Physics Letters B 773 (2017) 320–324

Contents lists available at ScienceDirect

Physics Letters B

www.elsevier.com/locate/physletb

Unified model of chaotic inflation and dynamical supersymmetry breaking

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ARTICLE INFO

Article history:
Received 27 July 2017
Accepted 22 August 2017
Available online 26 August 2017
Editor: G.F. Giudice

ABSTRACT

The large hierarchy between the Planck scale and the weak scale can be explained by the dynamical breaking of supersymmetry in strongly coupled gauge theories. Similarly, the hierarchy between the Planck scale and the energy scale of inflation may also originate from strong dynamics, which dynamically generate the inflaton potential. We present a model of the hidden sector which unifies these two ideas, i.e., in which the scales of inflation and supersymmetry breaking are provided by the dynamics of the same gauge group. The resultant inflation model is chaotic inflation with a fractional power-law potential in accord with the upper bound on the tensor-to-scalar ratio. The supersymmetry breaking scale can be much smaller than the inflation scale, so that the solution to the large hierarchy problem of the weak scale remains intact. As an intrinsic feature of our model, we find that the sgoldstino, which might disturb the inflationary dynamics, is automatically stabilized during inflation by dynamically generated corrections in the strongly coupled sector. This renders our model a field-theoretical realization of what is sometimes referred to as sgoldstino-less inflation.

1. Introduction

Cosmic inflation not only solves the flatness and horizon problems of big bang cosmology [1–4], but also explains the origin of the primordial density fluctuations that seed the large-scale structure of the universe [5–9]. To satisfy the upper bound on the tensor-to-scalar ratio in the power spectrum of the cosmic microwave background (CMB) [10], the potential energy during inflation must be much smaller than the scale of gravity, $\Lambda_{\rm inf} = V^{1/4} \lesssim 10^{-2} M_{\rm Pl}$. The smallness of the energy scale of inflation, $\Lambda_{\rm inf}$, is nicely explained if the inflaton potential V is generated by means of dimensional transmutation in a strongly coupled gauge theory. Refs. [11–14] and [15–18] proposed models of small-field and large-field inflation along this idea, respectively.

The electroweak scale also suffers from a hierarchy problem, $v_{\rm ew} \ll M_{\rm Pl}$, which can be solved by supersymmetry and its breaking at a low energy scale [19–22]. Again, a plausible explanation for the smallness of the supersymmetry breaking scale, $\Lambda_{\rm SUSY} \ll M_{\rm Pl}$, would be to presume that supersymmetry is broken dynamically by strong dynamics [21]. So far, no evidence for superpartners of the standard model particles has been found at the LHC, which has brought about the *little* hierarchy problem, $v_{\rm ew} \ll m_{\rm SUSY}$ (where $m_{\rm SUSY}$ denotes a typical soft superparticle mass). But supersymmetry nonetheless solves the *large* hierarchy problem, predicts the unification of the standard model gauge couplings and provides a particle candidate for dark matter. For these reasons, we take up the attitude that supersymmetry as well as its dynamical breaking are some of the leading candidates for new physics beyond the standard model.

In this letter, we propose a model of the hidden sector which unifies these two ideas of dynamically generated energy scales. The model resembles that of Refs. [15–17] during inflation; but the potential energy is non-zero even after the end of inflation, which breaks supersymmetry. The inflationary dynamics are those of chaotic inflation [23] with a fractional power-law potential. The model is thus free from an initial conditions problem; and it is consistent with the recent PLANCK data [10]. See Refs. [24,25] for other models of chaotic inflation with fractional power-law potentials. We also refer to Ref. [13] for an earlier proposal for the unified and dynamical generation of the energy scales of inflation and supersymmetry breaking, which results in a scenario of hybrid inflation [26,27]. This work has been followed up more

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http://dx.doi.org/10.1016/j.physletb.2017.08.050

recently in Refs. [28,29], where it is demonstrated how the dynamical breaking of supersymmetry at a very high energy scale may result in scenarios of F-term and D-term inflation, respectively. Finally, we refer to Ref. [30], which considers a perturbative model (as opposed to our strongly coupled models) in which the inflaton potential as well as the breaking of supersymmetry are both provided by the F term of a single chiral field.

2. Dynamical chaotic inflation

We first review the idea of dynamical chaotic inflation (DCI) proposed in Refs. [15–17]. We start from a strongly coupled gauge theory which generates a potential energy proportional to some power of the dynamical scale Λ,

$$V_{\rm dyn} \propto \Lambda^n . \tag{1}$$

To this theory, we add a pair of particles, q and $\bar{q}$, that obtain their mass from a coupling to the inflaton field ϕ,

$$\mathcal{L} = \lambda\phi\, q\bar{q}. \tag{2}$$

For a large field value of the inflaton, such that $\lambda\phi \gg \Lambda$, the fields q and $\bar{q}$ decouple; and around the dynamical scale the potential energy in Eq. (1) is generated. Since the energy scale at which $q\bar{q}$ decouples depends on the inflaton field value, the dynamical scale also depends on it, through the running of the gauge coupling constant,

$$\frac{\rm d}{{\rm d}\ln\mu}\frac{8\pi^2}{g^2(\mu)} = b, \tag{3}$$

where μ is the renormalization scale. We shall denote the beta function coefficient b in the high/low-energy theory with/without $q\bar{q}$ as $b_{\rm HE}$ and $b_{\rm LE}$, respectively. Then the effective dynamical scale $\Lambda(\lambda\phi)$ follows from

$$\begin{aligned} \frac{8\pi^2}{g^2(\mu_0)} - \frac{8\pi^2}{g^2(\lambda\phi)} &= b_{\rm HE}\ln\frac{\mu_0}{\lambda\phi}, \\ \frac{8\pi^2}{g^2(\lambda\phi)} &= b_{\rm LE}\ln\frac{\lambda\phi}{\Lambda(\lambda\phi)}, \end{aligned} \tag{4}$$

where g formally diverges, $g(\Lambda)\to\infty$, at the dynamical scale. Matching the running of the gauge coupling constant at the $q\bar{q}$ mass threshold, we obtain the dependence

$$\Lambda \propto \phi^{(b_{\rm LE}-b_{\rm HE})/b_{\rm LE}}. \tag{5}$$

Together with Eq. (1), this results in a power-law potential for the inflaton, ϕ^p, with the power p given as

$$p = n\,\frac{b_{\rm LE}-b_{\rm HE}}{b_{\rm LE}}. \tag{6}$$

This potential is suitable for inflation at large values of the inflaton field, $\phi \gg M_{\rm Pl}$, which is nothing but a (dynamical) realization of the idea of chaotic inflation.

The implementation of the above scheme into supersymmetric theories is straightforward. We start from a model of dynamical supersymmetry breaking, add chiral multiplets q and $\bar{q}$, and couple these chiral multiplets to the inflaton multiplet Φ. To avoid the eta problem in supergravity [31–34] for a large field value of the inflaton, we introduce an approximate shift symmetry $\Phi \to \Phi + iC$ in the Kähler potential [35,36]. The negative contribution to the potential energy is suppressed as long as the supersymmetry-breaking (Polonyi) field has a field value much smaller than the Planck scale during inflation.

Table 1
Matter content of the $SU(5)\times Sp(2)$ model.

	Q	$\bar{D}$	$\bar{U}$	$\bar{q}_{1,2}$	L
$SU(5)$	**5**	$\bar{\mathbf{5}}$	$\bar{\mathbf{5}}$	$\bar{\mathbf{5}}$	**1**
$Sp(2)$	**4**	**1**	**1**	**1**	**4**

3. Dynamical chaotic inflation and supersymmetry breaking unified

In this section, we propose a model of dynamical chaotic inflation in which the gauge dynamics also break supersymmetry in the true vacuum after the end of inflation. The basic idea is the following: We start from a dynamical supersymmetry breaking model with a product group $G_1\times G_2$, such that supersymmetry is broken by the strong dynamics of G_2, while the gauge interactions of G_1 merely lift flat directions by a classical D-term potential. To this model, we add G_2-charged matter fields ℓ and couple them to an inflaton multiplet, $W = \lambda\Phi\ell^2$. Supersymmetry is broken by the gauge dynamics of G_2 for large inflaton field values, where the new matter multiplets decouple. But for small field values, the gauge dynamics flow into a different phase; and the potential energy proportional to the dynamical scale of G_2 and hence the inflaton potential vanish. By a suitable choice of matter fields and couplings, supersymmetry is instead now broken by the strong dynamics of G_1 (or a subgroup of G_1, if the strong dynamics of G_2 partially break G_1). The supersymmetry breaking scale in the vacuum can be naturally much smaller than the scale of inflation, provided there is a hierarchy between the dynamical scales of G_1 and G_2 and/or the breaking of supersymmetry by the strong dynamics of G_1 involves particularly small couplings (realized, e.g., in the form of higher-dimensional operators). In this paper, we shall present a simple realization of this idea based on the groups $G_1 = SU(5)$ and $G_2 = Sp(2)$. Other examples will be given elsewhere.

3.1. $SU(5)\times Sp(2)$ model during inflation

Let us apply the idea described in Sec. 2 to the $SU(5)\times Sp(2)$ model of supersymmetry breaking [37], which is a generalization of the so-called 3–2 model [38]. The model is based on $SU(5)\times Sp(2)$ gauge dynamics and features chiral multiplets Q, $\bar{U}$, $\bar{D}$ L, $\bar{q}_{1,2}$ in representations of the gauge group as listed in Table 1. Our convention for $Sp(N)$ groups is such that $Sp(1)\cong SU(2)$. The theory contains the following flat directions,

$$Q\,\bar{Q}\,L,\;\; Q\,Q\,\bar{Q}\,\bar{Q}\,, \tag{7}$$

where $\bar{Q}\in\{\bar{D},\bar{U},\bar{q}_i\}$. The flat directions are lifted by introducing the following tree level superpotential,

$$W_{\rm tree} = y\,Q\,\bar{D}L + \frac{1}{M_*}Q\,Q\,\bar{q}_1\bar{q}_2\,. \tag{8}$$

In this paper, we concentrate on the case where the dynamical scale of $SU(5)$ is much smaller than that of $Sp(2)$, $\Lambda_{SU}\ll\Lambda_{Sp}$. Supersymmetry is then broken by the deformed moduli constraint [39] of the $Sp(2)$ dynamics, which results in non-zero F terms for $\bar{D}$ and the flat direction $Q\,Q\,\bar{q}_1\bar{q}_2$. The potential energy is given by [40]

$$V_{Sp} \sim y^{3/2}\left(\frac{\Lambda_{Sp}}{M_*}\right)^{1/2}\Lambda_{Sp}^4. \tag{9}$$

To turn this supersymmetry breaking model into a model of dynamical chaotic inflation, we add $Sp(2)$-charged chiral multiplets ℓ and couple them to the inflaton field Φ,

$$W = \lambda \Phi \ell^2 . \tag{10}$$

For $\lambda\Phi \gg \Lambda_{Sp}$ the extra multiplets ℓ decouple from the gauge dynamics. The theory then exhibits supersymmetry breaking and generates a non-zero potential energy.

The supersymmetry-breaking field is contained in $\bar{D}$ and $Q\,Q\,\bar{q}_1\bar{q}_2$. Its scalar component, the sgoldstino, is a flat direction at tree level, which could potentially disturb the inflationary dynamics. It, however, obtains a mass from strong-coupling corrections in the Kähler potential,

$$m \sim y^{7/8} \frac{\Lambda_{Sp}^{9/8}}{M_*^{1/8}}, \tag{11}$$

as is the case in generic models of dynamical supersymmetry breaking. Unless y is small, m is much larger than the Hubble scale, which provides a field-theoretical realization of the so-called sgoldstino-less inflation [41]. This is a generic feature in models of dynamical chaotic inflation. We note that the stabilization by a Hubble-induced mass would already be enough to ignore the sgoldstino dynamics [35,36]; but the stabilization via IR quantum corrections is advantageous in the sense that it is independent of the unknown UV physics which determine the sign and the magnitude of the Hubble-induced mass.

3.2. Flow into $SU(5)$ model in the vacuum

After inflation, at $\lambda\Phi \ll \Lambda_{Sp}$, the extra multiplets ℓ no longer decouple, but participate in the gauge interactions just like the other $Sp(2)$ flavors. In Refs. [15–17], the fields ℓ as well as their couplings were chosen so that the theory reaches a phase of s-confinement at low energies, where all flat directions are lifted and supersymmetry is restored. In this paper, we are instead going to chose the matter content and couplings such that supersymmetry remains broken even in the true vacuum after inflation.

We add a pair of $Sp(2)$ fundamentals, ℓ_1 and ℓ_2, and introduce a coupling to the inflaton multiplet Φ,

$$W = \lambda \Phi \ell_1 \ell_2 \tag{12}$$

The beta function coefficient of the $Sp(2)$ gauge coupling at high and low energies is then given as $b_{\rm HE} = 5$ and $b_{\rm LE} = 6$, respectively. The potential energy during inflation scales like Λ_{Sp} to the power $n = 9/2$, see Eq. (9), so that the exponent of the inflaton potential is given by $p = 3/4$, see Eq. (6). The dynamical scale around the vacuum, $\tilde{\Lambda}_{Sp}$, and the dynamical scale during inflation Λ_{Sp} are related to each other as follows, see Eq. (5),

$$\Lambda_{Sp} = \tilde{\Lambda}_{Sp} \left(\frac{\lambda \Phi}{\tilde{\Lambda}_{Sp}} \right)^{1/6} . \tag{13}$$

Around $\Phi = 0$, the $Sp(2)$ gauge theory reaches a phase of s-confinement; and the low-energy theory is described in terms of 28 gauge-invariant, composite meson fields,

$$M_{Q\,Q}\,,\ M_{Q\,L}\,,\ M_{Q\,\ell_{1,2}}\,,\ M_{L\ell_{1,2}}\,,\ M_{\ell_1\ell_2}\,. \tag{14}$$

The fields $(M_{Q\,L}, \bar{D})$ and $(M_{\ell_1\ell_2}, \Phi)$ obtain their masses from the superpotential in Eqs. (8) and (12), respectively. The inflaton mass around the origin is thus given by

$$m_\Phi \sim \lambda \tilde{\Lambda}_{\rm Sp}\,. \tag{15}$$

After those fields decouple, the theory still contains the following chiral multiplets

$$M_{Q\,Q}\ (\mathbf{10}),\ M_{Q\,\ell_{1,2}}\ (\mathbf{5}),\ M_{L\ell_{1,2}}\ (\mathbf{1}),\ \bar{U}\ (\bar{\mathbf{5}}),\ \bar{q}_{1,2}\ (\bar{\mathbf{5}}), \tag{16}$$

where the numbers in bold refer to representations of $SU(5)$. The superpotential in the s-confined phase reads

$$W \sim \frac{\tilde{\Lambda}_{\rm Sp}}{M_*} M_{Q\,Q}\,\bar{q}_1\bar{q}_2 + \frac{1}{\tilde{\Lambda}_{\rm Sp}} M_{Q\,Q}^2 \left(M_{Q\,\ell_1} M_{L\ell_2} + M_{Q\,\ell_2} M_{L\ell_1} \right) . \tag{17}$$

Here, the second line is generated by the $Sp(2)$ dynamics.

The theory now contains one **10**, two **5**'s, and three $\bar{\mathbf{5}}$'s of $SU(5)$. By giving masses to two pairs of $\mathbf{5}+\bar{\mathbf{5}}$, the theory becomes nothing but the chiral supersymmetry breaking model based on $SU(5)$, featuring one **10** and one $\bar{\mathbf{5}}$ of $SU(5)$ [42]. The vacuum energy is then given by

$$V_{\rm vac} \sim \tilde{\Lambda}_{SU}^4\,, \tag{18}$$

where $\tilde{\Lambda}_{SU}$ is the dynamical scale of $SU(5)$ in the low-energy effective theory containing only $\mathbf{10}+\bar{\mathbf{5}}$. We may obtain a hierarchy between the inflation scale and the supersymmetry breaking scale by choosing $\tilde{\Lambda}_{SU} \ll \Lambda_{Sp}$.

The $SU(5)$ singlets $M_{L\ell_1}$ and $M_{L\ell_2}$ remain massless. We can stabilize these fields by introducing $Sp(2)$ singlets and coupling them to $L\ell_1$ and $L\ell_2$ in the quark picture at high energies. Another possibility would be to simply introduce a higher-dimensional operator, $W = L\ell_i L\ell_j$.

Depending on how we give masses to the two pairs of $\mathbf{5}+\bar{\mathbf{5}}$, the inflaton potential could be affected. We may, e.g., remove the fields $M_{Q\,\ell_{1,2}}$ and $\bar{q}_{1,2}$ by adding the following superpotential in the quark picture,

$$W = \kappa_1\, Q\,\ell_1 \bar{q}_1 + \kappa_2\, Q\,\ell_2 \bar{q}_2\,, \tag{19}$$

such that the matter content of the $SU(5)$ supersymmetry breaking model is provided by the chiral fields $M_{Q\,Q}$ and $\bar{U}$. After s-confinement of $Sp(2)$, those terms give masses to the $(M_{Q\,\ell_1}, \bar{q}_1)$ and $(M_{Q\,\ell_2}, \bar{q}_2)$ pairs. At the same time, during inflation and after integrating out $\ell_1\ell_2$, this superpotential also generates the second term in Eq. (8) with $M_* \propto \Phi$. When this inflaton-dependent term dominates over the Φ-independent one, the inflaton potential becomes the one with $p = 3/4 - 1/2 = 1/4$. If they are comparable to each other, we have $p = 1/4$ for small field values and $p = 3/4$ for large field values.

4. Phenomenology of inflation

4.1. CMB observables

Taken all together, the model constructed in Sec. 3 results in an inflaton potential of the following form,

$$V = c \left| \frac{e^{i\alpha}}{M_*} + \frac{1}{\phi} \right|^{1/2} \left(\frac{\lambda \phi}{\Lambda_{Sp}} \right)^{3/4} \Lambda_{Sp}^{9/2} . \tag{20}$$

Here, we choose a convention in which both ϕ and M_* are real and positive; and the phase difference between these two complex parameters is accounted for by the phase α. The parameter c is a numerical constant, which we will set to $c = 1$ in the following. The scalar potential is only monotonically increasing for positive ϕ as long as $|\alpha/\pi| \leq 5/6$. For values of $|\alpha/\pi|$ closer to unity, the potential exhibits a false vacuum at small field values.

From the potential in Eq. (20), we derive the predictions for the CMB observables, i.e., for the scalar spectral index n_s as well as for the tensor-to-scalar ratio r. The result of our analysis is shown in Fig. 1. The predictions for both parameters only depend on M_* and α. If there is a clear hierarchy between M_* and ϕ for all times

Table 2
Charges under the Z_2 symmetry that forbids the decay of the inflaton into the supersymmetry breaking sector.

	Φ $\bar{D}$ $\bar{q}_1$ L ℓ_1	Q $\bar{U}$ $\bar{q}_2$ ℓ_2
Z_2	−	+

during inflation, we simply recover the predictions for chaotic inflation based on a standard power-law potential, $V \propto \phi^p$,

$$n_s = 1 - \frac{p+2}{2N_e} = 1 - 0.025\left(\frac{p+2}{3/4+2}\right)\left(\frac{55}{N_e}\right), \tag{21}$$

$$r = \frac{4p}{N_e} = 0.055\left(\frac{p}{3/4}\right)\left(\frac{55}{N_e}\right), \tag{22}$$

where N_e is the number of e-folds at the CMB pivot scale. For $M_*^{-1} \gtrsim M_{\rm Pl}^{-1}$, the M_*^{-1} term in Eq. (20) clearly dominates over the ϕ^{-1} term. In this case, we effectively obtain $p = 3/4$. On the other hand, if the M_*^{-1} term should be suppressed by a small coupling in Eq. (8) or by an (approximate) symmetry, such that $M_*^{-1} \lesssim 0.01 M_{\rm Pl}^{-1}$, it can be neglected throughout inflation and we can effectively work with $p = 1/4$. For intermediate values of M_*, the predictions for n_s and r are more complicated, as they become sensitive to the phase α. This is evident from Fig. 1, where we show the variation of n_s and r for different values of α. In particular, we observe how, for fixed α, the variation of M_* results in *orbits* in the n_s–r plane that connect the predictions for $p = 3/4$ and $p = 1/4$.

The parametric freedom of our model makes it easy to achieve consistency with the recent PLANCK data [10]. Our model predicts values of r in the $r \sim 0.01 \cdots 0.1$ range and is, therefore, in accord with the current upper bound, $r \lesssim 0.1$. In particular, close-to-maximal values of the phase, $\alpha \simeq 5/6\pi$, allow to achieve rather large values of r, which are going to be tested in future CMB experiments. Our model moreover prefers values of n_s in the $n_s \sim 0.97 \cdots 0.99$ range, which is slightly above the current best-fit value, $n_s \simeq 0.965$. It is however interesting to note that the data still admits such relatively large values of n_s, if it is fit by a $\Lambda{\rm CDM} + r + N_{\rm eff}$ model, which also accounts for the possibility of dark radiation.

For given values of M_* and α, the observed amplitude of the scalar power spectrum, $A_s \simeq 2 \times 10^{-9}$, fixes the parameter combination $\lambda^{1/5}\Lambda_{Sp}$ in Eq. (20). We find that, in the entire parameter space of interest, this product is required to take a value of around $\lambda^{1/5}\Lambda_{Sp} \sim 10^{16}$ GeV. At the same time, λ should not be too small, since otherwise the matter fields ℓ_1 and ℓ_2 do not decouple for the entire duration of inflation. We demand that $\lambda\phi \gtrsim \Lambda_{Sp}$ at all times during inflation, which roughly translates into $\lambda \gtrsim 10^{-2}$, see Ref. [17] for details. Given this lower bound on λ, we then find that the required value of Λ_{Sp} is always remarkably close to the scale of grand unification.

4.2. Reheating

After inflation, the energy density stored in the inflaton field must be transferred into standard model particles. In our model, the inflaton resides in the supersymmetry breaking sector, such that it may dominantly decay into particles in this sector. Those particles eventually decay into gravitinos, which easily leads to an overproduction of gravitinos. We can forbid the decay mode into the supersymmetry breaking sector by symmetry arguments. For example, we can impose the Z_2 symmetry shown in Table 2, under which the inflaton is odd. The particles in the $SU(5)$ model, M_{QQ} and $\bar{U}$, are Z_2-even and, hence, the inflaton does not decay into these states. The other Z_2-odd particles obtain masses proportional

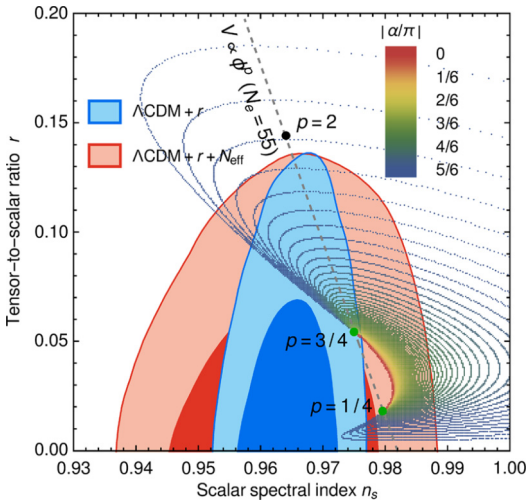

Fig. 1. Predictions of our model for n_s and r, compared with the latest constraints according to the PLANCK 2015 data (68% and 95% C.L., TT, TE, EE + lowP) [10]. The blue contours correspond to the standard $\Lambda{\rm CDM} + r$ fit, whereas the red contours also take into account the possibility of a non-standard number of relativistic degrees of freedom, $N_{\rm eff}$, at the time of photon decoupling. The color scale indicates the value of the phase α, which we vary on a linear scale. For each value of α, we vary the mass scale M_* in the interval $\left[10^{-3}, 10^{3}\right] M_{\rm Pl}$ on a logarithmic scale. This results in *orbits* in the n_s–r plane that smoothly connect the predictions of the pure power-law potentials $\phi^{3/4}$ and $\phi^{1/4}$. The local density of points in the above plot can be regarded as a measure for how "generic" or "typical" a certain prediction is. A low density of points indicates a rather special parameter choice, while a high density of points indicates that a prediction is stable under small variations of the input parameters M_* and α. (For interpretation of the references to color in this figure legend, the reader is referred to the web version of this article.)

to $\tilde{\Lambda}_{Sp}$. If λ is sufficiently small, the inflaton ends up being the lightest particle in the supersymmetry breaking sector, so that it does not decay into any particles in this sector.

The Z_2 symmetry also forbids the operator $Q\,Q\,\bar{q}_1\bar{q}_2$ in Eq. (8). Therefore, if the Z_2 is an *exact* symmetry, the M_*^{-1} term in Eq. (20) is actually no longer present. In our analysis, this corresponds to taking the limit $M_* \to \infty$, such that the scalar potential reduces to an exact power-law with $p = 1/4$. On the other hand, if the Z_2 is only an *approximate* symmetry, it only suppresses the M_*^{-1} term to some degree. In this case, we have to work with the full scalar potential in Eq. (20) and the predictions for the CMB observables depend on the exact hierarchy between M_*^{-1} and ϕ^{-1}, as discussed in the previous section.

The inflaton can decay, e.g., via a coupling to the Higgs multiplets $H_{u,d}$ in the minimal supersymmetric standard model, $W = \epsilon\,\Phi H_u H_d$. In this case, the μ term is generated via Z_2 symmetry breaking. We may also identify the Z_2 with R parity and introduce $W = \epsilon_i\,\Phi L_i H_u$, where the L_i denote the standard model lepton doublets [43].

5. Discussion

In this letter, we presented a strongly coupled model of the hidden sector based on $SU(5) \times Sp(2)$ gauge dynamics. Our model combines the ideas of dynamical supersymmetry breaking and dynamical chaotic inflation and, hence, explains the hierarchy between the scales of supersymmetry breaking, inflation, and gravity, $\Lambda_{\rm SUSY} \ll \Lambda_{\rm inf} \ll M_{\rm Pl}$. During inflation, supersymmetry is broken because of the $Sp(2)$ deformed moduli constraint. This results in an inflaton potential that interpolates between the power-law potentials $\phi^{3/4}$ and $\phi^{1/4}$, see Fig. 1. The pseudoflat sgoldstino direction is automatically stabilized during inflation by dynamically generated corrections in the Kähler potential. After inflation, the $Sp(2)$ sector reaches a phase of s-confinement and supersymmetry

is broken by the $SU(5)$ gauge interactions. In fact, at low energies, our model reduces to the chiral $SU(5)$ model of dynamical supersymmetry breaking.

In the $SU(5)$ model, some approximate global symmetries are believed to be spontaneously broken, which results in the presence of (pseudo-) Nambu–Goldstone bosons. These bosons obtain non-zero field values in the early universe and may affect the cosmological history. Among them, the R axion is potentially dangerous, since it has a mass squared of $\mathcal{O}(m_{3/2}\tilde{\Lambda}_{SU})$ through the explicit breaking of R symmetry [44] and because it dominantly decays into gravitinos. The gravitino eventually decays into the lightest supersymmetric particle (LSP), which may lead to its overproduction. Assuming that the initial amplitude of the R axion is as large as $\tilde{\Lambda}_{SU}$, the LSP abundance is estimated as

$$\frac{\rho_{\rm LSP}}{s} \sim m_{\rm LSP} \frac{T_{\rm RH}}{M_{\rm Pl}} \left(\frac{m_{3/2}}{M_{\rm Pl}}\right)^{1/4}, \tag{23}$$

where s is the entropy density and $T_{\rm RH}$ the reheating temperature. Here, we imposed the condition that the universe must reach a flat Minkowski vacuum after inflation, $m_{3/2}M_{\rm Pl} \sim \tilde{\Lambda}_{SU}^2$. Requiring that $\rho_{\rm LSP}/s < 4 \times 10^{-10}$ GeV, we obtain an upper bound on $T_{\rm RH}$,

$$T_{\rm RH} \lesssim 10^9\ {\rm GeV} \left(\frac{100\,{\rm TeV}}{m_{3/2}}\right)^{1/4} \left(\frac{1\,{\rm TeV}}{m_{\rm LSP}}\right). \tag{24}$$

In the $SU(5)$ model, the gaugino masses of the minimal supersymmetric standard model are generated only via anomaly mediation [45–50], meaning that they are loop-suppressed compared to the gravitino mass. The scalar masses, on the other hand, follow from the tree-level Kähler potential and are as large as (or larger than) the gravitino mass. For $m_{3/2} \sim \mathcal{O}(100 \cdots 1000)$ TeV, our model is thus compatible with the scenario of high-scale supersymmetry breaking [46,51,52], which has gained considerable interest after the discovery of the standard model Higgs boson with a mass of 126 GeV [53,54].

By choosing a different gauge group, we may also obtain a model of gauge mediation. For example, we can modify our model by gauging only the $SU(3) \times SU(2) \times U(1)$ subgroup of $SU(5)$. By adding an appropriate superpotential term, supersymmetry is broken via the 3–2 model in the vacuum. The $U(1)$ symmetry may be used as the messenger hypercharge [55]. We leave a detailed discussion of modifications of our model for future work.

Acknowledgements

The authors are grateful to the organizers of the 20th PLANCK conference held at the University of Warsaw in May 2017, where this project was initiated. The work of K.H. was supported in part by the Director, Office of Science, and Office of High Energy and Nuclear Physics of the U.S. Department of Energy under contract DE-AC02-05CH11231 and by the National Science Foundation under grants PHY-1316783 and PHY-1521446. This project has received funding from the European Union's Horizon 2020 research and innovation programme under the Marie Sklodowska-Curie grant agreement No. 674896 (K.S.). K.S. acknowledges the hospitality of the Laboratoire Univers et Particules de Montpellier during a stay at the University of Montpellier in July 2017, where this project was finished.

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