

Unified model of D-term inflationValerie Domcke¹ and Kai Schmitz²¹*APC/PCCP, Paris Diderot University, 75013 Paris, France*²*Max-Planck-Institut für Kernphysik (MPIK), 69117 Heidelberg, Germany*

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Hybrid inflation, driven by a Fayet-Iliopoulos (FI) D term, is an intriguing inflationary model. In its usual formulation, it however suffers from several shortcomings. These pertain to the origin of the FI mass scale, the stability of scalar fields during inflation, gravitational corrections in supergravity, as well as to the latest constraints from the cosmic microwave background. We demonstrate that these issues can be remedied if D-term inflation is realized in the context of strongly coupled supersymmetric gauge theories. We suppose that the D term is generated in consequence of dynamical supersymmetry breaking. Moreover, we assume canonical kinetic terms in the Jordan frame as well as an approximate shift symmetry along the inflaton direction. This provides us with a unified picture of D-term inflation and high-scale supersymmetry breaking. The D term may be associated with a gauged $U(1)_{B-L}$, so that the end of inflation spontaneously breaks $B - L$ in the visible sector.

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Cosmic inflation [1] is a successful paradigm in our understanding of the early universe. It is, however, still unclear how to correctly embed inflation into particle physics [2]. One promising ansatz is the idea of hybrid inflation [3], which establishes a connection between inflation and grand unification. Hybrid inflation exits into the subsequent radiation-dominated phase via a *waterfall transition*. In the context of a given grand unified theory (GUT), this phase transition may then be identified with the spontaneous breakdown of a local GUT symmetry.

Depending on the type of GUT symmetry, the waterfall transition may have important consequences for the particles of the standard model (SM). Here, a prominent example is the spontaneous breaking of $U(1)_{B-L}$ [4], i.e., the Abelian gauge symmetry associated with the difference between baryon number B and lepton number L . The end of inflation is then accompanied by the generation of large L -violating Majorana masses for a number of right-handed neutrinos, which sets the stage for the type-I seesaw mechanism [5] as well as for baryogenesis via leptogenesis [6]. Hybrid inflation ending in a $B - L$ phase transition, thus, promises to provide an appealing framework for the early universe that not only determines the initial conditions of the hot thermal phase, but which also explains the smallness of the SM neutrino masses.

In its simplest, nonsupersymmetric form, hybrid inflation predicts the primordial scalar power spectrum to be blue-tilted, which is by now observationally ruled out at a level of more than 5σ [7]. This problem can be avoided in supersymmetry (SUSY), where scalar and fermion loops generate a logarithmic effective potential. Supersymmetric hybrid inflation comes in two variants, depending on whether the vacuum energy density during inflation is dominated by a nonzero F term [8] or D term [9]. In F-term hybrid inflation (FHI), the inflaton field itself has a large

F term during inflation. In combination with R symmetry breaking, this results in a dangerous supergravity (SUGRA) tadpole term [10], which breaks the rotational invariance in field space, generates a false vacuum at large field values, and potentially spoils slow-roll inflation. In D-term hybrid inflation (DHI), the superpotential of the inflationary sector has, by contrast, zero vacuum expectation value (VEV) at all times, so that SUGRA corrections tend to become more manageable. Moreover, DHI is based on a nonzero Fayet-Iliopoulos (FI) D term [11] and, hence, does not require a dimensionful input parameter in the superpotential.

In this paper, we shall construct a consistent SUGRA model in which DHI is driven by the D term associated with a local $U(1)_{B-L}$ symmetry. Despite the absence of the inflaton tadpole term, this is still a difficult task for at least five reasons: (i) The consistent embedding of the FI term into SUGRA is a subtle issue that has been the subject of a long debate in the literature. On the one hand, constant, field-independent FI terms always require an exact global symmetry [12], which conflicts with the expectation that quantum gravity actually does not admit such symmetries [13]. On the other hand, field-dependent FI terms (such as those in string theory [14]) imply the existence of a shift-symmetric modulus field [15], which causes cosmological problems [16], as long as it is not properly stabilized [17]. (ii) The sfermions in the minimal supersymmetric standard model (MSSM) carry nonzero $B - L$ charges and, thus, acquire D-term-induced masses during inflation [18]. Some of these masses are tachyonic and may, hence, destabilize the corresponding directions in scalar field space [19]. (iii) General arguments in SUGRA [20] indicate that a nonzero D term is typically accompanied by a comparatively larger F term, $|F| \gtrsim |D|$. If SUSY breaking is mediated to the visible sector by ordinary gravity mediation [21], this implies that the inflaton picks up a

gravity-mediated soft mass of the order of the gravitino mass, $m_{3/2} \sim |F|/M_{\text{Pl}}$, that necessarily exceeds the Hubble rate during inflation, $H \sim |D|/M_{\text{Pl}}$. DHI in combination with ordinary gravity mediation, therefore, also faces the η slow-roll problem in SUGRA [22]. (iv) In the global-SUSY limit, DHI predicts a scalar spectral index of $n_s \approx 0.98$ [9]. This deviates from the latest value reported by the PLANCK collaboration, $n_s^{\text{obs}} = 0.9677 \pm 0.0060$ [7], by about 2σ . SUGRA corrections may help to reach better agreement with the data [23]. But in general, realizing a spectral index of $n_s \approx 0.96$ in DHI is a nontrivial task. (v) In its standard formulation, DHI is driven by a D term associated with a $U(1)$ symmetry that becomes spontaneously broken only during the waterfall transition at the end of inflation. This results in the production of cosmic strings, which impact the scalar power spectrum of the cosmic microwave background (CMB). The recent CMB bounds on the tension of such cosmic strings [24] severely constrain the parameter space of standard DHI.

We now argue that all of these issues can be remedied as soon as one makes the following three assumptions: (i) The FI term is dynamically generated in the context of dynamical SUSY breaking (DSB) [19]. (ii) DHI is embedded into Jordan-frame SUGRA with canonical kinetic terms for all fields [25]. (iii) There exists an approximate shift symmetry in the direction of the inflaton field [26]. For a more comprehensive account of our idea, see [27].

As far as the generation of the FI term is concerned, we follow the discussion in [19]. We assume that SUSY is broken in a hidden sector by the dynamics of a strongly interacting supersymmetric gauge theory. To this end, we shall employ the Izawa-Yanagida-Intriligator-Thomas (IYIT) model [28] in its $SU(2)$ formulation, i.e., the simplest conceivable DSB model with vectorlike matter fields. At high energies, the IYIT model consists of four quark fields Ψ^i in the fundamental representation of $SU(2)$. At energies below the dynamical scale Λ_0 , these quarks condense into six gauge-invariant meson fields, $M^{ij} \approx \langle \Psi^i \Psi^j \rangle / (\eta^2 \Lambda)$, where $\Lambda \approx \Lambda_0/\eta$ and $\eta \approx 4\pi$ [29]. The scalar mesons span a quantum moduli space of degenerate supersymmetric vacua, subject to a particular constraint on their Pfaffian, $\text{Pf}(M^{ij}) \approx \Lambda^2$ [30]. In order to break SUSY in this model, one couples the high-energy theory to a set of six $SU(2)$ singlets, Z_{ij} , so as to lift the flat directions in moduli space. At high and low energies, the IYIT superpotential respectively reads as follows,

$$W_{\text{hid}}^{\text{HE}} = \frac{1}{2} \lambda_{ij} Z_{ij} \Psi^i \Psi^j \rightarrow W_{\text{hid}}^{\text{LE}} \approx \frac{1}{2} \lambda_{ij} \Lambda Z_{ij} M^{ij}, \quad (1)$$

where λ_{ij} is a matrix of Yukawa couplings. SUSY is now broken à la O’Raifeartaigh [31], as the F-term conditions for the singlet fields Z_{ij} are incompatible with the Pfaffian constraint. A crucial observation for our purposes is that the IYIT model exhibits an axial $U(1)$ flavor symmetry associated with a quark phase rotation, $\Psi^i \rightarrow e^{iq_i \alpha} \Psi^i$. We

shall now identify this symmetry with $U(1)_{B-L}$ and promote it to a weakly gauged local symmetry. In doing so, we suppose that two quarks carry charge $q_0/2$, while the other two carry charge $-q_0/2$. In this case, we have to deal with six mesons (and similarly six singlets) with charges $q_0, -q_0$, and four times 0, respectively. Here, we assign the $B-L$ charges in such a way that the charged mesons, M_{\pm} , have the smallest Yukawa couplings, λ_{\pm} . During SUSY breaking, it is therefore the fields M_{\pm} that acquire nonzero VEVs. The neutral mesons and singlets remain stabilized at their origin. In the weakly gauged limit, one finds (see [18,19,32,33] for more details on the dynamics of the IYIT model and its applications),

$$\langle M_{\pm} \rangle = \frac{\lambda}{\lambda_{\pm}} \Lambda, \quad \lambda = \sqrt{\lambda_+ \lambda_-}. \quad (2)$$

These VEVs break $B-L$ spontaneously, which results in an effective FI term in the $B-L$ D-term scalar potential,

$$V_D = \frac{g^2}{2} [q_0(|M_-|^2 - |M_+|^2) + \dots]^2. \quad (3)$$

Here, g denotes the $B-L$ gauge coupling, while the ellipsis stands for all further fields that are charged under $U(1)_{B-L}$. One then obtains for the FI mass scale ξ ,

$$\xi = \langle M_- \rangle^2 - \langle M_+ \rangle^2 = \frac{2}{\rho^2} (1 - \rho^4)^{1/2} \Lambda^2, \quad (4)$$

where $\rho = [(\lambda_+/\lambda_- + \lambda_-/\lambda_+)/2]^{-1/2}$ is a measure for the degeneracy among the Yukawa couplings λ_+ and λ_- . For $\lambda_+ \approx \lambda_-$, ρ is close to unity; for a strong hierarchy among λ_+ and λ_- , it takes a value close to zero. In the following, we will assume that λ_+ and λ_- are both of the same order of magnitude. Averaging all possible values of ρ under this assumption (varying λ_+ and λ_- on a linear scale) then results in an expectation value of $\langle \rho \rangle \approx 0.80$.

The FI parameter in Eq. (4) is a field-dependent FI term, as it originates from the VEVs of the two meson fields M_{\pm} . The modulus field associated with this FI term is nothing but the $B-L$ Goldstone multiplet, $A = (\langle M_+ \rangle M_+ - \langle M_- \rangle M_-) / f_a$, where f_a is the Goldstone decay constant, $f_a^2 = \langle M_+ \rangle^2 + \langle M_- \rangle^2$. The pseudoscalar in A is absorbed by the massive $B-L$ vector boson, while the real scalar in A is stabilized by an F-term-induced mass, $m_F = \rho \lambda \Lambda$. The same holds true for the fermion in A . Owing to the fact that our FI term is generated in conjunction with dynamical SUSY breaking, we therefore do not face any modulus problem. Our model, thus, avoids the problems described in [12,15].

In the SUSY-breaking vacuum at low energies, the IYIT model effectively reduces to the Polonyi model [34] with an effective superpotential of the following form [33],

$$W_{\text{hid}} = \mu^2 X + w. \quad (5)$$

Here, $\mu^4 = \lambda_+^2 \langle M_+ \rangle^2 \Lambda^2 + \lambda_-^2 \langle M_- \rangle^2 \Lambda^2 = 2\lambda^2 \Lambda^4$ denotes the F-term SUSY-breaking scale, $X = (Z_+ + Z_-)/\sqrt{2}$ is the Polonyi field, and w is an R -symmetry-breaking constant that needs to be added to W_{hid} , so as to achieve zero cosmological constant in the true vacuum.

We now couple the effective Polonyi model in Eq. (5) to SUGRA. In doing so, we shall work in Jordan-frame supergravity with canonical kinetic terms [25]. The total Kähler potential of our theory is therefore given as

$$K_{\text{tot}} = -3M_{\text{Pl}}^2 \ln \left(-\frac{\Omega_{\text{tot}}}{3M_{\text{Pl}}^2} \right), \quad (6)$$

where $\Omega_{\text{tot}} = -3M_{\text{Pl}}^2 + F_{\text{tot}}$ is the frame function of the Jordan frame. We assume that the kinetic function F_{tot} can be split into separate contributions from the hidden, visible, and inflaton sector. Schematically, we may write

$$F_{\text{tot}} = F_{\text{hid}} + F_{\text{vis}} + F_{\text{inf}} + \frac{1}{M_*^2} F_{\text{hid}} F_{\text{vis}}, \quad (7)$$

such that the inflaton sector becomes sequestered from the hidden sector [35]. This serves the purpose to protect the inflaton from a SUGRA mass of the order of $m_{3/2} \gtrsim H$, which would otherwise spoil slow-roll inflation. Meanwhile, the MSSM sfermions do acquire soft masses via gravity mediation. These may be much larger than $m_{3/2}$, provided that the mass scale M_* is parametrically smaller than the reduced Planck mass $M_{\text{Pl}} \approx 2.44 \times 10^{18}$ GeV. In particular, by choosing an appropriate value of M_* , the MSSM sfermions are also sufficiently stabilized during inflation, even if inflation is driven by a $B-L$ D term.

For $F_{\text{hid}} = |X|^2 + [\text{other fields}]$, we obtain for the Polonyi VEV during ($f \neq 0$) and after ($f = 0$) inflation,

$$\langle X \rangle = \frac{(k-4/3)^{1/2}}{[(1-f)k-4/3]k^{1/2}} \frac{2M_{\text{Pl}}}{\sqrt{3}}. \quad (8)$$

Here, f is related to the kinetic function of the inflaton field S , $f = (F_{\text{inf}} - \partial_S F_{\text{inf}} \partial_{S^\dagger} F_{\text{inf}})/(3M_{\text{Pl}}^2)$. k is a ratio of different contributions to the total Polonyi mass,

$$k = [(m_{1\ell}^J)^2 + 2H_J^2] \frac{M_{\text{Pl}}^2}{\mu^4}, \quad (9)$$

which is typically very large, $k \gg 1$. This reflects the fact that X is stabilized by the strong dynamics close to the origin, $\langle X \rangle \ll M_{\text{Pl}}$. H_J is the Hubble rate in the Jordan frame, while $m_{1\ell}^J \approx 0.02\lambda^2 \Lambda$ denotes the effective Polonyi mass in the IYIT model. This mass is generated via meson loops at one-loop level [36]. The Polonyi field is stabilized at $\langle X \rangle$ as given in Eq. (8) only as long as $\langle X \rangle$ does not induce masses for the IYIT quarks larger than Λ_0 . This defines a critical field value, $X_c \approx \sqrt{2}/\lambda \Lambda_0$, above which the Polonyi potential changes from a quadratic to a logarithmic form. The requirement that $\langle X \rangle \lesssim X_c$ then

translates into a lower bound, $\lambda_{\text{min}}(\Lambda, \langle \rho \rangle) \sim 0.1 \dots 1$, on λ . At the same time, we impose an upper bound, $\lambda \lesssim \lambda_{\text{pert}} \approx 4$, so that noncalculable higher-dimensional terms in the Kähler potential, which scale like $\lambda^2/(16\pi^2)$ [36], are suppressed by at least half an order of magnitude. In fact, for definiteness, we will set $\lambda^2 = \lambda_{\text{min}} \lambda_{\text{pert}}$ in the following. Meanwhile, a given value of λ implies a lower bound on ρ , such that neither of λ_\pm becomes larger than λ_{pert} . For our choice of λ , we find $\rho^{-2} < \rho_{\text{pert}}^{-2} = (\lambda_{\text{min}}/\lambda_{\text{pert}} + \lambda_{\text{pert}}/\lambda_{\text{min}})/2$, where ρ_{pert} depends on the dynamical scale Λ via λ_{min} .

The Polonyi field is a linear combination of the charge eigenstates Z_\pm . It hence enters into the D-term potential, where it mixes with the field $Y = (Z_+ - Z_-)/\sqrt{2}$. This mixing destabilizes the vacuum in Eq. (8), unless $|\Delta m_{XY}|^2 = |g^2 q_0^2 \xi| < m_{1\ell}^J m_F$, which translates into an upper bound on g . In our analysis of the inflationary dynamics, we will evaluate the bounds on ρ and g numerically. The lesson from these two bounds is the following: $\rho > \rho_{\text{pert}}$ guarantees that we can neglect nonperturbative corrections to the Kähler potential, while $g < g_{\text{max}}$ ensures a stable vacuum in the SUSY-breaking sector.

The constant w in Eq. (5) needs to be tuned to

$$w_0 = \left(1 - \frac{4}{3k}\right)^{1/2} \frac{\mu^2 M_{\text{Pl}}}{\sqrt{3}}, \quad (10)$$

so as to reach zero cosmological constant in the SUSY-breaking vacuum. We then obtain for $m_{3/2} = \langle W \rangle / M_{\text{Pl}}^2$,

$$m_{3/2} = \left(1 - \frac{4}{3k}\right)^{1/2} \frac{(1-f)k + 2/3}{(1-f)k - 4/3} \frac{\mu^2}{\sqrt{3} M_{\text{Pl}}}. \quad (11)$$

Making use of our choices for the parameters λ and ρ , we arrive at the following phenomenological relation,

$$\frac{m_{3/2}}{10^{11} \text{ GeV}} \sim \left(\frac{\sqrt{\xi}}{10^{15} \text{ GeV}}\right)^{5/2} \sim \left(\frac{\Lambda}{10^{15} \text{ GeV}}\right)^{5/2}, \quad (12)$$

which illustrates that SUSY is broken at a high scale. All mass scales in the IYIT sector are now solely controlled by Λ . This scale is dynamically generated via dimensional transmutation, meaning that our model does not require any hard dimensionful input scale. Equation (12) sets the stage for our particular implementation of DHI.

We consider the following contributions to the kinetic function and superpotential from the inflationary sector,

$$F_{\text{inf}} = \frac{\chi}{2} (S^\dagger + S)^2 - \frac{1}{2} (1 - \chi) (S^\dagger - S)^2 + \Phi^\dagger \Phi + \bar{\Phi}^\dagger \bar{\Phi},$$

$$W_{\text{inf}} = \kappa S \Phi \bar{\Phi}. \quad (13)$$

Taking $|\chi| \ll 1$, the real scalar component σ of S acquires an approximate shift symmetry and will play the role of the

inflaton. The so-called waterfall fields $\Phi, \bar{\Phi}$ carry $B-L$ charges $\pm q$. Due to their σ -dependent mass spectrum, one of the scalar waterfall degrees of freedom becomes tachyonic at $\sigma \leq \sigma_c$, leading to a phase transition that ends inflation. The sequestered structure of Eq. (7) protects the waterfall fields from acquiring soft SUGRA masses of the order of $m_{3/2}$, which could prevent this phase transition. Note that since $B-L$ is already broken by $\langle M_{\pm} \rangle \neq 0$ during inflation, the production of cosmic strings at the end of inflation can be avoided [19].

The scalar inflaton potential at $\sigma \gg \sigma_c$ is given by

$$V(\sigma) \simeq \mathcal{C}^4(\sigma) \left[V_D^J + V_F^J(\sigma) + \frac{Q_J^4(\sigma)}{16\pi^2} \ln(x(\sigma)) \right], \quad (14)$$

with the tree-level D- and F-term contributions being

$$V_D^J = \frac{1}{2} q_0^2 g^2 \xi^2, \quad V_F^J(\sigma) = \frac{-f(\sigma)}{1-f(\sigma)} \mu^4. \quad (15)$$

Similarly as in global SUSY, inflation is driven by the constant D-term potential induced by the FI term (4). V_F^J arises due to F-term SUSY breaking in the hidden sector and vanishes in the true vacuum at $\sigma = 0$. The function $f = (1-2\chi)\chi\sigma^2/(3M_{\text{Pl}}^2) \ll 1$ is introduced below Eq. (8). The third term in Eq. (14) is the effective one-loop potential arising from integrating out the waterfall multiplets. Here, the renormalization scale $Q_J^4 = q^2 m_D^4 + \delta m^4$ and $x = (m_{\text{eff}}^2 + 2H_J^2)/Q_J^2$ are given in terms of the various contributions to the masses of the waterfall fields: $m_D^2 = g^2 q_0 \xi$ is induced by the D term, $m_{\text{eff}}^2 = \kappa^2 \sigma^2/2$ follows from the coupling in W_{inf} , and $\delta m^2 \simeq m_{3/2} m_{\text{eff}}$ is a bilinear soft mass in consequence of R symmetry breaking. The critical inflaton field value σ_c is obtained once $m_{\text{eff}}^2 + 2H_J^2 - Q_J^2 = 0$, which results in $\sigma_c^2 = 2/(\kappa^2(Q_J^2 - 2H_J^2))$. The conformal factor $\mathcal{C}^2 = -3M_{\text{Pl}}^2/\Omega_{\text{tot}}$ translates the Jordan-frame potentials to their counterparts in the Einstein frame.

Integrating out the heavy $B-L$ gauge fields results in additional gauge-mediated soft masses, m_{gm} , for the waterfall fields [37] (see also [18]). However, at the level of the effective one-loop potential for the inflaton, these radiative corrections merely represent a two-loop effect. The inflaton itself, being a gauge singlet, receives by contrast no gauge-mediated soft mass. Moreover, one can show that, in the parameter range of interest, m_{gm} is always outweighed by the tree-level mass induced by the D-term potential, $m_{\text{gm}} \lesssim 0.03 m_D$. For these reasons, we will neglect the effect of gauge mediation in the following.

We solve the slow-roll equation of motion numerically,

$$K_{S^i S^j}(\sigma) V(\sigma) \sigma'(N) = M_{\text{Pl}}^2 V'(\sigma), \quad (16)$$

to obtain the predictions for the CMB observables at $N_* \simeq 55$ e-folds before the end of inflation. With

$$\varepsilon = \frac{M_{\text{Pl}}^2}{2} \left(\frac{V'(\hat{\sigma})}{V} \right)^2, \quad \eta = M_{\text{Pl}}^2 \frac{V''(\hat{\sigma})}{V}, \quad (17)$$

where derivatives with respect to the canonically normalized field $\hat{\sigma}$ can be obtained by $\partial \hat{\sigma} / \partial \sigma = \sqrt{K_{S^i S^i}}$, the amplitude of the scalar perturbation spectrum, its tilt and the tensor-to-scalar ratio are obtained as

$$A_s = \frac{V}{24\pi^2 \varepsilon M_p^4}, \quad n_s = 1 - 6\varepsilon + 2\eta, \quad r = 16\varepsilon, \quad (18)$$

evaluated at $\sigma(N_*)$. Requiring $A_s = A_s^{\text{obs}} = 2.1 \times 10^{-9}$ [7] fixes Λ (or equivalently ξ). The parameter $\kappa \neq 0$ explicitly breaks the shift symmetry in the superpotential, which leads us to expect that $\kappa \lesssim 1$. On the other hand, for $\kappa \ll 1$, the correct spectral index can only be obtained if the SUGRA contributions become much larger than the one-loop contributions [27]. We thus set $\kappa = 0.1$. In this regime, inflation occurs at field values slightly below the Planck scale. For simplicity, we also fix the $B-L$ charges to $q = 2q_0 = -2$, inspired by neutrino mass generation (see below). We depict our results in the remaining (χ, g) plane in Fig. 1. r is of $\mathcal{O}(10^{-6} \dots 10^{-4})$, which is, similarly as in FHI, far below current bounds.

These results are very well reproduced by approximate analytical expressions for the slow-roll parameters,

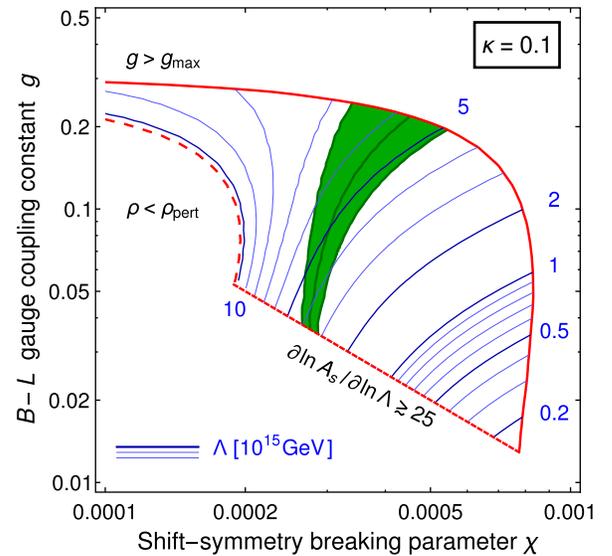


FIG. 1. Theoretical and experimental constraints. The green band indicates values of the spectral index in agreement (at 2σ) with the Planck 2015 data, $n_s = 0.9677 \pm 0.006$ [7]. The blue contours indicate the values of the dynamical scale Λ in the IYIT sector, such that $A_s = A_s^{\text{obs}}$. The red lines are boundaries of the parameter space due to the requirement of perturbativity ($\rho > \rho_{\text{pert}}$) and a stable vacuum ($g < g_{\text{max}}$) in the SUSY-breaking sector as well as not too large fine-tuning in the parameters of the inflation sector ($\partial \ln A_s / \partial \ln \Lambda \lesssim 25$).

$$\varepsilon \simeq \left(\frac{M_{\text{Pl}}}{\sigma/\sqrt{2}} \right)^2 \left[(1 + \delta_\varepsilon^4) \frac{q^2 g^2 D_0^2}{16\pi^2 V_D^J} - f \frac{F_0^2}{V_D^J} \right]^2, \quad (19)$$

$$-\eta \simeq \left(\frac{M_{\text{Pl}}}{\sigma/\sqrt{2}} \right)^2 \left[(1 - \delta_\eta^4) \frac{q^2 g^2 D_0^2}{16\pi^2 V_D^J} + f \frac{F_0^2}{V_D^J} \right], \quad (20)$$

with $F_0^2 = \mu^4 \gtrsim D_0^2 = q_0^2 g^2 \xi^2$ and

$$(\delta_\varepsilon/\delta)^4 = \ln x + \frac{1}{2}, \quad (\delta_\eta/\delta)^4 = \ln x + \frac{1}{2} + \frac{2 + \delta^4}{1 + \delta^4}. \quad (21)$$

Here, $\delta^2 = \delta m^2/(q m_D^2)$ parametrizes the effect of the soft B-term mass, δm , in the one-loop potential. Lowering the spectral index compared to DHI in global SUSY becomes possible due to the negative mass-squared induced by the SUGRA contributions to the tree-level F-term potential, reflected by the last term in Eq. (20),

$$\Delta\eta \simeq -\frac{2\chi}{3} \left(\frac{m_{3/2}}{H_J} \right)^2, \quad H_J^2 \simeq \frac{V_D^J}{3M_{\text{Pl}}^2}. \quad (22)$$

Successful inflation is thus due to the interplay of the one-loop contribution and the SUGRA-induced mass, with the latter being suppressed by an approximate shift symmetry, $\chi \sim 10^{-4}$. We note that $F_{\text{inf}} \supset \chi \sigma^2$ in Eq. (13) might, e.g., arise from further shift-symmetry breaking terms in the superpotential. Suppose the inflaton couples to superheavy multiplets with strength κ' . Integrating out these fields results in an effective Kähler potential, $K_{1\ell} \sim \kappa'^2/(16\pi^2)|S|^2$ [38]. With $\kappa' \sim \kappa \sim 0.1$, this is of just the right order to explain the required value of χ .

In the viable region of parameter space, inflation occurs either near a hill-top (i.e., a local maximum in the scalar potential) or near an inflection point, depending on the exact values of χ and g [27]. The hill-top regime may suffer from an initial conditions problem. For particular parameter values, there is however a false vacuum at large field values. From there, σ could tunnel to the correct side of the hill-top, thereby setting off inflation in our Universe. The inflection-point regime allows, by contrast, to start out at super-Planckian field values.

In FHI, $n_s \simeq 0.96$ is obtained from the interplay of the one-loop potential and the SUGRA tadpole, which is *linear* in the inflaton field. The tadpole also renders the question of initial conditions more subtle [10]. Its size is controlled by $m_{3/2}$, an independent parameter, which can be chosen in accord with low-scale SUSY breaking.

Let us conclude. We have presented a complete and phenomenologically viable SUGRA model of DHI, in which inflation is driven by the D term of a gauged $U(1)_{B-L}$ symmetry. Our model unifies the dynamics of dynamical SUSY breaking in the hidden sector, DHI, and spontaneous $B-L$ breaking. It links all relevant energy scales to the dynamical scale in the hidden sector, the magnitude of which is fixed by the amplitude of the CMB

power spectrum, $\Lambda \simeq 5 \times 10^{15}$ GeV. This value is remarkably close to the GUT scale, $\Lambda_{\text{GUT}} \sim 10^{16}$ GeV.

We based our construction on three assumptions: dynamical SUSY breaking as the origin of the FI term, Jordan-frame supergravity with canonical kinetic terms, and an approximate inflaton shift symmetry. These assumptions remedy all shortcomings of standard DHI: Our FI term is a field-dependent FI term. The associated modulus is stabilized via F-term SUSY breaking. The MSSM sfermions are stabilized against tachyonic D-term-induced masses thanks to their direct coupling to the hidden sector in the kinetic function F_{tot} (7). Owing to our Jordan-frame description, the inflaton sector sequesters from the hidden sector, such that the fields in the inflation sector pick up no dangerous gravity-mediated soft masses. Meanwhile, a slight breaking of the shift symmetry provides a small SUGRA correction to the inflaton potential, $V \supset -\chi m_{3/2}^2 \sigma^2$, that allows us to reproduce $n_s \simeq 0.96$. As $B-L$ is spontaneously broken in the hidden sector already during inflation, no dangerous cosmic strings are produced during the waterfall transition.

Our model has important phenomenological consequences. For instance, if we assign $B-L$ charge $q = -2$ to the waterfall field Φ , it can couple to the right-handed neutrinos N_i in the seesaw extension of the MSSM, $W \supset h_{ij}/2\Phi N_i N_j$. For $q_0 \xi < 0$, it is the field Φ that acquires a nonzero VEV during the waterfall transition, $\langle \Phi \rangle = |q_0/q\xi|^{1/2}$, whereas $\langle \bar{\Phi} \rangle$ remains zero. This VEV generates the Majorana mass matrix for the right-handed neutrinos, $M_{ij} = h_{ij}\langle \Phi \rangle$, and, hence, sets the stage for the seesaw mechanism and leptogenesis [4]. Besides, our model predicts a superheavy SUSY mass spectrum. Only the lightest neutralino may have a fine-tuned small mass, so as to evade overproduction in gravitino decays [39]. Together with the corresponding chargino, this neutralino is then expected to be the only superparticle at low energies. It constitutes thermal neutralino dark matter and can be searched for in direct detection experiments.¹ One may also hope to probe

¹In our model, the mediation of spontaneous SUSY breaking to the visible sector is essentially described by the framework of pure gravity mediation (PGM) [40]. In this mediation scheme, one is, e.g., able to achieve a light wino mass by tuning Higgsino threshold corrections against the usual gaugino masses from anomaly mediation [41]. For $\tan\beta = v_u/v_d \simeq 1$ and Higgs mass parameters of the order of the gravity mass, $|\mu_H| \simeq |B| \simeq m_{3/2}$, PGM readily allows us to push the wino mass down to $m_{\text{wino}} \simeq 2.7$ TeV, so that it may constitute ordinary thermal dark matter in the form of weakly interacting massive particles (WIMPs). This fine-tuning might be the result of anthropic selection. In addition, the wino mass may also receive threshold and anomaly-mediated corrections from heavy vector matter multiplets charged under $SU(2)_L$ and/or threshold corrections from the F terms of flat directions in Kim-Shifman-Vainshtein-Zakharov-type axion models [42]. In this case, these contributions would also play a role in tuning the wino mass.

our model in gravitational-wave (GW) experiments. Depending on further model assumptions, the $B-L$ phase transition may give rise to observable signals [43]. Likewise, if the shift symmetry is realized along the imaginary component of S , the inflaton may have an axion-like coupling to gauge fields, $\mathcal{L}_{\text{eff}} \supset \text{Im}(S)F\tilde{F}$. This could drastically enhance the GW signal from inflation [44].

Our model leaves open several questions that call for further exploration: For instance, one may ask what UV physics underlies the kinetic function in Eq. (7). It would be interesting to derive this structure from the viewpoint of a higher-dimensional brane-world scenario or from a strongly coupled conformal field theory [35]. We have only briefly sketched the mechanism of sfermion mass generation. It would, therefore, be desirable to devise a model that accounts for the origin of the scale M_* in Eq. (7). These questions are however beyond the scope of this

work. We conclude by stressing that our dynamical SUGRA model resolves all issues of standard DHI.

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