

## SQCD inflation & SUSY breaking

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## SQCD inflation & SUSY breaking

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**Philippe Brax, Carlos A. Savoy and Arunansu Sil**

*Institut de Physique Theorique, CEA/Saclay,  
F-91191 Gif-sur-Yvette Cedex, France*

*E-mail:* [philippe.brax@cea.fr](mailto:philippe.brax@cea.fr), [carlos.savoy@cea.fr](mailto:carlos.savoy@cea.fr), [arunansu.sil@cea.fr](mailto:arunansu.sil@cea.fr)

**ABSTRACT:** A model of generalized hybrid inflation in a supersymmetric QCD theory is proposed whose parameters are the gauge coupling and quark masses. Its gravitational coupling to another SQCD sector induces a metastable supersymmetry breaking vacuum of the ISS type as ISS quarks become massive at the end of inflation. Using a known mechanism with a gravitational breaking of the baryon number and the gauging of flavour symmetries, we find that gauge mediation of supersymmetry breaking is compatible with the dynamics of the inflation sector. Reheating proceeds via the thermalization of the ISS messengers into the standard model states. This setup contains a single dimensionful parameter in the form of a quark mass term in the inflationary sector, i.e. all other scales involved are either related to this single mass parameter or dynamically generated.

**KEYWORDS:** Cosmology of Theories beyond the SM, Supersymmetry Breaking

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## 1 Introduction

New large effective scales are suggested by many phenomenological approaches to physics beyond the Standard Model: neutrino masses, baryogenesis, FCNC, CP violations and, in particular, inflation. In supersymmetric theories - where the interplay between large and small scales is more comfortable and  $M_P = 2.4 \times 10^{18}$  GeV is a natural cutoff - the gravitational interactions in the effective supergravity lagrangians have been often taken into account to generate renormalizable operators when some fields take relatively large vacuum expectation values (v.e.v.'s). Of particular interest are the possibilities of gravity mediation of symmetry breaking effects from an otherwise hidden sector to the observable world. The classical example is supergravity mediation of supersymmetry breaking<sup>1</sup> to the supersymmetric extensions of the Standard Model (SM), generically referred to here as MSSM, where the effective supersymmetry breaking is suppressed by a  $M_P$  factor. Another kind of examples arises in models where the decay of the inflaton into SM states, the reheating, arises from gravitational operators.

In the present paper we investigate whether gravitational interactions between two hidden sectors responsible, respectively, for supersymmetry breaking and inflation could be relevant and whether the relatively large scales generated in supersymmetric inflation could be responsible for supersymmetry breaking in the other sector. This is done here in the context of two attractive hidden sector candidates: supersymmetry breaking in a metastable vacuum of a supersymmetric QCD (SQCD) theory [2] and inflation in another SQCD model, where the peculiar properties of the inflationary model are fixed by the choice of the numbers of colours and flavours. This novel inflationary model is interesting

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<sup>1</sup>For a review, see, e.g. [1].

*per se* because its effective superpotential arises from a SQCD as its UV completion. We also discuss the feedback to the inflation scenario: how the reheating could proceed through the states of the ISS sector.

In dynamical supersymmetry breaking (DSSB) by means of the Intriligator-Seiberg-Shih mechanism [2] (ISS), the Universe lives in the metastable vacuum of a asymptotically free SQCD where the overall chiral symmetry is broken by a quark mass term. The resulting superpotential in the dual magnetic theory is of the O’Rafertaigh type, where the linear term, proportional to the quark mass, defines the supersymmetric breaking scale. The  $R$ -symmetry of the superpotential is only broken by its non-perturbative term that induces a supersymmetric vacuum. The (meta-)stability is ensured if the DSSB scale - hence the quark mass in the electric theory - is small enough as compared to the SQCD scale.

The effective supersymmetry breaking in the standard supersymmetric gauge and matter sector (MSSM) requires a mediation mechanism of the DSSB from the otherwise hidden ISS sector. However, in its original version, ISS does not fulfill the well-known requirements to implement gaugino masses, namely,  $R$ -symmetry breaking [3] and scalars with charges  $R \neq 0, 2$  [4]. The original ISS picture must be enriched [5] and we follow here the elegant suggestion of [6] and introduce a baryonic term in the superpotential to fix the magnetic quarks to have charges  $R = 1$  and to produce spontaneous  $R$ -symmetry breaking. Then, it is possible to identify a subgroup of the flavour symmetry with the GUT  $SU(5)$  group (or the corresponding SM gauge symmetry subgroups). The magnetic quarks become plausible messenger candidates for gauge mediation of DSSB (GMSB)<sup>2</sup> as confirmed by a phenomenological analysis [8], which finds, in particular, that  $O(\text{TeV})$  gaugino masses, require a relatively low supersymmetry breaking scale,  $O(10^8 \text{ GeV})$ .

Although quite attractive, the resulting model has two free scale parameters much smaller than the scale associated to the gauge coupling: the DSSB one - related to the electric quark mass in ISS - and the coefficient of the baryonic term, which controls the  $R$ -symmetry breaking. Furthermore, in the effective magnetic theory both parameters must be of the same order of magnitude to produce realistic mass spectra in the MSSM through gauge mediation [8]. These low scales are protected by global symmetries, equivalent to a quark chirality and a baryon-number and we would like here to propose a possible origin for their breaking at relatively low scales.

In recent works [9, 10], the authors have shown that the scales of inflation and supersymmetry breaking in a metastable vacuum can be naturally connected from the assumption that the two sectors are coupled only by (super)gravity. Of course, for the suggestion to be meaningful, the inflationary model must be also endowed with some basic properties of the ISS model: supersymmetry,  $R$ -symmetry and, possibly, other global symmetries, which should also be present in the gravitational interactions. As a case study for providing such relations between inflation and supersymmetry breaking, supersymmetric hybrid inflation was discussed in [9]: the scale of hybrid inflation is constrained by experimental data to be very large and the resulting supersymmetry breaking is also large, just consistent with supergravity mediation. Then, with a new model for inflation, which includes a shift

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<sup>2</sup>For a review, see, *e.g.* [7].

symmetric Kähler potential along with an inflationary potential of the hilltop type, and leads for a much lower scale for inflation, we showed in [10] that the scale of supersymmetry breaking comes out proportionally smaller.<sup>3</sup>

In this paper, we want to investigate a similar scenario with a further motivation to include the  $R$ -symmetry breaking which is necessary for making gauginos massive. Since in this scheme the DSSB scale is related to the inflation scale, we need an inflation model such that the ISS effective scales come out naturally low. To implement this requirement and to carry on the parallelism with the ISS model and, also the MSSM, we build a SQCD inflationary model where, apart from the dynamically generated scale, the dimensional parameters have a UV interpretation in terms of quark masses. Besides this theoretical motivation, the model turns out to be quite consistent with inflation data. We show that, at the end of inflation, its  $R$ -symmetric gravitational couplings with the fields in the ISS sector can generate low scale DSSB and  $R$ -symmetry breaking, suitable for gauge mediation.

For not too small quark masses in the inflation SQCD, the mesons in its IR phase are massive enough for the inflaton to mainly decay into ISS magnetic quarks, in particular those coupled to the ordinary matter through the SM gauge sector which play the role of messengers in the model. The calculated reheating temperature is quite acceptable. Therefore, in this scheme, one has a sort of “ISS mediation” of the inflaton energy to the SM or, more precisely, the MSSM particles.

As a step further, we suggest an upgrade of our inflationary model where the dimensional parameters in the latter are obtained as quark masses of the former. This is done by increasing the number of flavours by one unit, from  $N_f = N$  to  $N_f = N + 1$ , with one heavier quark whose mass then corresponds to the second parameter needed in the smooth hybrid model, so completing our SQCD picture of inflation.

The paper is organized as follows. The basic assumptions are presented in section II and are followed, in section III, by a brief review of the deformed ISS model of [6] including the phenomenological constraints on both parameters in the effective O’Raifeartaigh superpotential. The SQCD inflationary model is described and confronted to data in section IV to obtain the constraints on the parameters. It is then coupled through supergravity to the ISS fields in section V to produce the deformation of the ISS model with the low scales in the magnetic phase fixed up to  $O(1)$  factors by both the dynamical SQCD scales involved. Section VI explains how the ordinary matter and gauge particles are produced through the production and decays of the GMSB messengers in the ISS sector during the reheating period. The last section summarizes our conclusions.

## 2 The basic setup: a 3-fold Universe

Before presenting our model thoroughly, it will prove useful to give a simpler version of the essential ingredients. The more complete picture will follow in the rest of this section. In particular, we will pay attention to the embedding of the following simple model into a full-fledged gauge theoretical description in both the inflation and the supersymmetry breaking

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<sup>3</sup>A different inflation scenario where the inflaton rolls down to the ISS metastable vacuum has been proposed in [11].

sectors. On the inflation side, the effective description below the scale  $\Lambda_0 = O(10^{17}\text{GeV})$  is given in terms of two fields  $S$ , the inflaton field, and  $\chi$ , the waterfall field, whose dynamics are defined by the superpotential [15]

$$W_{\text{Infl}} = S \left( \frac{\chi^4}{\Lambda_0^2} - \Lambda_{\text{eff}}^2 \right). \tag{2.1}$$

The model has a supersymmetric minimum for  $\langle S \rangle = 0$  and  $\langle \chi \rangle = (\Lambda_{\text{eff}} \Lambda_0)^{1/2}$  where  $\Lambda_{\text{eff}} = O(10^{14.7}\text{GeV})$ . There is a flat valley in the  $(S, \chi)$  plane (will be discussed in detail later) along which inflation takes place with a spectral index of order 0.97 and a reheate temperature of order  $T_R = O(10^{2-3}\text{GeV})$  is generated at the end of inflation.

The inflation sector is gravitationally coupled to a supersymmetry breaking sector of the ISS type so that the supersymmetry breaking scale ( $\mu$ ) depends on  $\chi$  as

$$\mu^2 \approx \frac{\Lambda_0^4 \Lambda_s}{M_P^2} \langle \chi \rangle^4, \tag{2.2}$$

where  $\Lambda_s = O(10^{15}\text{GeV})$  is the cut-off of the effective O’Raifeartaigh model that describes the ISS sector at low energy. Once inflation stops and  $\chi$  settles at the global minimum of the inflation potential, supersymmetry is broken with  $\mu = O(10^{8-9})\text{GeV}$ , of the right order of magnitude to be mediated to the MSSM by gauge interactions. In the following we will motivate and present the details leading to the above results. In particular, we shall justify the dependence of eqs. (2.1) and (2.2) on the field  $\chi$  and on the three scales above.

As in [9], the set-up consists of three components namely, the inflationary sector (Infl), the supersymmetry breaking sector (here the deformed ISS sector, dISS) and, of course, the MSSM sector. Each sector is phenomenologically very constrained. For the MSSM, we only consider here general constraints related to scalar and gaugino masses.

First of all, fields in different sectors have only (super)gravitational interactions described by an effective supergravity theory where the superpotential and a Kähler potential respect the symmetries of the different sectors. We want to study the possible interferences between the phenomenologies in the different sector due to those gravitational interactions. They are generically present in supergravity and are restricted by some symmetries common to both sectors, e.g.,  $R$ -symmetry. Actually, the connection between the potential flatness and  $R$ -symmetry makes its choice rather natural as flatness is the basic idea in hybrid inflation models as well as in building ISS metastable vacua. They could break less essential accidental symmetries, e.g., chiral symmetries broken by quark mass terms.

However we are forced to break these seclusion rules in one respect. Indeed DSSB must be mediated from dISS to MSSM. Gravity mediation was envisaged in [9], but the fundamental scales come out all very high, the gaugino masses were not accounted for in that first analysis. In GMSB one assumes that gauge multiplets are coupled to both sectors and the minimal choice [6] is that the SM gauge particles couple to some of the (so-called) quarks in the magnetic phase of the dISS sector which then become messengers. Therefore the SM gauge symmetries must be embedded into the ISS chiral flavour group (in the present paper in a vector-like way). This model for GMSB with ISS is also adopted

here. The consequences have been widely discussed in the literature [8] and are summarized below. In our setup, it could lead to a novel reheating mechanism.

As mentioned in the Introduction, we assume here that each sector is described by its own SQCD theory, with its gauge sector, its chiral multiplets called quarks and, by dimensional transmutation, a characteristic scale. In practice, we take simpler (in particular simple) gauge groups and representations than in the MSSM, basically one  $SU(N)$  with  $N_f$  vector-like pairs of fundamental representations for the chiral multiplets, or “quarks”. For ISS, the IR or magnetic theory with a metastable vacuum has a known UV completion,<sup>4</sup> an (electric) asymptotically free SQCD if it satisfies  $N + 1 \leq N_f \leq \frac{3}{2}N$ . The implementation of  $R$ -symmetry breaking as in [6] then requires  $N = 5$  and  $N_f = 7$  where the  $SU(7)$  factor in the flavour chiral symmetry is explicitly broken by a baryonic term to  $SU(5) \otimes SU(2)$ .

SQCD with  $N = N_f = 2$  has been discussed in the literature [13] because its low energy superpotential is of the supersymmetric hybrid inflation type [14] with a flat inflaton direction in the potential. It leads to a less satisfactory model of inflation and a large supersymmetry breaking [9] from its gravitational coupling to ISS. Here we alter the inflation scenario by increasing the number of colours,  $N$  and also by considering  $N_f = N + 1$  SQCD inflation. For  $N = 4$ , the low energy effective superpotential in terms of so-called mesons and baryons has a similarity with the so-called smooth hybrid inflation model (SHI) [15].

However, there are two key differences: *a)* there are two waterfall fields in SHI to be compared to the  $N^2 - 1$  mesons in  $N = N_f$  SQCD, which in the vanishing quark mass limit are massless goldstone bosons; *b)* SHI has two free scales, a cutoff  $\Lambda_0$  somewhat below  $M_P$  and the inflation scale  $\Lambda_{\text{eff}}^2$  which in our model are defined as the strong coupling scale and  $m_Q \Lambda_0$ , respectively, where  $m_Q$  is the mass of the additional quark in the UV region. The addition of this quark promotes the theory to the  $N_f = N + 1$  theories. The SHI superpotential and its free parameters are justified as the IR dynamics of a SQCD theory. Actually, this fact is the main motivation (on the inflation side) for the new inflationary model proposed in section VII.

The  $R$ -symmetric supergravity coupling between the dISS and the Infl mesons which produce a linear term in the ISS superpotential, namely, the mass term for the electric quarks that induces DSSB. The DSSB scale is then controlled by the two inflation parameters. The CMB data leaves only one free parameter: the inflation scale can be tuned down by lowering the scale  $\Lambda_0$  without affecting the agreement with data. Because it also controls the DSSB scale in the dISS sector, this one can be lowered as well. As a result,  $\Lambda_0$  turns out to be fixed by the feedback from the MSSM phenomenology after GMSB. This shows a tight connexion between the dISS and Infl phenomenologies in our scenario.

We also find a new reheating mechanism. Reheating presupposes some link between the Infl and the MSSM sectors. This was possible in [9] by a gravitational coupling of the waterfall fields to right-handed neutrinos. In the present case this coupling is reduced by the required symmetries. Nevertheless, the inflaton field now decays into dISS quarks and the MSSM particles are produced by thermalization since the two sectors intersect through the MSSM gauge fields. Reheating is mediated by ISS quarks.

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<sup>4</sup>For a review, see, *e.g.*, [12].

### 3 ISS sector: supersymmetry breaking in a SQCD metastable vacuum

In this section we briefly review the ISS approach and the modification suggested in [6]. In the UV region, the dISS sector corresponds to an asymptotic free SQCD with  $N = 5$  colours and  $N_f = 7$  flavours for the quark superfields,  $Q_i^\alpha, \tilde{Q}_i^\alpha$  (the electric theory) where  $i = 1, \dots, 7$  and  $\alpha = 1, \dots, 5$ . The GMSB kit, consists of flavoured particles in the dISS sector as messengers that couple to the MSSM only through these gauge interactions: a subgroup of the flavour group containing the SM group is to be gauged. For simplicity we gauge a whole  $SU(5)$  and identify it with the GUT  $SU(5)$ . The chiral  $SU(7)$  flavour symmetries are then broken to  $SU(2) \otimes SU(5)$  by construction. For simplicity we still refer to the gauged  $SU(5)$  as a flavour group in this section.

Below the strong coupling scale  $\Lambda_s$  the model has a dual description in terms of a IR free magnetic theory with magnetic gauge group  $SU(2)$  and  $N_f = 7$  quarks and antiquarks,  $q_i^a, \tilde{q}_i^a$ , where  $a = 1, 2$  denotes the magnetic colours, together with the so-called mesons which define a matrix  $\Phi_{ij}$ . The matching of the degrees of freedom and of the preserved flavour symmetries is ensured by a superpotential in the magnetic theory. The deformed ISS model has two more terms in the superpotential and reads:

$$W_{\text{dISS}} = \Phi_{ij} q_i \tilde{q}_j - \mu_{ij}^2 \Phi_{ji} + m_q \epsilon_{ab} \epsilon_{rs} q_r^a q_s^b + (N_f - N) \left( h^{N_f} \frac{\det \Phi}{\Lambda_s^{3N - 2N_f}} \right)^{\frac{1}{N_f - N}}, \quad (3.1)$$

where  $\mu_{ij}^2 = \text{diag}(\mu_2^2, \mu_2^2, \mu_5^2, \dots, \mu_5^2)$ , and the indices  $r, s = 1, 2$  denote the quarks that are singlets under the GUT  $SU(5)$ . The first and last term constitute the superpotential of the magnetic dual of the UV SQCD. The second term was introduced in ISS and amounts to a breaking of two chiral symmetries: it is responsible for the supersymmetry breaking in a metastable vacuum at a scale proportional to the parameters  $\mu_2$  and/or  $\mu_5$ .

The third term that “deforms” the ISS model, behaves as a  $SU(2)$  baryon, hence it breaks a baryon number and leads to the dynamical breaking of  $R$ -symmetry. Indeed, if one neglects the last term, for  $\Phi \ll \Lambda_s$ , the remaining terms have all  $R = 2$  if one defines the  $R$ -charges of  $\Phi, q$  and  $\tilde{q}$  as 2, 1 and  $-1$ , respectively (in the original ISS model, the latter can be taken to be 0). Therefore the  $R$ -symmetry can be broken at a scale  $O(m_q)$ . Actually, this baryon term introduces a runaway direction towards a non-supersymmetric vacuum at infinity, but the potential is stabilized by the Coleman-Weinberg radiative corrections, proportional to the supersymmetry breaking parameters,  $\mu_2$  and/or  $\mu_5$ .

It was shown [6] that the introduction of this term in eq. (3.1) of the magnetic dual theory is responsible for shifting  $\langle \Phi \rangle$  away from zero, thereby breaking the  $U(1)_R$ . Without going into many details [6], we need to recall a few points that are crucial in our discussion. It has been shown that the potential is minimized when  $\Phi$  is diagonal and does not break  $SU(2) \otimes SU(5)$ , its v.e.v.'s being  $O(m_q)$ , while only the  $SU(5)$  singlet quarks get a v.e.v so that the SM (here the GUT) symmetry is preserved as it must be. With some technically natural choice of the parameters,

$$\mu_2 \simeq \mu_5 = O(m_q), \quad (3.2)$$



one obtains a viable model with DSSB and enough  $R$ -symmetry breaking by fields with  $R \neq 0, 2$  and coupled to the  $SU(5)$  as required to give masses to gauginos and implement GMSB. The mass degeneracy of the dISS squarks with non-trivial SM quantum numbers is broken giving masses to the MSSM gauginos. The requirement that gaugino masses are of order  $O(\text{GeV})$  leads to the constraint  $\mu_5 = O(10^8 \text{ GeV})$ .

Our goal is to provide a mechanism to generate these parameters and reproduce those constraints from the coupling the model for inflation. For an analysis of the dISS phenomenology see [8].

#### 4 Inflation sector: a SQCD model

We discuss here an inflationary model represented by a strongly coupled supersymmetric  $SU(\mathcal{N})$  gauge group with  $\mathcal{N}_f = \mathcal{N}$  flavours of quark superfields  $\mathcal{Q}_i$  and  $\bar{\mathcal{Q}}_i$  ( $i = 1, \dots, \mathcal{N}_f$ ) in the  $\mathcal{N}$  and  $\bar{\mathcal{N}}$  representations of the gauge group. The system has a non-anomalous global symmetry  $G = SU(\mathcal{N}) \otimes SU(\mathcal{N}) \otimes U_B(1) \otimes U_R(1)$ . It turns out that the case  $\mathcal{N} = 4$ , which is a SQCD generalization of the smooth hybrid inflation model [15, 16], fits better the data and we concentrate on it from now on.

Below the scale  $\Lambda_0$  (where the  $SU(4)$  gauge coupling becomes large),<sup>5</sup> the theory is described by an effective theory of composite mesons in the representation  $(\underline{4}, \underline{4})$  of the chiral flavour group, one baryon and one anti-baryon,

$$T_{ij} = \frac{1}{2} \Lambda_0^{-1} \mathcal{Q}_i^a \bar{\mathcal{Q}}_j^a, \quad B = \frac{1}{3} \Lambda_0^{-3} \epsilon_{ijkl} \mathcal{Q}_i^1 \mathcal{Q}_j^2 \mathcal{Q}_k^3 \mathcal{Q}_l^4, \quad \bar{B} = \frac{1}{3} \Lambda_0^{-3} \epsilon_{1234} \bar{\mathcal{Q}}_i^1 \bar{\mathcal{Q}}_j^2 \bar{\mathcal{Q}}_k^3 \bar{\mathcal{Q}}_l^4, \quad (4.1)$$

where the superscripts are colour indices. All these fields have  $R = 0$  and another field, denoted by  $S$ , must be introduced to implement the charge  $R = 2$  for the superpotential, which is then fixed as the flavour symmetry invariant<sup>6</sup>

$$W_{\text{Inf}} = S \left( \frac{\det T}{\Lambda_0^2} - B\bar{B} - \Lambda_{\text{eff}}^2 \right), \quad (4.2)$$

Notice the presence of an additional scale,  $\Lambda_{\text{eff}}$  that we treat here as a parameter and refer to [12] and section VI for a discussion of its meaning.

In the so-called meson branch of the theory, where  $T$  does not vanish, it can be represented in terms of a non-linear realization by Nambu-Goldstone bosons (NGB) of the global symmetry  $SU(4) \times SU(4)$  broken down to  $SU(4)_V$  as follows,

$$T = \chi \exp \left( i \frac{t^\alpha \lambda^\alpha}{\langle \chi \rangle} \right), \quad (4.3)$$

with  $\alpha = 1, \dots, 15$  where  $t^\alpha$  represents the NGB superfields and  $\lambda^\alpha$  the  $SU(4)$  generators. Obviously these NGB are cosmologically relevant and will be treated in the reheating

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<sup>5</sup>More precisely (see below)  $\Lambda_0$  is the strong coupling scale of the parent SQCD theory with one more flavour.

<sup>6</sup> For reviews see, e.g., [12]. The relationship between the existence of the superpotential, and the matching of anomalies and degrees of freedom is also discussed in [17] on more general grounds.

section. The superpotential for  $B = \bar{B} = 0$  and replacing  $\det T = \chi^4$  in eq. (4.2) coincides with the smooth hybrid inflation one, namely,

$$W_{\text{Infl}} = S \left( \frac{\chi^4}{\Lambda_0^2} - \Lambda_{\text{eff}}^2 \right), \tag{4.4}$$

which has a supersymmetric minimum at

$$\langle \chi \rangle = (\Lambda_{\text{eff}} \Lambda_0)^{1/2}, \quad S = 0. \tag{4.5}$$

But there is a relevant difference between the two model besides the presence of the NGB: the scales in eq. (4.4) have now a physical meaning as the gauge theory scale for the cutoff  $\Lambda_0$  and the scale  $\Lambda_{\text{eff}}$  can be related to the explicit breaking of a chiral symmetry if one massive flavour is added to the  $\mathcal{N}_f = \mathcal{N}$  SQCD [12]. To see it, let us start with a  $\mathcal{N}_f = \mathcal{N} + 1 = 5$  SQCD, and add a superpotential corresponding to a quark mass term,

$$W_m = \text{Tr} \hat{m} Q \bar{Q}, \tag{4.6}$$

where the trace is taken over the five flavours and four colours. At low energies this theory is better formulated in terms of the mesons and baryons which are defined analogously to eq. (4.1) but, since now  $i = 1, \dots, 5$ , the baryons carry a free flavour index,  $\hat{B}^i$  and the meson matrix  $\hat{T}_{ij}$  is correspondingly larger. The low energy superpotential is,

$$\hat{W} = \hat{B} \hat{T} \hat{B} - \frac{\det \hat{T}}{\Lambda_0^2} + \Lambda_0 \text{Tr} \hat{m} \hat{T}, \tag{4.7}$$

where  $\Lambda_0$  is the strong coupling scale of the  $\mathcal{N}_f = \mathcal{N} + 1$  SQCD and the last term is the counterpart of eq. (4.6).

The  $R$ -symmetry, that plays an important rôle for the flatness of the potential, is preserved if one chooses  $\hat{m}_Q = \text{diag}(0, 0, 0, 0, m_Q)$  corresponding to only one massive quark.<sup>7</sup> The heavy degrees of freedom can be integrated out to define an effective theory at low energies (the descent relation discussed, e.g., in [12]) in terms of the baryons  $B = \hat{B}^5, \bar{B} = \hat{\bar{B}}^5, T_{ij} = \hat{T}_{ij}$  for  $i = 1, 2, 3, 4$ , and  $S = \hat{T}_{55}$ . The effective superpotential coincides with eq. (4.2) with the identification

$$\Lambda_{\text{eff}}^2 = m_Q \Lambda_0, \tag{4.8}$$

Therefore the smooth hybrid inflation superpotential can be dynamically generated from a UV gauge theory completion rather than justified by *ad hoc* global symmetries. The cutoff scale is defined from the gauge coupling, while the vacuum energy during inflation,  $\Lambda_{\text{eff}}^4$  is determined by a quark mass. Furthermore, the power of the waterfall field  $\chi$  in eq. (4.4) is given by the number of colours,  $\mathcal{N}$ .

Now we turn to discuss how the scales  $\Lambda_0$  and  $\Lambda_{\text{eff}}$  are constrained by the Cosmic Microwave Background (CMB) experimental data [18] and WMAP [19]. We start with the superpotential in eq. (4.4) with  $B = \bar{B} = 0$ , the so-called meson branch, for simplicity.

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<sup>7</sup>Though in section VI, we shall complete the model by adding additional masses to the remaining quarks, hence to the associated goldstone bosons.

The scalar potential in terms of the real normalized fields  $\sigma = \sqrt{2}\Re(S)$  and  $\xi = \sqrt{2}\Re(\chi)$  is given by

$$V_{\text{Inf}}(\sigma, \phi) = \left( \frac{\xi^4}{4\Lambda_0^4} - \Lambda_{\text{eff}}^2 \right)^2 + \frac{\sigma^2 \xi^6}{\Lambda_0^4}, \quad (4.9)$$

where scalar components are described by the same notations as superfields. Following the standard smooth hybrid inflation [15], we see here that the flat direction at  $\xi = 0$  is now a local maximum for all values of  $\sigma$  and there are two symmetric valleys of minima present at  $\xi \simeq \pm \frac{\Lambda_{\text{eff}} \Lambda_0}{\sqrt{3}\sigma}$ .

An interesting point is that the valleys contain the global SUSY minimum which lie at  $\xi = \sqrt{2\Lambda_{\text{eff}}\Lambda_0}$ ,  $\sigma = 0$  and have a slope which can actually drive the inflaton  $\sigma$  towards the right vacuum. The potential along this valley is

$$V(\sigma) \simeq \Lambda_{\text{eff}}^4 \left( 1 - \frac{1}{54} \frac{\Lambda_{\text{eff}}^2 \Lambda_0^2}{\sigma^4} + \dots \right), \quad \text{for } \sigma \gg \sqrt{\Lambda_{\text{eff}} \Lambda_0}. \quad (4.10)$$

We identify the slow roll parameters

$$\epsilon = \frac{M_P^2}{2} \left( \frac{V'(\sigma)}{V(\sigma)} \right)^2 \simeq \frac{2}{729} \left( \frac{\Lambda_{\text{eff}}^2 \Lambda_0^2 M_P}{\sigma^5} \right)^2, \quad (4.11)$$

$$|\eta| = M_P^2 \left| \frac{V''(\sigma)}{V(\sigma)} \right| \simeq \frac{10}{27\pi} \frac{\Lambda_{\text{eff}}^2 \Lambda_0^2 M_P^2}{\sigma^6}, \quad (4.12)$$

where we have used the approximated  $V$  with  $\sigma \gg \sqrt{\Lambda_{\text{eff}} \Lambda_0}$ . The number of  $e$ -folds during inflation and the temperature fluctuation are estimated to be

$$N_l = \frac{1}{M_P^2} \int_{\sigma_0}^{\sigma_l} \frac{V(\sigma) d\sigma}{V'(\sigma)} \simeq \frac{5}{6|\eta|}, \quad (4.13)$$

$$\Delta = \left( \frac{\delta T}{T} \right)_Q = \left( \frac{32\pi}{45} \right)^{1/2} \frac{V^{3/2}(\sigma)}{V'(\sigma) M_P^3} \simeq \left( \frac{9}{8\sqrt{5}\pi} \right) \frac{1}{M_P^3} \frac{\sigma_l^5}{\Lambda_0^2}, \quad (4.14)$$

where  $\sigma_l$  and  $\sigma_0$  indicate the values of the inflaton field when the ‘comoving’ scale  $l$  crossed outside the event horizon and the end of inflation (corresponding to the slow roll parameter,  $|\eta| = 1$ ) respectively. Using these parameters, the spectral index of density fluctuations,  $n_s$  is estimated as

$$n_s = 1 - 6\epsilon + 2\eta \simeq 1 - \frac{5}{3N_l} \simeq 0.97, \quad (4.15)$$

for  $N_l \sim 56$ .

We use the data from COBE [18] ( $\delta T/T \simeq 6.6 \times 10^{-6}$ ) and the WMAP 5 results [19] ( $n_s \simeq 0.964 \pm 0.014$ ) to fix the scales. The value of  $n_s$  of eq. (4.15) is in agreement with the experimental result. From eqs. (4.13), (4.14), we can write the scale  $\Lambda_{\text{eff}}$ , or equivalently,  $m_Q$ , in terms of  $\Lambda_0$  as follows,

$$\Lambda_{\text{eff}} \simeq 0.6 \Delta^{3/5} \Lambda_0^{1/5} M_P^{4/5}, \quad m_Q \simeq 2.2 \times 10^{-7} \left( \frac{M_P}{\Lambda_0} \right)^{3/5} M_P, \quad (4.16)$$

where we put  $N_l = 57$ . One important point should be observed. The  $\epsilon$  parameter at the time of horizon exit is given by

$$\epsilon(\sigma_l) = \frac{1}{72} \frac{1}{N_l^2} \left( \frac{\sigma_l}{M_P} \right)^2 \lesssim 10^{-5}, \quad (4.17)$$

it depends upon  $\sigma_l$  which is not fixed by the present data. We will use this freedom, in the next section, to relate the inflationary scale with the scale of supersymmetry breaking once we have fixed the interaction between the two sectors.

So far we have neglected the supergravity corrections. With a canonical Kahler potential, the effective scalar potential for  $\sigma$  in supergravity is given by

$$V(\sigma) = \Lambda_{\text{eff}}^4 \left[ 1 - \frac{1}{54} \frac{\Lambda_{\text{eff}}^2 \Lambda_0^2}{\sigma^4} + \frac{\sigma^4}{M_P^4} \right] \quad \text{for } \sigma \gg \sqrt{\Lambda_{\text{eff}} \Lambda_0}. \quad (4.18)$$

As long as  $\sigma \ll M_P$ , the inflationary dynamics are dominated by the false vacuum energy density  $\sim \Lambda_{\text{eff}}^4$  and the supergravity corrections do not modify it much [16]. However the supergravity correction during inflation is important for the fields in the dISS sector as we discuss now. The superfields  $\Phi, q, \tilde{q}$  all get Hubble induced mass-square terms  $\sim H^2$  and thereby settles to zero. In [9], this coincides with the supersymmetry breaking minimum for  $\Phi$ . Here during inflation  $\mu$  is very small but nonzero. As soon as  $H$  is decreasing after inflation and the Coleman-Weinberg correction stabilizes the runaway direction towards supersymmetry breaking minimum of the dISS sector, the component fields of  $\Phi$  will roll down to their minima since those are close to the origin rather than to roll towards the supersymmetric minimum which is far away. At the end of inflation, the inflaton field performs damped oscillations about the supersymmetric minimum of the inflationary sector and decays. We will discuss this part in section VI.

## 5 Supersymmetry breaking as a remnant of inflation

We now turn to discuss how the parameters  $\mu^2$ , responsible for spontaneous supersymmetry breaking, and  $m_q$ , which drives  $R$ -symmetry breaking, could originate from gravitational interactions in the UV superpotential. In particular, how terms suppressed by inverse powers of  $M_P$  can couple a pure SQCD in the dISS sector to the SQCD fields of the Infl sector to produce the superpotential of the ISS model.

These couplings are controlled by the  $R$ -symmetry present in both sectors. In the Infl sector, the  $R$ -charges are only fixed as 2, 0,  $1-x$ ,  $1+x$  ( $\forall x$ ), for  $S, T, Q$  and  $\bar{Q}$ , respectively. In the original ISS model, the fields  $\Phi, q, \bar{q}$  have  $R$ -charges 2,  $r, -r$ , where  $r = 0$  can be chosen, so that  $R$  is unbroken at the metastable minimum, while the  $R$ -charges of the UV fields  $Q, \bar{Q}$  are not uniquely fixed. In the dISS generalization, the presence of the baryon term in eq. (3.1) fixes  $r = 1$ . Then, identifying the IR baryons (two quarks  $q$ ) to the UV ones (five quarks  $Q$ ), one finds  $R = \frac{2}{5}$  for  $Q$ 's and, from  $R = 2$  for  $\Phi$ , one gets  $R = \frac{8}{5}$  for  $\bar{Q}$ 's.

Let us now construct new UV superpotential interactions with  $R = 2$ . Assuming the explicit breaking of the baryon number in dISS by the third term in eq. (3.1),  $m_q q e q$  one can write its  $R = 2$  avatar in the UV completion as

$$W_{\text{BV}} = \frac{1}{M_P^2} \epsilon_{ijklm} Q_i^1 Q_j^2 Q_k^3 Q_l^4 Q_\mu^5, \quad (5.1)$$

The supergravity cutoff  $O(M_P)$ , characterizes a gravitational coupling, the only one allowed at low orders and generically present in the superpotential unless the baryon symmetry is

imposed. This gives the relation

$$m_q = O\left(\frac{\Lambda_s^3}{M_P^2}\right). \tag{5.2}$$

which fixes  $\Lambda_s$  once  $m_q$  is fixed by the MSSM phenomenology. Therefore the  $R$ -symmetry breaking is controlled by the scale associated to the dISS coupling,  $\Lambda_s$ .

Now we turn to the main point of this section, the generation of the DSSB scale from its gravitational coupling to the inflation sector. We take for granted that the dISS sector is secluded enough so that quark chiral symmetries are not produced from supergravity couplings or other sources.<sup>8</sup> Therefore, in contrast to the original ISS model [2], we do not assume any explicit mass terms for the quarks in the UV completion. Instead, as in [9], we consider that the two sectors, dISS and Infl, communicate only via gravity. The lowest dimensional UV term which respects  $U(1)_R$  and the other chiral  $SU(4) \otimes SU(4)$  flavour symmetries in the inflation sector is given by<sup>9</sup>

$$W_{\text{int}} = \frac{\det Q \bar{Q}}{M_P^7} \text{Tr} f Q \tilde{Q} \quad \longrightarrow \quad \frac{\Lambda_0^4 \Lambda_s}{M_P^7} \det T \text{Tr} \Phi, \tag{5.3}$$

where the IR avatar of the UV operator is also indicated by an arrow. Once the field  $\det T$  gets a vacuum expectation value from eq. (4.5) at the end of inflation,  $W_{\text{int}}$  generates the linear terms in the dISS superpotential and gives for their coefficients:

$$\mu_i^2 = f_i \frac{\Lambda_{\text{eff}}^2 \Lambda_0^6 \Lambda_s}{M_P^7} = O\left(\frac{\Lambda_0^7}{M_P^7}\right) \Lambda_s m_Q = O(10^{-7}) \left(\frac{\Lambda_0}{M_P}\right)^{27/5} \Lambda_0 \Lambda_s. \tag{5.4}$$

where we have inserted the phenomenological constraint from eq. (4.16) between the two inflation parameters to obtain the last equality.

The phenomenological aspect of the model follows immediately since in order to have the right amount of gauge mediation from the dISS sector to the MSSM sector we must impose that both the scales of supersymmetry ( $\mu$ ) and  $R$ -symmetry ( $m_q$ ) breaking should be  $O(10^8 \text{ GeV})$ . This fixes the strong coupling scales,  $\Lambda_0$  and  $\Lambda_s$ , of both the Infl and dISS sectors. Note that  $\sigma_l$  is  $O(10^{16-17} \text{ GeV})$  and therefore  $\epsilon$  turns out to be  $O(10^{-8})$ .

In table 1, we have summarized the different scales involved in the problem. We have given two examples corresponding to two values of the supersymmetry breaking scale. It should be noted that the scale  $\Lambda_0$  is larger than  $\Lambda_s$ . Inflation must be valid to higher energies than the ISS sector as it gives a mass to the electric quarks, this mass becoming the effective supersymmetry breaking scale at low energy. In a sense, inflation is a precursor to supersymmetry breaking. As the inflation scale  $\Lambda_0$  is close to the Planck scale, its UV completion is a gauge theory whose domain of validity must be compatible with physics at

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<sup>8</sup> The quark-antiquark representation being vector-like with respect to the gauged flavour symmetries, the GUT or SM ones, this is an open issue, analogous to the  $\mu$ -problem, and we are making a similar assumption.

<sup>9</sup>One can think of  $\text{Tr} T Q \tilde{Q}$  term instead of the determinant. But this would not be invariant under the bigger symmetry group,  $SU(4) \otimes SU(4) \otimes U_B(1) \otimes U_R(1)$ . Chiral symmetry breaking in the inflationary SQCD will be introduced in the next section, but appear at a much lower scale.

$\mu$	$\Lambda_0$	$m_Q$	$\Lambda_s$	$\langle\chi\rangle$	$\Lambda_{\text{eff}}$	$T_R$
$2.5 \times 10^8$	$7.1 \times 10^{16}$	$4.3 \times 10^{12}$	$7.6 \times 10^{14}$	$6.2 \times 10^{15}$	$5.5 \times 10^{14}$	100
$10^9$	$10^{17}$	$3.5 \times 10^{12}$	$1.2 \times 10^{15}$	$7.75 \times 10^{15}$	$5.9 \times 10^{14}$	1330

**Table 1.** Different scales involved in the scenario in units of GeV.

energies close to the Planck scale. One enticing possibility would be to realise the inflation sector in a brane construction [20]. The same could also be true of the ISS sector [21]. In this case, the coupling between the inflation and the ISS sector could be understood as springing from gravitational effects in the bulk. The construction of explicit brane models is of course beyond the scope of the present paper.

## 6 Reheating from gauge mediation

Let us now discuss the inflaton decay and reheating. The inflaton fields smoothly enter an era of damped oscillation about the supersymmetric vacuum of the Infl sector. The oscillating system has a common mass  $m_{\text{inf}} = 2\sqrt{2}\frac{\Lambda_{\text{eff}}^2}{\langle\chi\rangle}$  and will decay eventually to reheat the universe. There will be Nambu-Goldstone bosons (NGB) present from the Inf sector due to spontaneous breaking of the global symmetry as we discussed before. Their derivative type of coupling therefore indicates that the inflaton system could decay into those goldstons with a decay width  $\Gamma_{\text{inf}\rightarrow\text{NGB}} \simeq \frac{1}{64\pi}\frac{m_\chi^3}{\Lambda^2}$ . These particles would be produced copiously during reheating and their abundance could spoil the success of big-bang nucleosynthesis (otherwise with massless NGB, it could just be part of the radiation component of the universe). Here we prescribe a resolution of this cosmological problem and show that effectively the reheating will take place via *ISS mediation*.

Without spoiling our description of the inflation sector, we perform a small deformation of the set-up by choosing  $\hat{m}_Q = \text{diag}(m, \dots, m, m_Q)$  with  $m \ll m_Q$  in  $N_f = N + 1$  case to we end up with an extra  $m\Lambda_0\text{Tr}T$  term in the superpotential of eq. (4.6). In the limit  $m \rightarrow 0$ , it can be shown that the supersymmetric minimum coincides with  $B\bar{B} = 0$ ,  $\det\langle T \rangle = m_Q\Lambda_0^3$ ,  $S = 0$ . The insertion of the  $m\Lambda_0\chi$  term in the superpotential of Infl sector will induce a tadpole term for  $S$  in the scalar potential that would shift the vev of  $S$  from zero to  $\langle S \rangle \simeq m\langle\chi\rangle/m_Q$  at the end of inflation. The inclusion of this new mass term for quarks would break the chiral symmetry in the Infl sector explicitly and we would expect a mass term for the pseudo-NGBs. Following an analogy with the pion mass and using a variant of Dashen formula [22], we can argue that the these pseudo-NGBs will get a mass,

$$m_t^2 = O(1)m\frac{\Lambda_0^3}{\langle\chi\rangle^2}. \tag{6.1}$$

In order to forbid the decay of the inflaton into the NGBs kinematically we impose the constraint,  $m_{\text{inf}} < 2m_t$ , namely

$$O(1)\frac{m_Q^2}{\Lambda_0} < m < m_Q, \tag{6.2}$$

where we have included the fact that  $m_Q < \Lambda_0$  as turned out from our analysis. We conclude that  $10^{-4} < m/m_Q < 1$  (see table 1) resolves the problem of NGBs.

Then we see that the inflaton decays into the magnetic quarks  $q, \tilde{q}$  of the dISS sector,

$$V \ni \left| \frac{\partial W}{\partial \Phi} \right|^2 = \left| q\tilde{q} + f_{2,5} \frac{\chi^4 \Lambda_0^4}{M_P^7} \Lambda_s \right|^2. \quad (6.3)$$

Since part of these  $q, \tilde{q}$  are charged under SU(5) of MSSM after gauging, the particular decay mode  $\chi \rightarrow q\tilde{q}$  is instrumental for the production of MSSM particles through their subsequent annihilation. The corresponding decay width is therefore given by

$$\Gamma_{\text{inf} \rightarrow q\tilde{q}} = \frac{5}{8\pi} \left( \frac{\mu^2}{\langle \chi \rangle} \right)^2 \frac{1}{m_{\text{inf}}}. \quad (6.4)$$

This particular way of reheating (we phrase it as *ISS mediation*) is a general feature of our scenario. This reheating mechanism is a new feature of our scenario.<sup>10</sup> The reheating temperature is given in table 1. As one can see, the reheat temperature is very sensitive to the details of the models although it is always larger than the electroweak scale.

## 7 Conclusions

The origin of the supersymmetry breaking scale is mysterious as no precise model has been derived so far after more than 25 years of intensive efforts. This lack of understanding of the supersymmetry breaking mechanism plagues the possibility of carrying out a predictive comparison of the sparticle spectrum with present and future particle data coming from LEP and soon the LHC. Indeed, depending on the breaking scale many methods can be envisaged in order to mediate the supersymmetry breaking to the observable sector. Two main scenarios offer widely different results, gravity mediation requires a large breaking scale while gauge mediation can accommodate a much smaller scale. The latter allows a description at low energy without the need to invoke what happens close to the Planck scale (or the GUT scale). The absence of a well founded description plagues the analysis of inflation too. There again despite 25 years of efforts, hundreds of models have been proposed although none can be deemed as fully fundamental and problem-free. Recently, supergravity models derived from string theory following the original proposal by KKLT [24] and KKLLMT [25] have tried to tackle both inflation and supersymmetry breaking at the same time. Unfortunately, the simplest models require a large gravitino mass compared to the Hubble rate as well as a certain amount of fine-tuning [26]. It is nevertheless compelling that both inflation and supersymmetry breaking can be treated within the same framework.

Recently ISS [2] have proposed that supersymmetry breaking could occur in a long-lived metastable state. One salient point of their analysis is the fact that supersymmetry breaking is a low energy phenomenon occurring in the low energy regime of a SQCD theory. The transmission to the observable sector has been analysed in [8] where gauge mediation

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<sup>10</sup> We are aware that mediation is not mandatory for reheating, see, e.g., [23].

appears as a natural candidate, only relying on low energy physics too. From the effective theory point of view, these facts are very appealing as both supersymmetry breaking and its mediation can be described at low energy. In this paper we extend this analysis to the inflationary sector by requiring that it can also be described by the low energy effective theory of a SQCD theory. More explicitly, we have focused on an inflationary sector with 4 colours and 5 flavours. As usual, the inflation scale is derived from the COBE bound on the CMB anisotropies. As advocated in [9], the scale of supersymmetry breaking can be linked to the inflationary scale. This happens when the two sectors are coupled via gravitational interactions. As a result we find that the supersymmetry breaking scale can be low enough for gauge mediation. We also find that reheating at the end of inflation is obtained via the decay of the inflaton in the quarks of the ISS sector. In conclusion, supersymmetry breaking is mediated from the inflationary sector and the reheating of the universe is due to the coupling of the inflaton with the magnetic quarks in the supersymmetry breaking ISS sector.

Of course, it is also relevant that both the supersymmetry breaking and inflationary sectors have an ultra violet completion. Moreover, this completion appears to be gauge theories with non-Abelian interactions. An interesting extension of our work would be to analyse the embedding of both the ISS sector and the inflation sector in string theory. This should be describable within the brane engineering framework [20, 21]. In such a description, both inflation as a result of brane motion and the interaction between the ISS and the inflation sectors should have a direct interpretation. This is left for future work.

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