat potential for inflaton without any small dimensionless parameters, which is a totally new mechanism.

The exact values of the dimensionless parameters in our model are determined by observations. But, this does not mean we have to tune those parameters to obtain a plateau-like infalton potential. As we discussed, a plateaulike inflaton potential is a direct consequence in our model. A mild tuning is required as we need $c + 2g \sim 1$. We do not consider this (about 5% tuning) as a severe fine tuning. With $c + 2g \sim 1$, the value of n_s is robust against c. Indeed, even setting $c = 0$ (with a suitable but flexible adjustment to g) only makes some change to predicted spectral index (i.e., n_s would become ∼0.95).

With the concrete inflaton potential specified above, now we discuss the appealing points of our inflation model. As is clear from Eq. (6), the potential is of the typical new inflation type. Albeit similar in the shape, there is a crucial distinction of Eq. (6) as compared to the original new inflation [3,4]: our inflation model is featured by the field value at the global minimum of $\phi_{\rm min} \sim 2.5$ and the inflation scale $H_{\rm inf} \simeq \Lambda_*^2 \sim 10^{-6} \ll 1$. Thanks to the (super-Planckian) large enough $\phi_{\rm min}$, the inflationary expansion in a false vacuum residing in a wall (the spatial region in-between domains with $\langle\phi\rangle = \pm\phi_{\rm min}$) is guaranteed when $\phi_{\rm min} > 1$ holds [23,24].[7] Hence, our model does not suffer from the initial condition problem. What is remarkable is that the separation of the scales of $V^{1/4}$ and $\phi_{\rm min}$, which is required by consistency with CMB observables and the use of logics in the topological inflation, is achieved without any fine-tuning of parameters in $V(\phi)$.

[7]When the wall thickness $\delta \sim \phi_{\rm min} V(0)^{-1/2}$ is greater than $H^{-1} \sim V(0)^{-1/2}$, the false vacuum region located in the wall experiences the inflationary expansion [23,24]. The condition $\delta > H^{-1}$ is converted to $\phi_{\rm min} > 1$.

IV. PREINFLATION ERA

At the Planck time $t \sim 1/M_p^{-1}$, there might be particle creation in an expanding background [25–27]. Afterwards, when $Sp(2)$ gauge theory becomes strong enough, the thermal bath made of multiplets of Q_i, S and $Sp(2)$ gluons is expected to form. For a given spatial distribution of T of thermal baths in mind, we may consider for simplicity two classes of region with $T < \Lambda_*$ and $T > \Lambda_*$. For the former, $V(\phi)$ in Eq. (6) applies since the horizon is already in the confined phase of $Sp(2)$. In contrast, for the later case, starting from zero potential, $V(\phi)$ eventually develops to the form in Eq. (6) at $t \sim (M_P/\Lambda_*)\Lambda_*^{-1}$.

Once the inflaton potential is described by Eq. (6) everywhere, we can expect that there arise a pair of neighboring Hubble patches with $\langle\phi\rangle \sim \pm\phi_{\rm min}$ after a bit of time for homogenizing ϕ within horizons. Then inflation is initiated at $\phi = 0$ residing in the wall between the Hubble patches (the space with ϕ field variation by $\sim 2\phi_{\rm min}$) [23,24]. Hence our model serves as a UV model for topological inflation scenario, providing the concrete picture for development of $V(\phi)$ at the preinflation era in accordance with the $Sp(2)$ strong dynamics.

V. SUSY BREAKING ($W_{\not{\rm SUSY}}$)

The last term in Eq. (4) explains the (F-term) SUSY breaking. Now that the Polonyi field Z is charged under both Z_{6R} and Z_6, there cannot be a marginal or relevant operator purely composed of Z. This fact makes the last operator in Eq. (2), i.e., $\mathcal{O} \sim QQQQZ$, the most relevant in the F-term contribution of Z to the scalar potential[8] and thus SUSY-breaking is indeed safely guaranteed once $Sp(2)$ enters the confined phase.[9] Even if the Polonyi field has no antisymmetric flavor indices of $Sp(2)$ quarks from the beginning, we see that this way of SUSY-breaking is similar to IYIT dynamical SUSY-breaking [19,20] in that $F_Z \neq 0$ is attributable to quark condensation with the deformed moduli constraint. Given $F_Z \sim \Lambda_*^4$, the model's prediction for the gravitino mass becomes $m_{3/2} = F_Z/(\sqrt{3}M_P) \sim \Lambda_*^4 = \mathcal{O}(10^{-12})$.

[8]We introduce a pair of massive chiral multiplets X and $\bar{X}$ which have a superpotential $W = ZX^2 + MX\bar{X}$. One-loop diagrams of the X and $\bar{X}$ give a large positive soft mass squared for Z to stabilize the potential of Z at the origin [28].

[9]When either of marginal or relevant operator purely made up of Z is allowed in the superpotential, the scalar potential $|F_Z|^2$ can possibly have a SUSY-preserving field value at the global minimum. Then SUSY remains unbroken. Were it not for Z_6, for instance, the operator $m_Z Z^2$ with a dimensionful parameter m_Z is allowed. Therefore, having other discrete symmetry than the discrete R-symmetry is the crucial point in constructing the SUSY-breaking sector in the model. One may be concerned about the unwanted operator $\sim ZH_uH_dN$. This one is not a problem since there is the parity as the remnant of $U(1)_{B-L}$ to suppress $\sim ZH_uH_dN$ to stabilize the SUSY vacuum. Under the parity, only $\mathbf{5}^*$, $\mathbf{10}$, and N transform as odd.

VI. HIGGSINO (W_H) AND RH NEUTRINO MASS (W_N)

With $\lambda_H, \lambda_N = \mathcal{O}(1)$, the second and third term in Eq. (4) can account for the Higgsino mass (μ_H) term in the MSSM and the right-handed (RH) neutrino masses. Given $\Lambda_* \sim 10^{-3}$ inferred from CMB observables, the model predicts $\mu_H = \mathcal{O}(10^{-12})$ and $m_N = \mathcal{O}(10^{-6})$. Below we discuss implications of the prediction.

Now that the soft masses are of the order $m_{3/2} = F_Z/(\sqrt{3}M_P)$ in supergravity (SUGRA) [29], we expect all of m_{H_u}, m_{H_d} and B to be $\mathcal{O}(m_{3/2})$. Hence, the prediction $\mu_H = \mathcal{O}(10^{-12})$ renders the electroweak symmetry breaking (EWSB) condition $(|\mu_H|^2 + m_{H_u}^2)(|\mu_H|^2 + m_{H_d}^2) \simeq (B\mu_H)^2$ nicely satisfied.

On the other hand, the heavy Majorana mass term $m_N NN$ and the symmetry allowed Yukawa interaction $\mathbf{5}^* H_u N$ in the superpotential can explain the tiny mass for the active neutrinos in the MSSM via the seesaw mechanism [30–32]. Also we note that $m_N = \mathcal{O}(10^{-6})$ is sufficiently heavy enough not to exceed the maximal baryon asymmetry in the leptogenesis [33,34].

VII. CONSTANT TERM IN W

The constant term in the superpotential (W_0) needs to be generated so as to properly cancel the SUSY-breaking F-term contributions to the scalar potential of the model. To this end, we may consider $SU(3)$ pure gauge theory, i.e., without any matter fields, with which Z_{6R} still remains gauge anomaly free. Then, once $SU(3)$ pure gauge theory enters its confinement at the scale Λ', there arises the gaugino condensation [35,36], i.e., $\langle \lambda^a \lambda^a \rangle = 32\pi^2 \Lambda'^3$. This results in the effective superpotential which take responsible for the constant term in the superpotential, i.e.,

$$W_0 = 3\Lambda'^3 = m_{3/2} \qquad \Rightarrow \qquad \Lambda' = \mathcal{O}(10^{-4}). \quad (7)$$

Given that Λ' is one order of magnitude smaller than Λ_*, we expect that W_0 is generated after PT of $Sp(2)$ and during inflation. $\langle \lambda^a \lambda^a \rangle$ respecting Z_{2R} but not Z_{6R}, there could be formation of the domain wall due to the gaugino condensation. This wall, however, will be diluted away during inflation.

In a SUGRA model, the difference between $|F_Z|^2$ and $m_{3/2}^2$ determines the scalar potential at the leading order. In our model, since both of F_Z and $m_{3/2}$ are of the order $\mathcal{O}(\Lambda_*^4)$, the required cancellation among the two can be achieved. We note that we are not resolving the fine-tuning of the cancellation to get a small cosmological constant. Rather, we use the condition of the vanishing cosmological constant to set the exact relation between F_Z and $m_{3/2}$. Interestingly, the resultant relation gives the correct μ_H-parameter along with the EWSB scale as previously discussed, which in turns supports the $Sp(2)$ strong dynamics as the common origin of operators $QQQQH_uH_d$ and $QQQQZ$.

The above constant term does not affect the inflation dynamics. The inflaton potential $V_{\rm inf} \sim \Lambda_*^4$ while the SUSY breaking potential is $V_{\rm SUSY} \sim \Lambda_*^8$ which is negligible. However, the constant term is $W_0 = \Lambda_* 4$, then we have $W_{\rm eff} = \Lambda_*^2 S + \Lambda_*^4$. From that we have a S linear potential $V = \Lambda_*^6 S$. But since S have a positive mass induced by Q loops as mentioned, $V = \Lambda_*^2 |S|^2 + \Lambda_*^6 S + \text{H.c.}$ Then, S will shift from the origin by $|\delta S| = \Lambda_*^4$ which is order of the gravitino mass $\sim 10^6$ GeV and thus negligible.

VIII. OTHER DYNAMICAL COMPONENTS

Here we comment on the dynamics of S and the imaginary part of Φ. First of all, the loops of the dynamical quarks give rise to a term of $-\frac{\eta}{4\times 16\pi^2}\frac{|S|^4}{\Lambda_*^2}$ in the effective Kähler potential [37] and we have taken $\lambda = 1$ here. $\eta = \mathcal{O}(1)$ accounts for the loop momentum factor. To the leading order, such a term gives a positive mass squared term $\eta \frac{\Lambda_*^2}{(4\pi)^2}|S|^2$ to the effective potential. Therefore, $m_S \simeq \Lambda_*/(4\pi) \sim 10^{-4}$ satisfies $m_S > H_{\rm inf} \simeq \Lambda_*^2 \sim 10^{-6}$ during the inflation and thus S will be stabilized at the origin. For the imaginary part of Φ, the relevant leading order term in the effective potential is $-2g\Lambda_*^4 |\Phi|^2 \cos(2\theta)$, where θ is the phase of Φ. Thus, as $|\Phi|$ becomes nonzero, θ get stabilized at 0 or π and hence only the real part of Φ is important.

IX. DISCUSSION AND OUTLOOK

In this paper, we proposed a new inflation model with the huge separation between the scale of the false vacuum ($V(\phi)^{1/4} \sim 10^{15}$ GeV) and the field value at the global minimum ($\phi_{\rm min} \simeq M_P$). We discussed how the physics in the preinflationary era in our model can justify the occurrence of inflation in a similar way to topological inflation. Thereby the model was shown free from the initial condition problem. What is new as compared to the arguments made in the topological inflation scenarios [23,24] lies in $Sp(2)$ strong dynamics-driven inflaton potential generation which importantly does not have any fine-tuned parameter.

The interesting observation that Z_{6R} and Z_6 are gauge anomaly free with respect to $Sp(2)$ within the $Sp(2)$ sector was invoked for achieving $V(\phi)$ without any tuning. The inflation scale was determined by the R-charged quark condensate and thus R-symmetry breaking scale, $\Lambda_* \sim 10^{-3}$, could be inferred from CMB observables. As the spurion field of Z_{6R} breaking, the quark condensate $\langle QQ \rangle \sim \Lambda_*^2$ also determines the Higgsino mass and the RH neutrino mass and thereby unifies the origin of various energy scales (the inflation scale, SUSY-breaking scale, R-symmetry breaking scale, Higgsino mass, and RH neutrino mass).

We also emphasize the key idea that the model underlies. The specialty of R-symmetry to apply to every operator in a superpotential W was used for explaining various energy scales and dimensionful parameters appearing in W. When every physics and operator yielding certain energy scales of interest is associated with a common symmetry, if experimental consistency allows, one may dream of explaining those energy scales as proper powers of spurion field of the symmetry. For our work, the symmetry was Z_{6R}.

With all phenomenologies handled in this work described by a single energy scale Λ_*, the model proposes the strong dynamics of $Sp(2)$ as the fundamental underlying physics governing the universe. Albeit ambitious, there are still open questions to be answered. Z_6 in Eq. (1) still remains anomalous for $SU(2)_L$ and thus poses the question about extension of $SU(2)_L$-charged particle contents.[10] On the other hand, as the SUGRA inflation model, the model is still subject to the long-standing η-problem, i.e., how to justify suppression of all the unwanted higher dimension operators in Eq. (5) contributing to $V(\phi)$. We leave these structural problems of the model to future work.

ACKNOWLEDGMENTS

T. T. Y. appreciates Li Fu for encouraging him to consider the origin of the universe. T. T. Y. is supported in part by the China Grant for Talent Scientific Start-Up Project and by Natural Science Foundation of China (NSFC) under grant No. 12175134 as well as by World Premier International Research Center Initiative (WPI Initiative), MEXT, Japan.

[10]One way to eliminate such an anomaly is to introduce H'_d and H'_u whose product $H'_u H'_d$ has charges (2,4) for Z_{6R} and Z_6, respectively.

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