Dynamical supersymmetry breaking and late-time *R* symmetry breaking as the origin of cosmic inflation

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Spontaneously broken supersymmetry (SUSY) and a vanishingly small cosmological constant imply that *R* symmetry must be spontaneously broken at low energies. Based on this observation, we suppose that, in the sector responsible for low-energy *R* symmetry breaking, a discrete *R* symmetry remains preserved at high energies and only becomes dynamically broken at relatively late times in the cosmological evolution, i.e., after the dynamical breaking of SUSY. Prior to *R* symmetry breaking, the Universe is then bound to be in a quasi–de Sitter phase—which offers a dynamical explanation for the occurrence of cosmic inflation. This scenario yields a new perspective on the interplay between SUSY breaking and inflation, which neatly fits into the paradigm of high-scale SUSY: inflation is driven by the SUSY-breaking vacuum energy density, while the chiral field responsible for SUSY breaking, the Polonyi field, serves as the inflaton. Because *R* symmetry is broken only after inflation, slow-roll inflation is not spoiled by otherwise dangerous gravitational corrections in supergravity. We illustrate our idea by means of a concrete example, in which both SUSY and *R* symmetry are broken by strong gauge dynamics and in which late-time *R* symmetry breaking is triggered by a small inflaton field value. In this model, the scales of inflation and SUSY breaking are unified, the inflationary predictions are similar to those of F-term hybrid inflation in supergravity, reheating proceeds via gravitino decay at temperatures consistent with thermal leptogenesis, and the sparticle mass spectrum follows from pure gravity mediation. Dark matter consists of thermally produced winos with a mass in the TeV range.

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I. INTRODUCTION: INFLATION AND SUPERSYMMETRY BREAKING UNIFIED

A. High-scale supersymmetry breaking as the origin of inflation

The paradigm of cosmic inflation [1,2] is one of the main pillars of modern cosmology. Not only does inflation account for the vast size of the observable Universe and its high degree of homogeneity and isotropy on cosmological scales, it also seeds the postinflationary formation of structure on galactic scales. In this sense, inflation is a key aspect of our cosmic past and part of the reason why our Universe is capable of harboring life. From the perspective of particle physics, the origin of inflation is, however, rather unclear. After decades of model building, there exists a plethora of inflation models in the literature [3]. But a consensus about how to embed inflation into particle physics is out of sight. In this situation, it seems appropriate to take a step back and ask ourselves what avenues have been left unexplored so far. In particular, we should question our dearly cherished prejudices and re-examine whether inflation might not be connected to other high-energy phenomena which, up to now, have been taken to be mostly unrelated to inflation. As we are going to demonstrate in this paper, an important example in this respect might be the interplay between inflation and the spontaneous breaking of supersymmetry (SUSY).[1]

In recent years, the picture of supersymmetry as a solution to the hierarchy problem has become increasingly challenged by the experimental data. The null results of SUSY searches at the Large Hadron Collider (LHC) [8] and the rather large standard model (SM) Higgs boson mass of 125 GeV [9] indicate that supersymmetry, if it exists in nature, must be broken at a high scale [10]. Based on this observation, one could feel tempted to give up on supersymmetry as an extension of the standard model altogether. But this would not do justice to supersymmetry's other virtues. Provided that supersymmetry is broken at a high scale [11,12], such as in the minimal framework of pure gravity mediation (PGM) [13,14],[2] it may no longer be responsible for stabilizing the electroweak scale. But in this case, supersymmetry is still capable of providing a viable candidate for dark matter [14,16,17], ensuring the unification of the SM gauge couplings [18] and setting the stage for a UV completion of the standard model in the context of string theory. In addition, high-scale

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[1]For related earlier work on the relation between inflation and supersymmetry breaking, see, e.g., [4–7].

[2]For closely related schemes for the mediation of supersymmetry breaking to the visible sector, see [15,16].

supersymmetry does not suffer from a number of phenomenological problems that low-scale realizations of supersymmetry breaking are plagued with. A high SUSY breaking scale does away with the cosmological gravitino problem [19] and reduces the tension with constraints on flavor-changing neutral currents and *CP* violation [20]. Moreover, in PGM, the SUSY-breaking (or "Polonyi") field is required to be a nonsinglet [21], which solves the cosmological Polonyi problem [22].

In this paper, we will now concentrate our attention to yet another intriguing feature of supersymmetry which comes into reach, once we let go of the notion that supersymmetry's main purpose is to solve the hierarchy problem in the standard model. The spontaneous breaking of supersymmetry at a scale $\Lambda_{\rm SUSY}$ results in a nonzero contribution to the total vacuum energy density, $\Lambda_{\rm SUSY}^4$. If we allow $\Lambda_{\rm SUSY}$ to take values as large as, say, the unification scale, $\Lambda_{\rm GUT} \sim 10^{16}$ GeV, this SUSY breaking vacuum energy density might, in fact, be the origin of the inflationary phase in the early Universe. Such a connection between inflation and supersymmetry breaking not only appears economical, but also very natural.

First of all, supersymmetry tends to render inflation technically more natural, independent of the scale at which it is broken. Thanks to the SUSY nonrenormalization theorem [23], the superpotential W in supersymmetric models of inflation does not receive any radiative corrections in perturbation theory. This represents an important advantage in preserving the required flatness of the inflaton potential. Besides, all remaining radiative corrections (which can be collected in an effective Kähler potential K to leading order [24]) scale with the soft SUSY-breaking mass scale [25] and are, thus, under theoretical control. Supersymmetry, therefore, has the ability to stabilize the inflaton potential against radiative corrections, and it is, thus, conceivable that supersymmetry's actual importance may lie in the fact that it is capable of taming the hierarchy among different mass scales in the inflaton sector rather than in the standard model. Second of all, the spontaneous breaking of global supersymmetry via nonvanishing F-terms, i.e., via the O'Raifeartaigh mechanism [26], always results in a pseudoflat direction in the scalar potential [27]. Together with the constant vacuum energy density $\Lambda_{\rm SUSY}^4$, such a flat potential for a scalar field is exactly one of the crucial requirements for the successful realization of an inflationary stage in the early Universe. In principle, the necessary ingredients for inflation are, therefore, already intrinsic features of every O'Raifeartaigh model. Inflation may be driven by the SUSY-breaking vacuum energy density $\Lambda_{\rm SUSY}^4$ and the inflaton field may be identified with the pseudoflat direction in the scalar potential.

The main obstacle in implementing this idea in realistic models is gravity. Here, the crucial point is that the vanishingly small value of the cosmological constant (CC) tells us that we live in a near-Minkowski vacuum with an almost zero total vacuum energy density, $\langle V\rangle \simeq 0$. Note that, as pointed out by Weinberg, this not a mere observation, but a necessary condition for a sufficient amount of structure formation in our Universe, so that it can support life [28]. In the context of supergravity (SUGRA) [29], the fact that $\langle V\rangle \simeq 0$ means that the SUSY-breaking vacuum energy density $\Lambda_{\rm SUSY}^4$ must be balanced, with very high precision, by a nonvanishing vacuum expectation value (VEV) of the superpotential, $\langle W\rangle$,

$$\langle V\rangle = \langle |F|\rangle^2 - 3\exp\left[\frac{\langle K\rangle}{M_{\rm Pl}^2}\right]\frac{\langle |W|\rangle^2}{M_{\rm Pl}^2} \simeq 0,$$
$$\langle |F|\rangle = \Lambda_{\rm SUSY}^2, \tag{1}$$

where $M_{\rm Pl} = (8\pi G)^{-1/2} \simeq 2.44\times 10^{18}$ GeV denotes the reduced Planck mass. If the SUSY breaking scale $\Lambda_{\rm SUSY}$ is indeed of $\mathcal{O}(10^{16})$ GeV, the requirement of a zero CC results in a huge VEV of the superpotential, which, in turn, generates dangerously large SUGRA corrections to the scalar potential. These corrections then easily spoil the flatness of the potential and render inflation impossible [30].

The most attractive and, in fact, only way out of this problem is R symmetry. Under R symmetry, the superpotential carries charge 2, so that $\langle W\rangle = 0$, as long as R symmetry is preserved [31]. In other words, by imposing R symmetry, we promote $\langle W\rangle$ to the order parameter of spontaneous R symmetry breaking, in analogy to $\langle F\rangle$ which acts as the order parameter of spontaneous supersymmetry breaking. In the true vacuum, R symmetry must be broken, so that $\langle W\rangle \neq 0$, in order to satisfy the condition in Eq. (1). But this does not necessarily mean that $\langle W\rangle$ must be nonzero during the entire cosmological evolution. It is conceivable that, at early times, R symmetry is, in fact, a good symmetry, so that $\langle W\rangle \approx 0$, thereby "switching off" the most dangerous SUGRA corrections to the Polonyi potential. The main intention of our paper now is to present a minimal dynamical model, in which this is indeed the case, so that inflation driven by the SUSY-breaking vacuum energy density, i.e., Polonyi inflation, becomes a viable option.

B. Two avenues towards a vanishing cosmological constant

Let us now outline our general philosophy in, to some extent, simplified terms. We suppose that, from the perspective of the low-energy effective theory, the breaking of supersymmetry and R symmetry are two distinct dynamical processes, taking place in two different hidden sectors at two different times, $t_{\rm SUSY}$ and t_R. In particular, we assume that the parameters $t_{\rm SUSY}$ and t_R may be sampled in the UV theory, so that in some patches of the early Universe (or in the landscape of string vacua [32] for that purpose) $t_{\rm SUSY} < t_R$ and in other patches $t_{\rm SUSY} > t_R$. Note that this differentiation equally includes the case of a maximally symmetric initial state, $\min\{t_{\rm SUSY}, t_R\} > t_{\rm ini}$, as well as the possibility that either supersymmetry or R symmetry is already broken from the very beginning, $\min\{t_{\rm SUSY}, t_R\} \equiv t_{\rm ini}$. Our crucial observation is that regions in space where R symmetry is broken before supersymmetry correspond to bubbles of anti–de Sitter (AdS) space with an AdS radius $R_{\rm AdS}$ equal to the inverse of the gravitino mass $m_{3/2}$, while regions in space where supersymmetry is broken before R symmetry correspond to bubbles of de Sitter (dS) space with a dS radius $R_{\rm dS}$ equal to the inverse of the Hubble parameter H,

$$\begin{aligned} t_{\rm SUSY} > t_R &\Rightarrow \text{AdS bubbles with radius } R_{\rm AdS} = m_{3/2}^{-1}, \\ m_{3/2} &= \exp\left[\langle K\rangle/2/M_{\rm Pl}^2\right]\langle|W|\rangle/M_{\rm Pl}^2, \\ t_{\rm SUSY} < t_R &\Rightarrow \text{dS bubbles with radius } R_{\rm dS} = H^{-1}, \\ H &= \langle|F|\rangle/\sqrt{3}/M_{\rm Pl}. \end{aligned} \tag{2}$$

As long as supersymmetry is unbroken, the AdS bubbles represent open Friedmann-Lemaître-Robertson-Walker (FLRW) universes with a negative CC and an oscillatory scale factor. The dS bubbles, on the other hand, turn, as long as R symmetry is unbroken, into (asymptotically) flat FLRW universes with a positive CC and an exponentially growing scale factor [33]. This is to say that the dS bubbles experience inflation, while the AdS bubbles remain limited in their spatial extent. Furthermore, if we suppose that there is no other source of inflation present in the theory, this means that the AdS bubbles will never develop into habitable universes. Intelligent observers and humans can, therefore, only live in regions where, initially, R symmetry remains unbroken up to a certain high energy scale. This applies, in particular, to our own observable Universe. Under the above assumptions, the inflationary period in the early history of our Universe must have been a consequence of spontaneous supersymmetry breaking and R symmetry must have been broken only at late times, i.e., after a sufficient amount of inflation. Put differently, we can say that supersymmetry breaking and Weinberg's argument regarding the size of the CC postdict a period of inflation and late-time R symmetry breaking in our cosmic past.

This conclusion also sheds new light on the role of the CC itself. The fine-tuning of the CC is now a dynamical process that takes place only after inflation. In order to obtain zero CC, we have to require that the gravitino mass generated during R symmetry breaking matches the inflationary Hubble rate,[3]

$$\langle V\rangle \simeq 0 \Rightarrow m_{3/2} \simeq H_{\rm inf}. \tag{3}$$

If R symmetry breaking results in a gravitino mass smaller than $H_{\rm inf}$, inflation never ends; if it "overshoots" and the gravitino mass is eventually larger than $H_{\rm inf}$, our Universe becomes AdS. We, thus, recognize the requirement that inflation must terminate at one point or another as part of the reason why the CC in our Universe is fine-tuned. A certain amount of fine-tuning during late-time R symmetry breaking is inevitable, as inflation would otherwise not exit into a near-Minkowski vacuum. This situation needs to be contrasted with standard scenarios of inflation, such as chaotic [34] or hybrid [35] inflation, where the vacuum energy density driving inflation is neither related to R symmetry breaking nor to low-energy supersymmetry breaking. These scenarios require an independent reason for the vanishing of the inflationary vacuum energy density (e.g., a tuning of the inflaton potential or some kind of waterfall transition), whereas in our case this reason is already inherent to the fine-tuning of the CC in the course of spontaneous R symmetry breaking at the end of inflation.

C. Ingredients for a realistic model of Polonyi inflation

The above reasoning is just a rough sketch. To construct a realistic model of Polonyi inflation, we need to be more specific. This pertains, first of all, to the kind of R symmetry that we have in mind. Naively, our first attempt might be to protect $\langle W\rangle$ by means of a global $U(1)_R$ symmetry. On general grounds, quantum gravity is, however, expected to explicitly break all global symmetries (see [36] and references therein), so that a global $U(1)_R$ does not appear to be a viable possibility. Meanwhile, gauging a continuous R symmetry is a subtle issue that easily results in conflicts with anomaly constraints at low energies (see [37] for a recent discussion). This leaves us with a *discrete* gauged (i.e., anomaly-free [38]) R symmetry, Z_N^R, as a unique choice to ensure the vanishing of the superpotential at early times.

Interestingly enough, such a discrete R symmetry readily comes with a number of other advantages in the context of SUSY phenomenology: (i) A discrete R symmetry prevents too-rapid proton decay via perilous dimension-5 operators [39], (ii) it may give rise to an accidental approximate global Peccei-Quinn symmetry and, thus, help in solving the strong CP problem [40,41], and (iii) it may account for the approximate global continuous R symmetry which is required to realize stable [42] or metastable [43] SUSY-breaking vacua in a large class of models of dynamical supersymmetry breaking (DSB). Moreover, if we restrict ourselves to the special case of a Z_4^R symmetry, R symmetry can also help us in solving the μ problem [44] in the minimal supersymmetric standard model (MSSM). Any discrete R symmetry suppresses the bilinear Higgs mass term (i.e., the μ term) in the superpotential. Only in the case

[3]Note that the relation between $m_{3/2}$ and $H_{\rm inf}$ in our scenario is conceptually quite different from other inflation models, such as, e.g., the one in [4], where $H_{\rm inf} \simeq m_{3/2}$ on purely phenomenological grounds. In our case, $H_{\rm inf} \simeq m_{3/2}$ is a property of the low-energy effective Lagrangian of our Universe that ensures that we live in a vacuum with an almost zero CC.

of a Z_4^R symmetry, we are, in addition, allowed to include a Higgs bilinear term in the Kähler potential, $K \supset H_u H_d$ (see [41] and references therein for an extended discussion of this point). We are then able to generate the μ term in the course of R symmetry breaking [45], which directly relates the μ parameter to the gravitino mass.[4] It is for this reason that we will assume a discrete Z_4^R symmetry in the following.

For our purposes, it will not be necessary to specify the origin of this Z_4^R symmetry. But it is interesting to note that orbifold compactifications of the heterotic string have the ability to yield discrete R symmetries in the low-energy effective theory. In this case, the discrete R symmetry at low energies is nothing but a remnant of the higher-dimensional Lorentz symmetry that survives the compactification of the internal space (for early as well as more recent work on this topic, see [47] and [48], respectively). We also mention that the special case of a discrete Z_4^R symmetry has received particular attention in the context of orbifold compactifications in recent years [49]. The assumption of a discrete Z_4^R symmetry in the low-energy effective theory is, therefore, well motivated and stands theoretically on a sound footing.

Second of all, inflation is more than just a pure dS phase. It represents a stage of quasi-dS expansion, in the course of which the Hubble parameter slowly varies. We, thus, need to specify the dynamics of supersymmetry breaking more precisely and check whether the corresponding Polonyi potential is, in fact, suitable for slow-roll inflation. Here, we shall work within the framework of dynamical supersymmetry breaking [50], in which supersymmetry is assumed to be broken by the dynamics of a strongly coupled SUSY gauge theory. This gives us the advantage that the SUSY breaking scale $\Lambda_{\rm SUSY}$ is generated dynamically via dimensional transmutation. As far as supersymmetry breaking is concerned, we will, therefore, not have to rely on any dimensionful input parameters. Instead, $\Lambda_{\rm SUSY}$ and, thus, the energy scale of inflation, $\Lambda_{\rm inf}$, will be controlled by the dynamical scale Λ of the SUSY-breaking hidden sector,

$$\Lambda_{\rm SUSY} \equiv \Lambda_{\rm inf} \sim \Lambda. \tag{4}$$

In this sense, our inflation model should be regarded as a variant of dynamical inflation [51–53], as we assume the energy scale of inflation to be generated by strong dynamics. Similarly, our model is closely related to natural inflation [54], which treats the inflaton as an axionlike field that is likewise subject to a scalar potential generated by nonperturbative dynamics (for recent dynamical implementations of natural inflation in field theory, see [55]), as well as to modulus inflation [56], where one identifies the inflaton as a (composite) modulus in the effective low-energy regime of strongly coupled gauge theories.

In our case, the role of the inflaton is played by the scalar component of the chiral Polonyi field Φ, which breaks supersymmetry via its nonzero F-term. In global supersymmetry and at the classical level, the scalar Polonyi potential is exactly flat and, thus, an ideal starting point for the realization of inflation. At the quantum level and in supergravity, the Polonyi potential, however, receives corrections, which may or may not spoil the flatness of the potential. Here, the SUGRA corrections lead, in particular, to the notorious η problem [57], which typically requires a parameter fine-tuning at the level of 1…10% or so. The quantum and gravity corrections to the effective scalar potential scale with the coupling strengths of the Yukawa interactions between the Polonyi field and matter fields, λ, as well as with the coefficients of the higher-dimensional Polonyi terms in the effective Kähler potential, ϵ, respectively. In order to assess the prospects of successful Polonyi inflation, one, therefore, has to study the viability of inflation as a function of the parameters λ and ϵ and identify those parameter ranges that lead to consistency with the observational data on the cosmic microwave background (CMB) [58].

D. Our setup: Minimal model based on two strongly coupled $SU(2)$ gauge theories

To this end, we will present in this paper a *minimal* realization of the idea of Polonyi inflation. Our model is based on two strongly coupled $SU(2)$ hidden gauge sectors (featuring two quark/antiquark pairs each), which we take to be responsible for the dynamical breaking of supersymmetry and R symmetry, respectively. We supplement both sectors with an appropriate number of gauge singlet fields, so that the SUSY-breaking sector becomes identical to the simplest version of the Izawa-Yanagida-Intriligator-Thomas (IYIT) DSB model [59], while the R-symmetry-breaking sector turns into a strongly coupled SQCD theory with a quantum mechanically deformed moduli space and field-dependent quark masses [60]. More precisely, we assume the quark masses in the R-symmetry-breaking sector to be controlled by the VEV of a singlet, which we call P.

Inflation is then driven by the SUSY-breaking vacuum energy density in the IYIT sector, which results in an inflaton potential equivalent to that of supersymmetric F-term hybrid inflation (FHI) [61], including corrections from supergravity [62] as well as from higher-dimensional terms in the tree-level Kähler potential [63,64]. This is reminiscent of the inflation models presented in [6] and [51]. The model in [51], however, corresponds to a reduced version of the IYIT model with fewer singlets, which leads to the vanishing of the inflationary vacuum energy at the end of inflation. Contrary to our approach, it,

[4]This solution to the μ problem is not to be confused with the Giudice-Masiero mechanism [46], which relates the generation of the μ term to the spontaneous breaking of supersymmetry rather than to R symmetry breaking.

thus, does not establish a connection between inflation and supersymmetry breaking. Meanwhile, the model in [6] identifies a different singlet field, other than the Polonyi field, as the inflaton. This allows the author of [6] to separate the scales of inflation and supersymmetry breaking by imposing a hierarchy among the Yukawa couplings of the inflaton and the SUSY-breaking field. We, on the other hand, will show how to implement inflation into the full IYIT model, sticking to the notion that the best motivated inflaton candidate in the IYIT model is still the Polonyi field itself. An important consequence of this approach is that, in our case, F-term hybrid inflation does *not* end in a waterfall transition in the inflationary sector. Instead, we simply retain the inflationary vacuum energy density at low energies, which then continues to act as the vacuum energy density associated with the spontaneous breaking of supersymmetry. This also means that our "waterfall-transition-free" scenario of F-term hybrid inflation does not suffer from the usual production of topological defects, such as cosmic strings in the case of a $U(1)$ waterfall transition, which would otherwise exert some serious phenomenological pressure on our model (see, e.g., [65]).

We assume that the SUSY-breaking sector and the R-symmetry-breaking sector only communicate with each other via interactions in the Kähler potential (i.e., not via interactions in the superpotential). This is sufficient to stabilize the scalar field P during inflation at $\langle P\rangle = 0$ by means of its Hubble-induced mass. During inflation, the (fermionic) quarks in the R-symmetry-breaking sector are, therefore, massless and the discrete Z_4^R symmetry remains unbroken in this sector. As anticipated, the superpotential then lacks a constant term during inflation, which relieves the inflationary dynamics from the most dangerous gravitational corrections in supergravity. In particular, the inflaton potential is free of the notorious "tadpole term" linear in the inflaton field [7,65,66]. Towards the end of inflation, the Hubble-induced mass of the scalar field P decreases. Adding an appropriately chosen superpotential for the field P, we can use this fact to trigger a waterfall transition in the R-symmetry-breaking sector at small inflaton field values. The field P then acquires a large VEV and the quarks in the R-symmetry-breaking sector become very massive. Consequently, the R-symmetry-breaking sector turns into a pure super-Yang-Mills (SYM) theory and R symmetry becomes spontaneously broken via gaugino condensation [67]. This external waterfall transition is associated with the breaking of a Z_2 parity, which is, however, only an approximate symmetry. For this reason, we do not have to fear the production of topological defects (i.e., domain walls) during the waterfall transition in the R-symmetry-breaking sector.

After these introductory remarks, we are now in the position to present our analysis. The remainder of this paper is organized as follows. In the next section, we will show how the IYIT DSB model may give rise to dynamical F-term hybrid inflation. Here, we will first argue why the original Polonyi model [68] of supersymmetry breaking is *not* sufficient for a successful realization of Polonyi inflation, even if we assume zero constant in the superpotential during inflation. We, therefore, conclude that we only have a chance of successfully realizing Polonyi inflation in the presence of radiative corrections—which leads us to consider the IYIT model as a possible UV completion of the original Polonyi model. We then derive the scalar potential of F-term hybrid inflation in the IYIT model and discuss its embedding into supergravity. As an interesting aside, we demonstrate that Polonyi inflation is incompatible with the concept of an approximate shift symmetry in the inflaton direction [69]. Instead, it turns out that Polonyi inflation requires a near-canonical Kähler potential. In Sec. III, we expand on our mechanism of late-time R symmetry breaking, showing how a small inflaton field value may trigger gaugino condensation in a separate hidden sector. Related to that, we comment on the backreaction of the R-symmetry-breaking sector on the inflationary dynamics and discuss how the two sectors of supersymmetry and R symmetry breaking have to conspire to yield a zero CC in the true vacuum after inflation. In Sec. IV, we turn to the phenomenological implications of our scenario. Here, we identify the viable region in parameter space that leads to agreement with the latest Planck data on the inflationary CMB observables [58]. As we are able to show, a scalar spectral index of $n_s \simeq 0.968$ can be easily achieved for an $\mathcal{O}(1)$ Yukawa coupling, $\lambda \simeq 2$, and a slightly suppressed coefficient in the noncanonical Kähler potential, $\epsilon \simeq 0.2$. The amplitude of the scalar power spectrum, A_s, fixes the dynamical scale of the SUSY-breaking sector to a value close to the unification scale, $\Lambda \simeq 1 \times 10^{16}$ GeV, suggesting that our setup may eventually be part of a grand unified theory (GUT). As a characteristic feature of Polonyi inflation, we highlight the fact that the relation between the gravitino mass and the inflationary Hubble rate in Eq. (3) directly translates into a one-to-one correspondence between the gravitino mass $m_{3/2}$ and the tensor-to-scalar ratio r,

$$m_{3/2} \simeq \frac{\pi}{\sqrt{2}}(rA_s)^{1/2}M_{\rm Pl} \sim 10^{12}\ \mathrm{GeV}\left(\frac{r}{10^{-4}}\right)^{1/2}. \tag{5}$$

At this point, it is interesting to note that the observed value of the scalar spectral amplitude, $A_s \sim 10^{-9}$, might be the result of anthropic selection [70]. Together with the paradigm of slow-roll inflation (which implies $r \ll 1$), the anthropic value of A_s could, therefore, explain why the soft SUSY mass scale is so much higher than the electroweak or TeV scale. Moreover, we point out in Sec. IV that, after inflation, the Polonyi field mostly decays into gravitinos. Polonyi inflation is, thus, followed by a phase of gravitino domination [71], which leads to reheating around temperatures of $\mathcal{O}(10^8)$ GeV. This paves the way for thermal wino dark matter as well as thermal leptogenesis [72] enhanced by resonance effects [73].

Appendix A contains some technical details regarding the derivation of the effective inflaton potential. In particular, we show how the one-loop corrections to the effective potential may be obtained, to leading order, from an effective Kähler potential. In Appendix B, we explain in more detail why Polonyi inflation in the IYIT model does not work, if the superpotential already contains a constant term from the very beginning. This completes our argument that, in the context of our minimal model, successful Polonyi inflation requires (i) radiative corrections, (ii) a near-canonical Kähler potential, and (iii) late-time R symmetry breaking. Among all possible choices regarding (i) the type of interactions that the Polonyi field participates in, (ii) the shape of the Kähler potential, and (iii) the chronology of supersymmetry and R symmetry breaking, this leave one unique possibility for how to realize Polonyi inflation.

II. DYNAMICAL INFLATION IN THE IYIT SUPERSYMMETRY BREAKING MODEL

A. Inflation in the original Polonyi model and the need for radiative corrections

The Polonyi model [68] is the simplest O'Raifeartaigh model of supersymmetry breaking via a nonvanishing F-term. Its superpotential consists of a tadpole term and a constant,

$$W = \mu^2 \Phi + w. \tag{6}$$

Here, Φ denotes the chiral Polonyi field, μ is the scale of supersymmetry breaking and w is a constant that breaks R symmetry and which determines the potential energy density in the ground state. For a canonical Kähler potential and fine-tuning w, so that it takes the particular value $w_0 = (2-\sqrt{3})\mu^2 M_{\rm Pl}$, this model has a Minkowski vacuum at $\langle\Phi\rangle = (\sqrt{3}-1)M_{\rm Pl}$, in which supersymmetry is broken by the Polonyi F-term, $|F_\Phi| = \mu^2$. It has been known for a long time that the scalar potential for the Polonyi field around this vacuum is unfortunately too steep to support slow-roll inflation [30]. In Appendix B, we review and extend this argument, showing for various choices of the Kähler potential that, with $w = w_0$ from the very beginning, the Polonyi model does not give rise to inflation. We consider, in particular, a canonical Kähler potential supplemented by higher-dimensional corrections as well as Kähler potentials featuring an approximate shift symmetry either along the real or the imaginary axis in the complex plane. In none of the cases under consideration inflation is viable—either because we fail to satisfy the slow-roll conditions or because the scalar potential does not exhibit a global Minkowski vacuum in the first place.

This immediately raises the question whether inflation might perhaps become possible in the Polonyi model, if we impose a discrete R symmetry at high energies, so that $w = 0$ initially. Let us address this question for a canonical Kähler potential supplemented by a higher-dimensional correction,[5]

$$K = \Phi^\dagger\Phi + \frac{\epsilon}{(2!)^2}\left(\frac{\Phi^\dagger\Phi}{M_{\rm Pl}}\right)^2 + \mathcal{O}(\epsilon^2, M_{\rm Pl}^{-4}), \quad \epsilon \lesssim 1. \tag{7}$$

Here, we assume that the Kähler potential is always dominated by the canonical term, $K \supset \Phi^\dagger\Phi$, also at field values above the Planck scale. An exhaustive study of arbitrary choices for the Kähler potential is beyond the scope of this paper. Under this assumption, the scalar potential always picks up a SUGRA correction, $V \propto e^{K/M_{\rm Pl}^2}$, which spoils the flatness of the potential at super-Planckian field values. For this reason, we only have a chance of realizing slow-roll inflation at field values below the Planck scale. For the Kähler potential in Eq. (7), the scalar potential in supergravity then takes the following form:

$$V(\varphi) = V_0\left[1 - \frac{\epsilon}{2}\left(\frac{\varphi}{M_{\rm Pl}}\right)^2 + \frac{1}{8}\left(1 - \frac{7\epsilon}{2} + \frac{8\epsilon^2}{3}\right)\left(\frac{\varphi}{M_{\rm Pl}}\right)^4 + \mathcal{O}(\varphi^6)\right],$$
$$V_0 = \mu^4, \tag{8}$$

where the real scalar field φ denotes the canonically normalized radial component of the complex Polonyi scalar $\tilde{\phi} = \tilde{\varphi}/\sqrt{2}e^{i\tilde{\theta}}$ contained in Φ (see Sec. II D) and where we have introduced V_0 as the SUSY-breaking vacuum energy density at $\varphi = 0$. From the form of the scalar potential in Eq. (8), it is evident that, even with w being set to zero, the Polonyi model fails to yield successful inflation. For instance, if we choose ϵ to be negative, the field φ is driven towards the origin by a positive mass squared, similarly as in chaotic inflation [34]. Inflation may then take place at small field values close to the origin—but not in accord with the observational data. To see this, consider the slow-roll parameters ε and η,

[5]In Sec. II D, we will discuss the same question for an approximately shift-symmetric Kähler potential. In this case, inflation turns out be unfeasible because, with the superpotential being given as $W = \mu^2\Phi$, the $-3\exp[K/M_{\rm Pl}^2]|W|^2/M_{\rm Pl}^2$ SUGRA term in the scalar potential induces a tachyonic mass for the Polonyi field. This results in a global AdS minimum.

$$\varepsilon = \frac{M_{\rm Pl}^2}{2}\left(\frac{V'}{V}\right)^2, \qquad \eta = M_{\rm Pl}^2 \frac{V''}{V},$$
$$V' = \frac{dV}{d\varphi}, \qquad V'' = \frac{d^2V}{d\varphi^2}. \tag{9}$$

Independent of the sign of ϵ, we have $\eta \geq -\epsilon$. For negative ϵ, the slow-roll parameter η is, therefore, bound to be positive, while the slow-roll parameter ε turns out to be negligibly small during inflation, $\varepsilon \ll |\eta|$. According to the slow-roll formula for the scalar spectral index, $n_s = 1 + 2\eta_* - 6\varepsilon_*$, we will then always obtain a blue-tilted scalar spectrum ($n_s > 1$). In view of the latest best-fit value for n_s reported by the Planck Collaboration, $n_s^{\rm obs} = 0.9677 \pm 0.0060$ [58], such a spectrum is clearly ruled out by the observational data. On the other hand, if we choose ϵ to be positive, the Polonyi field acquires a tachyonic mass around the origin and inflation proceeds from small to large field values, similarly as in new inflation [2]. It is then hard to imagine how the generation of the constant w in the superpotential should be triggered, after a sufficient amount of inflation, at field values close or even above the Planck scale. But more than that, even if we assume that this problem could somehow be solved, the scalar potential in Eq. (8) still does not lead to an acceptable phenomenology. In new inflation, the scalar spectral index turns out to be bounded from above, $n_s \lesssim 0.95$ [74,75], which deviates from the observed value by at least 3σ. Therefore, also for $\epsilon > 0$, we fail to reach consistency with the observational data.

In summary, we conclude that the bare Polonyi model based on the superpotential in Eq. (6)—and for reasonable *Ansätze* regarding the shape of the Kähler potential—does not allow for a successful realization of slow-roll inflation. In respect of this null result, two comments are in order. (i) We emphasize that our analysis of the Polonyi Kähler potential in this paper does *not* mount up to a general no-go theorem. In the most general case, the Kähler potential for the Polonyi field is given by an arbitrary function F of the field Φ and its conjugate, $K = M_{\rm Pl}^2 F(\Phi/M_{\rm Pl}, \Phi^\dagger/M_{\rm Pl})$. In absence of any other scale, it is, in particular, clear that $M_{\rm Pl}$ can be the only relevant scale in the Kähler potential. It may then well be that certain fine-tuned functions F do allow for successful inflation in the Polonyi model, after all (see, e.g., [76] for a discussion of fine-tuned Kähler potentials in the context of SUGRA models of inflation). In the following, we will, however, ignore the possibility of such a biasedly chosen Kähler potential and focus on the usual suspects: Kähler potentials that are either near-canonical or approximately shift-symmetric. (ii) The fact that we are unable to realize successful inflation in the bare Polonyi model is not a serious problem, as the Polonyi model is not expected to be a fundamental description of spontaneous supersymmetry breaking, anyway. It should rather be seen as the effective theory resulting from some UV dynamics that provide a dynamical explanation for the origin of the parameters μ and w in Eq. (6). From this perspective, it is then more likely than not that the Polonyi field is not only subject to its gravitational self-interaction, but that it also participates in Yukawa interactions with heavy matter fields in the UV theory. In the corresponding effective Polonyi model at low energies, these matter fields are integrated out, so that they no longer appear in the superpotential. But their couplings to the Polonyi field still yield radiative corrections to the scalar potential in Eq. (8), which affect the inflationary dynamics. In such a modified setup, i.e., in the original Polonyi model supplemented by radiative corrections, successful Polonyi inflation may, therefore, very well be an option. In the following, we will construct a minimal extension of the Polonyi model where this is indeed the case. We shall consider a minimal UV completion of the Polonyi model—consisting of two strongly coupled $SU(2)$ sectors that account for the dynamical origin of the parameters μ and w, respectively—and demonstrate that, in the presence of radiative corrections, successful Polonyi inflation is indeed feasible for a natural choice of parameter values.

B. Dynamical supersymmetry breaking in the low-energy regime of the IYIT model

One of the simplest ways to generate the SUSY breaking scale μ in Eq. (6) is to identify the Polonyi field as part of the IYIT model—the simplest vectorlike model of dynamical supersymmetry breaking [59]. In its most general formulation, the IYIT model is based on a strongly coupled $Sp(N)$ gauge theory featuring $N_f = N + 1$ pairs of "quark fields" Ψ^i that transform in the fundamental representation of $Sp(N)$. The gauge dynamics of this model are associated with a dynamical scale Λ, which denotes the energy scale at which the $Sp(N)$ gauge coupling formally diverges. The low-energy effective theory below the dynamical scale exhibits a quantum moduli space of degenerate supersymmetric vacua, which is spanned by $N_f(2N_f - 1)$ gauge-invariant composite flat directions (or "meson fields") M^{ij},

$$M^{ij} \simeq \frac{1}{\eta\Lambda}\langle \Psi^i \Psi^j \rangle, \quad i, j = 1, 2, \ldots, 2N_f. \tag{10}$$

Here, the parameter η is a dimensionless numerical factor, which ensures the canonical normalization of the meson fields M^{ij} at low energies. Naive dimensionful analysis (NDA) [77] leads us to expect that η should be of $\mathcal{O}(4\pi)$ and, for definiteness, we will, therefore, simply set $\eta = 4\pi$ in the following. The quantum moduli space of the IYIT model is subject to the following constraint pertaining to the meson VEVs:

$$\mathrm{Pf}(M^{ij}) \simeq \left(\frac{\Lambda}{\eta}\right)^2, \tag{11}$$

which represents the quantum mechanically deformed version of the classical constraint $\mathrm{Pf}(M^{ij}) = 0$ [60]. In order to break supersymmetry in the IYIT model, one has to lift the flat directions M^{ij}, so that Eq. (11) is longer compatible with a vanishing vacuum energy density. This is readily done by coupling the quark pairs $\Psi^i\Psi^j$ to a corresponding number of $Sp(N)$ singlet fields, Z_{kl}, in the tree-level superpotential,

$$W_{\mathrm{IYIT}}^{\mathrm{tree}} = \frac{1}{4}\lambda_{ij}^{kl} Z_{kl}\Psi^i\Psi^j, \qquad Z_{kl} = -Z_{lk},$$
$$\lambda_{ij}^{kl} = -\lambda_{ji}^{kl} = \lambda_{ji}^{lk}, \quad i,j,k,l = 1,2,\ldots,2N_f, \tag{12}$$

where λ_{ij}^{kl} denotes a matrix of Yukawa couplings with at most $N_f(2N_f-1)$ independent eigenvalues. These Yukawa couplings induce Dirac mass terms for the meson and singlet fields at low energies,

$$W_{\mathrm{IYIT}}^{\mathrm{eff}} \simeq \frac{1}{4}\lambda_{ij}^{kl}\frac{\Lambda}{\eta}Z_{kl}M^{ij}, \tag{13}$$

so that the singlet F-term conditions, $\lambda_{ij}^{kl}M^{ij} = 0$, are incompatible with the deformed moduli constraint, $\mathrm{Pf}(M^{ij}) \neq 0$. Supersymmetry is then broken *à la* O'Raifeartaigh via nonvanishing F-terms.

In global supersymmetry and for all Yukawa couplings in Eq. (12) being equal, $\lambda_{ij}^{kl} = \lambda\delta_i^k\delta_j^l$, the anomaly-free global flavor symmetry G_F of the IYIT model is given as follows:

$$G_F = SU(2N_f) \times Z_{2N_f} \times U(1)_R. \tag{14}$$

Here, the discrete Z_{2N_f} symmetry is the anomaly-free subgroup of the anomalous $U(1)$ that is contained in the full flavor symmetry at the classical level, $U(2N_f)\times U(1)_R \cong SU(2N_f) \times U(1) \times U(1)_R$. Under the Z_{2N_f} symmetry, all quarks carry charge 1, while the singlet fields carry charge -2. Meanwhile, the presence of the global continuous R symmetry is characteristic for a large class of DSB models [42]. Under $U(1)_R$, the quark and singlet fields carry charges 0 and 2, respectively. In the SUSY-breaking vacuum of the IYIT model, where $\langle M\rangle \neq 0$ and $\langle Z\rangle = 0$, R symmetry, therefore, remains unbroken. As a global symmetry, the continuous R symmetry of the IYIT model is, of course, only an approximate symmetry, which we expect to be broken by quantum gravitational effects [36]. On the other hand, recall that we assume the discrete subgroup $Z_4^R \subset U(1)_R$ to be gauged (see our discussion in Sec. I C). This protects the quality of the $U(1)_R$ symmetry, and it is reasonable to assume that all gravity-induced $U(1)_R$-breaking effects in the IYIT sector are suppressed. In the following, we will, therefore, stick to the effective superpotential in Eq. (13) and neglect the possibility of small R symmetry-breaking corrections.

From now on, let us restrict ourselves to the simplest version of the IYIT model: a strongly coupled $Sp(1) \cong SU(2)$ gauge theory featuring four matter fields Ψ^i and six singlet fields Z_{kl}. The non-Abelian flavor symmetry $SU(2N_f)$ then corresponds to a global $SU(4)$, under which M^{ij} and Z_{kl} transform as six-dimensional antisymmetric rank-2 tensor representations. Here, note that $SU(4)$ is the double cover of $SO(6)$. This allows us to rewrite the meson and singlet fields as vector representations of $SO(6)$,

$$\begin{pmatrix} X^0 \\ X^1 \\ X^2 \\ X^3 \\ X^4 \\ X^5 \end{pmatrix} = \frac{1}{\sqrt{2}}\begin{pmatrix} +1(M^{12}+M^{34}) \\ -1(M^{13}-M^{24}) \\ +1(M^{14}+M^{23}) \\ -i(M^{14}-M^{23}) \\ -i(M^{13}+M^{24}) \\ +i(M^{12}-M^{34}) \end{pmatrix},$$
$$\begin{pmatrix} S_0 \\ S_1 \\ S_2 \\ S_3 \\ S_4 \\ S_5 \end{pmatrix} = \frac{1}{\sqrt{2}}\begin{pmatrix} +1(Z_{12}+Z_{34}) \\ -1(Z_{13}-Z_{24}) \\ +1(Z_{14}+Z_{23}) \\ +i(Z_{14}-Z_{23}) \\ +i(Z_{13}+Z_{24}) \\ -i(Z_{12}-Z_{34}) \end{pmatrix}. \tag{15}$$

As we will see shortly, it will turn out to be convenient to study the SUSY-breaking dynamics of the IYIT sector in terms of this "$SO(6)$ language" rather than in terms of the original "$SU(4)$ language".

Trading the mesons M^{ij} for the new fields X^a, the Pfaffian constraint in Eq. (11) can be written as

$$\mathrm{Pf}(M^{ij}) \equiv \frac{1}{2}(X\cdot X) \equiv \frac{1}{2}\sum_{a=0}^{5} X^aX^a$$
$$\equiv \frac{1}{2}(X^a)^2 \simeq \left(\frac{\Lambda}{\eta}\right)^2, \tag{16}$$

which defines a sphere in the six-dimensional space spanned by the six meson coordinates X^a. An elegant way to enforce this constraint is to directly incorporate it into the effective superpotential [59,60],

$$W_{\mathrm{eff}}^{\mathrm{dyn}} \simeq \frac{\kappa}{\eta}\Lambda^2 T C(x^a), \qquad C(x^a) = \frac{1}{2}(x^a)^2 - 1,$$
$$x^a = \frac{X^a}{\Lambda/\eta}. \tag{17}$$

Here, the field T denotes a Lagrange multiplier, the corresponding F-term condition of which, $C(x^a) = 0$, is nothing but a reformulation of the moduli constraint in

Eq. (16). The physical status of the field T is unfortunately rather unclear and depends on whether (uncalculable) strong-coupling effects induce a kinetic term for T or not. If the field T should be physical, it might represent a dynamical glueball field, $T \sim \langle gg \rangle$, and the dimensionless coupling constant κ in Eq. (17) would be expected to take some value of $\mathcal{O}(1)$. In that case, the Pfaffian constraint in Eq. (16) would be satisfied only approximately, depending on the competition between the different F-term conditions that enter into the determination of the true ground state. If, on the other hand, T should be unphysical, we would have to treat it as a mere auxiliary field. In that case, the Pfaffian constraint should be satisfied exactly, which would require us to eventually take the limit $\kappa \to \infty$ in our analysis. In the following, we will suppose that the Lagrange multiplier field T is, indeed, a physical (glueball) field and set $\kappa = 1$ for definiteness.[6]

In terms of the fields X^a and S_a, the effective tree-level superpotential takes the following form:

$$W_{\rm eff}^{\rm tree} \simeq \frac{\lambda_a}{\eta} \Lambda S_a X^a, \tag{18}$$

where we assume, without loss of generality, the Yukawa couplings λ_a to be ordered by size, $\lambda_b \le \lambda_{b+1}$ for all $b = 0, \ldots, 4$. Note that we will refer to the smallest Yukawa coupling, λ_0, simply as λ in the following, $\lambda \equiv \lambda_0$. For generic values of the six Yukawa couplings λ_a, the non-Abelian flavor symmetry is completely broken,

$$\lambda_a \text{ all different} \Rightarrow G_F = SO(6) \times Z_4 \times U(1)_R \to Z_4 \times U(1)_R. \tag{19}$$

The total effective superpotential is given by the sum of $W_{\rm eff}^{\rm dyn}$ in Eq. (17) and $W_{\rm eff}^{\rm tree}$ in Eq. (18),

$$W_{\rm eff} \simeq \frac{\lambda_a}{\eta} \Lambda S_a X^a + \frac{\kappa}{\eta} \Lambda^2 T C(x^a). \tag{20}$$

We mention once more that, as a consequence of the nonrenormalization theorem, the effective superpotential does not receive radiative corrections in perturbation theory [23]. Given this form of the total superpotential and assuming a quadratic Kähler potential for all meson and singlet fields, the global minimum of the resulting F-term scalar potential is located at

$$\langle X^0 \rangle = \pm \sqrt{2}(1-\zeta)^{1/2} \frac{\Lambda}{\eta},$$
$$\langle S_0 \rangle = \mp \sqrt{2} \left(\frac{1-\zeta}{\zeta}\right)^{1/2} \langle T \rangle,$$
$$\langle X^n \rangle = \langle S_n \rangle = 0, \quad n = 1, \ldots, 5, \tag{21}$$

where $\langle T \rangle$ is undetermined at tree level. The sign ambiguity is a consequence of the Z_4 flavor symmetry [see Eq. (19)]. Moreover, the parameter ζ measures how well the deformed moduli constraint is satisfied,

$$\zeta = \langle |C(x^a)| \rangle = \left(\frac{\lambda}{\kappa\eta}\right)^2 = \frac{\lambda^2}{16\pi^2}. \tag{22}$$

For perturbative values of the Yukawa coupling λ, i.e., for $\lambda \sim 1$ or smaller, we have $\zeta \ll 1$, which tells us that the deformed moduli constraint is fulfilled almost exactly. On the other hand, for nonperturbative values of λ, i.e., for λ as large as $\lambda \sim 4\pi$, the parameter ζ becomes of $\mathcal{O}(1)$, indicating that the constraint function $C(x^a)$ significantly deviates from zero. To put this result into perspective, we must remember that, for nonperturbative values of the Yukawa coupling, incalculable corrections to the effective Kähler potential due to strong-coupling effects become important. In fact, as pointed out by Chacko et al., these nonperturbative corrections are only negligible as long as $\lambda \ll 4\pi$ [79]. We can, therefore, trust our above analysis, based on a canonical Kähler potential, only as long as λ remains in the perturbative regime. For this reason, we will, from now on, only consider λ values at most as large as $\lambda_{\rm max} \simeq 4$, so that we may always maintain a hierarchy among λ and η (i.e., so that $\lambda/\eta \lesssim 10^{-0.5}$). For the parameter ζ, this then means that it can take values at most as large as $\zeta_{\rm max} \simeq 0.1$. This translates into the statement that the moduli constraint is always satisfied in our analysis—up to a deviation of at most 10%, $\langle |C(x^a)| \rangle \lesssim 0.1$.

In passing, we also mention that the scalar potential exhibits a saddle point at the origin in field space as well as two saddle points along each direction X^n in moduli space. Here, the loci of the saddle points away from the origin have the same functional form as $\langle X^0 \rangle$ and $\langle S^0 \rangle$ in Eq. (21), the only difference being that λ in Eq. (22) needs to be exchanged with the respective Yukawa coupling λ_n. The low-energy vacuum along the X^0 axis, therefore, not only marks the global (and only local) minimum of the scalar potential, it is also the stationary point at which the deformed moduli constraint is fulfilled best.

In the true vacuum, supersymmetry is broken by the nonvanishing F-terms of the fields S_0 and T,

$$\langle |F_{S_0}| \rangle = \sqrt{2}(1-\zeta)^{1/2} \lambda \left(\frac{\Lambda}{\eta}\right)^2,$$
$$\langle |F_T| \rangle = \zeta^{1/2} \lambda \left(\frac{\Lambda}{\eta}\right)^2. \tag{23}$$

[6]For an extended discussion of this point, see also [41,78].

The F-term of the meson field X^0, on the other hand, (which seems to be nonzero at first sight) cancels,

$$\langle |F_{X^0}|\rangle = \frac{\lambda}{\eta}\Lambda\langle S_0\rangle + \kappa\eta\langle T\rangle\langle X^0\rangle = \left(\frac{\lambda}{\eta} - \kappa\zeta^{1/2}\right)\Lambda\langle S_0\rangle = 0. \quad (24)$$

In order to identify the mass eigenstates around the true vacuum, we now shift the field X^0 by its VEV,

$$X^0 = \langle X^0\rangle + \Xi^0, \quad (25)$$

and rotate the SUSY-breaking fields S_0 and T by their mixing angle β,

$$\begin{pmatrix}\Phi\\ \Sigma\end{pmatrix} = \begin{pmatrix}\cos\beta & -\sin\beta\\ \sin\beta & \cos\beta\end{pmatrix}\begin{pmatrix}S_0\\ T\end{pmatrix}, \qquad \tan\beta = \frac{\langle|T|\rangle}{\langle|S_0|\rangle} = \frac{1}{\sqrt{2}}\left(\frac{\zeta}{1-\zeta}\right)^{1/2}. \quad (26)$$

After these field transformations, the effective superpotential for Φ, Σ, Ξ^0, S_n, and X^n reads as follows:

$$W_{\rm eff} \simeq \mu^2\Phi + m_0\Sigma\Xi^0 + m_n S_n X^n - (\kappa_\Phi\Phi - \kappa_\Sigma\Sigma)\left[\frac{1}{2}(\Xi^0)^2 + \frac{1}{2}(X^n)^2\right], \quad (27)$$

which is (apart from the missing constant term w) exactly of the form anticipated at the end of Sec. II A.

First of all, note that the first term on the right-hand side of Eq. (27), $\mu^2\Phi$, is nothing but the dynamical realization of the SUSY-breaking tadpole term in Eq. (6) within the IYIT model. The field Φ is, thus, to be identified as the chiral Polonyi field which breaks supersymmetry via its nonzero F-term,

$$\langle|F_\Phi|\rangle = \mu^2, \qquad \mu = (2-\zeta)^{1/4}\lambda^{1/2}\frac{\Lambda}{\eta}, \quad (28)$$

where the parameter μ denotes again the SUSY breaking scale. In this sense, Eqs. (27) and (28) show that the IYIT model serves, indeed, as a viable UV completion for at least half the Polonyi model: The SUSY-breaking dynamics of the IYIT model manage to provide a dynamical explanation for the SUSY breaking scale μ. However, as the IYIT model preserves R symmetry in its ground state, it is not capable of accounting for the origin of the R symmetry-breaking constant w in the Polonyi superpotential.

Second of all, the effective superpotential in Eq. (27) contains (also as envisaged at the end of Sec. II A) Yukawa couplings between the Polonyi field Φ and a number of massive matter fields, Ξ^0 and X^n. Here, the meson masses, m_0 and m_n, follow from Dirac mass terms together with the singlet fields Σ and S_n,

$$m_0 = \frac{m}{r_0}, \qquad m_n = \frac{m}{r_n}, \qquad m = \kappa_\Phi^{1/2}\mu = \lambda\frac{\Lambda}{\eta}, \qquad r_0 = \left(\frac{\zeta}{2-\zeta}\right)^{1/2} = \sin\beta, \qquad r_n = \frac{\lambda}{\lambda_n}, \quad (29)$$

where we have introduced the flavor-independent mass scale m as well as the respective ratios r_0 and r_n between this scale m and the masses m_0 and m_n. From the fact that $m = \kappa_\Phi^{1/2}\mu$ [see Eq. (31) below], it immediately follows that the scale m represents the amount of SUSY-breaking mass splitting within the respective meson and singlet multiplets that is induced by the tadpole term in Eq. (27), see Appendix A 1 for details. Given the definition of ζ in Eq. (22) and recalling that we assume λ to be the smallest among all Yukawa couplings, $\lambda \le \lambda_n$, we also find that the ratios r_0 and r_n are bounded from above,

$$r_0 = \frac{m}{m_0} \le 1, \qquad r_n = \frac{m}{m_n} \le 1, \quad (30)$$

so that the SUSY-breaking mass splitting m never exceeds the supersymmetric Dirac masses m_a. Moreover, for the parameter range of interest, $\lambda \lesssim 4$ and $\lambda_n \lesssim 4\pi$, the ratio r_0 always turns out to be the smallest, $r_0 < r_n$, which leads to the interesting (and to some extent counterintuitive) result that the zeroth flavor, i.e., the flavor with the smallest Yukawa coupling, ends up being stabilized the most. In addition to that, the Ξ^0 flavor is also singled out by the fact that its Dirac mass partner is none of the original singlet fields S_a, but the linear combination Σ that we introduced in Eq. (26) and which, for small values of ζ, mostly consists of the Lagrange multiplier field T. This also explains why, for large κ (and, hence, small ζ), the Ξ^0 mass m_0 diverges. In this limit, the Pfaffian constraint is fulfilled exactly, which results in the decoupling of T and removes one meson multiplet (i.e., Ξ^0) from the spectrum.

Next to the Polonyi field Φ, also the "stabilizer field" Σ couples to the meson fields Ξ^0 and X^n. Here, the strengths of the respective Yukawa couplings, κ_Φ and κ_Σ, are given by $\kappa\eta$ and the mixing angle β,

$$\kappa_\Phi = \sin\beta\,\kappa\eta = \frac{\lambda}{(2-\zeta)^{1/2}}, \qquad \kappa_\Sigma = \cos\beta\,\kappa\eta = \left(\frac{2}{\zeta}\right)^{1/2}\left(\frac{1-\zeta}{2-\zeta}\right)^{1/2}\lambda. \quad (31)$$

Just like the mass m_0, the Yukawa coupling κ_Σ diverges in the limit $\kappa \to \infty$. This is a trivial consequence of its proportionality to $\kappa\eta$. As for the Polonyi coupling κ_Φ, this divergence is, however, canceled out by the $\sin\beta$ factor. In contrast to κ_Σ, the coupling κ_Φ, therefore, always remains finite. In the limit $\kappa \to \infty$, it reproduces, in particular, the

Yukawa coupling of the singlet S_0 at energies above the dynamical scale,

$$W_{\rm IYIT}^{\rm tree} \supset \frac{\lambda}{\sqrt{2}} S_0 (\Psi^1\Psi^2 + \Psi^3\Psi^4). \tag{32}$$

Finally, we mention that the four new parameters μ, m_0, κ_Φ, and κ_Σ introduced in Eq. (27) are not linearly independent. In fact, they must be dependent, as they can all be expressed in terms of the three old parameters λ, Λ, and κ. By making use of Eqs. (28), (29), and (31), one easily convinces oneself that

$$\mu^2 = \frac{\kappa_\Phi}{\kappa_\Phi^2 + \kappa_\Sigma^2} m_0^2. \tag{33}$$

In order to see how the superpotential in Eq. (27) may give rise to Polonyi inflation, it is instructive to forget about the Yukawa couplings of the field Σ for a moment and to rewrite Eq. (27) as follows:

$$W_{\rm eff} \simeq \kappa_\Phi \Phi \left[v^2 - \frac{1}{2}(\Xi^0)^2 + \frac{1}{2}(X^n)^2 \right] + m_0 \Sigma \Xi^0 + m_n S_n X^n + \ldots, \tag{34}$$

where the ellipsis stands for the Yukawa couplings involving the stabilizer field Σ and where we have introduced the mass scale v,

$$v = \frac{m}{\kappa_\Phi} = (2-\zeta)^{1/2}\frac{\Lambda}{\eta}. \tag{35}$$

Remarkably enough, the first part of the superpotential in Eq. (34) has the same form as the superpotential of supersymmetric F-term hybrid inflation [61] based on $SO(6)$. In the context of this interpretation, the Polonyi field plays the role of the chiral inflaton singlet, while the meson fields Ξ^0 and X^n act as a multiplet of FHI waterfall fields that transform in the vector representation of $SO(6)$. The mass scale v is then to be identified as the energy scale of the waterfall transition at the end of inflation, while the mass splitting m should be understood as the tachyonic mass of the FHI waterfall fields at $\Phi = v$. We note that it is this picture that the authors of [51] arrive at. In their model, no other singlets except for the Polonyi field Φ are introduced. The meson fields Ξ^0 and X^n, therefore, lack their Dirac mass partners, so that the resulting effective superpotential is exactly identical to the one of F-term hybrid inflation. This allows the authors of [51] to realize F-term hybrid inflation at large inflaton field values, $|\Phi| > v$, where the flatness of the inflaton potential is lifted by logarithmic loop corrections.

In our case, the situation at large field values is quite similar, as we will discuss shortly, but at small field values, it is drastically different. The second part of the superpotential in Eq. (34) introduces explicit Dirac mass terms for the "would-be waterfall fields" Ξ^0 and X^n that are absent in standard F-term hybrid inflation (as well as in the model in [51]). Accounting for the presence of these Dirac masses in the low-energy effective theory, the tachyonic waterfall mass m is always compensated [see Eq. (30)], so that none of the meson fields ever becomes destabilized. Because of that, the total superpotential in Eq. (34) fails to give rise to a waterfall transition and, even in the low-energy vacuum, we retain the vacuum energy density resulting from the nonzero Polonyi F-term, $V_0 = \langle |F_\Phi| \rangle^2 = \mu^4$. As anticipated in Sec. I D, this is a characteristic feature of our construction, in which we intend to use one and the same vacuum energy density for driving inflation and breaking supersymmetry. Moreover, independent of the symmetry group under which the meson fields transform, the absence of the waterfall transition automatically implies that the end of inflation is not accompanied by the production of topological defects. This may be regarded as a significant phenomenological advantage of our scenario over standard F-term hybrid inflation.

C. Pseudomodulus potential in global supersymmetry

To the best of our knowledge, the superpotential in Eq. (34) has not been considered as the dynamical origin of inflation, so far. Here, part of the reason certainly is that successful inflation based on Eq. (34) is bound to require a rather high SUSY breaking scale μ. As explained in the Introduction, in supergravity, this necessitates a large constant in the superpotential to cancel the CC in the true vacuum, which then spoils slow-roll inflation (see Sec. I A). In the rest of this paper, we will, however, show that the superpotential in Eq. (34) can yield successful Polonyi inflation after all, if we generate the constant in the superpotential only towards the end of inflation. To this end, we shall now examine the effective one-loop potential for the complex Polonyi scalar $\phi = \varphi/\sqrt{2}e^{i\theta}$ in global supersymmetry more closely. In the next sections, we will then turn to the embedding of the IYIT model into supergravity (see Sec. II D) as well as to the generation of the constant term in a separate hidden sector (see Sec. III).

The Yukawa interactions between the Polonyi field Φ and the meson fields Ξ^0 and X^n in Eq. (27) lead to radiative corrections to the Polonyi potential, $V_{1-{\rm loop}}(\varphi)$, that may be calculated according to the Coleman-Weinberg (CW) formula for the effective one-loop potential [80]. The details of our calculation may be found in Appendix A; in the following, we will merely summarize our results. Generally speaking, the effective potential may be divided into two regimes: (i) At large field values, all of the meson fields acquire a large inflaton-dependent Majorana mass $M(\varphi) = \kappa_\Phi|\phi|$. For $M(\varphi) \gg m_a$, the supersymmetric Dirac masses m_a in Eq. (27) are, hence, negligible and the effective potential takes the usual logarithmic form as in standard F-term hybrid inflation, $V_{1-\mathrm{loop}}(\varphi) \propto m^4 \ln M(\varphi)$. (ii) On the other hand, at small Polonyi field values, such that $M(\varphi) \ll m_a$, the Dirac masses m_a become more relevant. Integrating out the "heavy fields" then leads to a quadratic Polonyi potential around the origin, $V_{1-\mathrm{loop}}(\varphi) \propto m^2 M^2(\varphi)$. We note that, as has been shown for the first time in [79], the effective potential around the origin has *positive* curvature. The low-energy vacuum at $\langle\Phi\rangle = 0$ is, therefore, indeed stable. Moreover, we find that, at large as well as at small field values, the effective potential scales with the soft SUSY-breaking mass scale m, i.e., the smallest mass scale in our model [see Eq. (30)]. This illustrates how supersymmetry succeeds in protecting the inflaton potential from picking up too-large radiative corrections [25].

To quantify the above statements, it is convenient to introduce the following mass ratios:

$$R_a(\varphi) = \frac{M(\varphi)}{m_a}, \quad M(\varphi) = \kappa_\Phi |\phi| = \frac{\lambda}{(2-\zeta)^{1/2}} \frac{\varphi}{\sqrt{2}}. \tag{36}$$

For large values of $R_a(\varphi)$, we then obtain a logarithmic one-loop potential, while for small values of $R_a(\varphi)$, the one-loop corrections take a quadratic form. This behavior can be captured by studying the effective potential as a function of a single order parameter $x(\varphi)$, the geometric mean of all ratios $R_a(\varphi)$,

$$x(\varphi) = \left(\prod_a R_a(\varphi)\right)^{1/N_X} = \frac{M(\varphi)}{\bar{m}},$$
$$\bar{m} = \left(\prod_a m_a\right)^{1/N_X}, \qquad N_X = 6. \tag{37}$$

Here, $N_X = 6$ counts the number of meson fields in the IYIT sector, while $\bar{m}$ stands for the geometric mean of all explicit mass parameters in Eq. (27). In this sense, $\bar{m}$ denotes the "supersymmetric mass scale" of the IYIT sector, i.e., a characteristic value for the nonperturbatively generated Dirac masses in the low-energy effective theory. In the following, we will set $\bar{m}$ to the dynamical scale Λ, for definiteness,

$$\bar{m} = \Lambda. \tag{38}$$

This mainly serves the purpose to account, in an effective way, for heavy composite states with masses around the dynamical scale Λ that we expect to be present, but which we can unfortunately not explicitly describe in terms of our perturbative language at low energies. Formally, we can always set $\bar{m}$ to Λ in our calculation by choosing an appropriate value for the effective heavy-flavor Yukawa coupling $\tilde{\lambda}$,

$$\bar{m} = m_0^{1/6} \tilde{m}^{5/6}, \qquad \tilde{m} = \tilde{\lambda}\frac{\Lambda}{\eta}, \quad \tilde{\lambda} = (\lambda_1\lambda_2\lambda_3\lambda_4\lambda_5)^{1/5}. \tag{39}$$

By fixing $\tilde{\lambda}$ at a nonperturbative value, we are, therefore, able to enforce our designated value for $\bar{m}$,

$$\tilde{\lambda} = \left(\frac{r_0}{\lambda}\right)^{1/5} \eta^{6/5} \simeq \eta. \tag{40}$$

The value of $x(\varphi)$ indicates whether the Dirac masses m_a are negligible or not and, thus, decides whether we are in the logarithmic or the quadratic part of the effective potential. The transition between both regimes takes place at field values close to what we shall refer to as the critical field value φ_c,

$$x(\varphi_c) = 1 \Leftrightarrow M(\varphi_c) = \bar{m} \Rightarrow \varphi_c = \sqrt{2}\frac{\bar{m}}{\kappa_\Phi}. \tag{41}$$

This implies that the order parameter $x(\varphi)$ can also be regarded as the ratio of the actual and the critical field value, $x(\varphi) = \varphi/\varphi_c$. Far away from the critical field value, i.e., at $x(\varphi) \ll 1$ and $x(\varphi) \gg 1$, we now find the following expressions for the effective potential (see Appendix A for details):

$$x(\varphi) \ll 1 \Rightarrow V_{1-\mathrm{loop}}^{\mathrm{LE}}(\varphi) = \frac{1}{2} m_{\mathrm{eff}}^2 \varphi^2 + \mathcal{O}(x^4),$$
$$x(\varphi) \gg 1 \Rightarrow V_{1-\mathrm{loop}}^{\mathrm{HE}}(\varphi) = \frac{N_X}{16\pi^2} m^4 \ln x(\varphi) + \mathcal{O}(x^{-4}), \tag{42}$$

with m_{eff} denoting the effective one-loop mass of the Polonyi field around the origin,[7]

$$m_{\mathrm{eff}}^2 = (2\ln 2 - 1) N_X^{\mathrm{eff}}(r_a) \frac{\kappa_\Phi^2}{16\pi^2} m^2,$$
$$N_X^{\mathrm{eff}}(r_a) = \sum_a \omega(r_a), \qquad \omega(r_a) \approx r_a^2. \tag{43}$$

Here, $N_X^{\mathrm{eff}}(r_a)$ counts the effective number of mesons that contribute to the effective Polonyi mass. The full functional form of $N_X^{\mathrm{eff}}(r_a)$ is a sum of complicated loop factors $\omega(r_a)$. To good approximation, these loop functions, however, happen to coincide with the mass ratios r_a squared, $\omega(r_a) \approx r_a^2$. We can write the result in Eq. (42)

[7]A similar expression has been derived for the first time in [79]. Our result differs from the one in [79] to the extent that we allow for nonzero ζ (and, hence, nonzero r_0), which means that we do not necessarily enforce the moduli constraint exactly. The calculation in [79], on the other hand, is based on the assumption that the moduli constraint *is* fulfilled exactly, so that $r_0 \equiv 0$. For a recent derivation and discussion of the effective Polonyi mass in $SU(4)$ language, see [41,78].

more compactly, if we make use of the following two potential energy scales:

$$\Lambda_{\rm LE}^4 = \sum_a \frac{m^4}{16\pi^2}(2\ln 2 - 1)\left(\frac{\bar{m}}{m_a}\right)^2 = \frac{1}{2}m_{\rm eff}^2\varphi_c^2,$$
$$\Lambda_{\rm HE}^4 = \sum_a \frac{m^4}{16\pi^2} = N_X\frac{m^4}{16\pi^2}. \quad (44)$$

The effective potential far away from the critical field value φ_c then takes the following form (see Fig. 1):

$$V_{1-\rm loop}(\varphi) \approx \begin{cases} \Lambda_{\rm LE}^4 x^2(\varphi) & ;\ x \ll 1 \\ \Lambda_{\rm HE}^4 \ln x(\varphi) & ;\ x \gg 1 \end{cases},$$
$$x(\varphi) = \frac{\varphi}{\varphi_c} = \frac{|\phi|}{\Lambda/\kappa_\Phi}. \quad (45)$$

The crucial question which we need to answer in the following is: Can we use either the low-energy or the high-energy part of this effective potential to realize successful Polonyi inflation? Let us first investigate whether inflation might occur in the quadratic part close to the origin. As we know from standard chaotic inflation [34], the effective inflaton mass then needs to take a value of $\mathcal{O}(10^{13})$ GeV to ensure the correct normalization of the scalar power spectrum. This requires the coupling λ to take a value at least as large $\lambda \simeq 0.2$, since otherwise the dynamical scale Λ would have to be super-Planckian,

$$m_{\rm eff} \simeq 10^{13}\ {\rm GeV}\left(\frac{\lambda}{0.2}\right)^3\left(\frac{\Lambda}{M_{\rm Pl}}\right) + \mathcal{O}(\lambda^5). \quad (46)$$

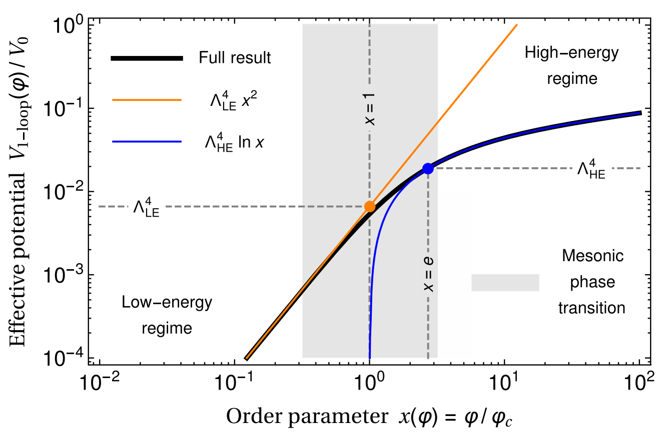

FIG. 1. Effective CW potential for the Polonyi field, $V_{1-\rm loop}$, as a function of the order parameter x for $\lambda = 1$, see Eq. (45). The potential energy scales $\Lambda_{\rm LE}$ and $\Lambda_{\rm HE}$ are given in Eq. (44). For values $x \sim 1$, we can only have limited trust in our perturbative result because of potentially important strong-coupling effects at energies close to the dynamical scale Λ.

At the same time, we know that chaotic inflation requires a large super-Planckian field excursion to yield a sufficient number of e-folds. The scalar perturbations probed in CMB observations, e.g., cross outside the Hubble horizon at a field value $\varphi_* \sim 15M_{\rm Pl}$. However, in the context of our SUSY-breaking model, this large field range does not "fit" into the low-energy part of the effective potential. This follows from the fact that, for $\lambda \simeq 0.2$ and $\Lambda \simeq M_{\rm Pl}$, the critical field value φ_c only becomes as large as $\varphi_c \simeq 10M_{\rm Pl}$,

$$\varphi_c \simeq 10M_{\rm Pl}\left(\frac{0.2}{\lambda}\right)\left(\frac{\Lambda}{M_{\rm Pl}}\right) + \mathcal{O}(\lambda). \quad (47)$$

Therefore, to raise φ_c, so as to make the field range required for chaotic inflation fit into the quadratic part of the effective potential, $\varphi_c \gtrsim \varphi_*$, we would have to go to smaller values of λ. But then, we are either forced to push Λ beyond the Planck scale or we fail to reproduce the correct scalar spectral amplitude. This eliminates the possibility of Polonyi inflation in the low-energy part of the effective potential, which is why we will focus on inflation in the logarithmic part of the effective potential from now on.

Before continuing, we, however, point out that inflation close to the origin might become possible, after all, if we relax our assumptions. That is, if we allowed for values of the dynamical scale as large as, say, $\Lambda \sim (8\pi)^{1/2}M_{\rm Pl}$, we would, in fact, be able to raise φ_c above φ_*. If we then trusted the full effective potential also at field values close to φ_c (see Fig. 1), inflation in the transitioning regime between the quadratic and the logarithmic part of the effective potential might become feasible. Such a scenario would promise to interpolate between the predictions of chaotic inflation and F-term hybrid inflation, so that we would expect it to result in interesting predictions for the tensor-to-scalar ratio, $r \sim 0.1$. Because of the uncertainties involved in such a scenario, we, however, do not pursue this idea any further in this paper and leave a more detailed study for future work. In closing, we remark that a similar model of subcritical hybrid inflation, based on a dynamically generated D-term [78], may be found in [81]. This model illustrates how to realize chaotic inflation after the waterfall transition of D-term hybrid inflation.

Let us now turn to the possibility of inflation in the logarithmic part of the effective potential. As we will show, in this part of the potential, successful Polonyi inflation is indeed feasible. To be on the safe side, we will limit our analysis in the following to field values that are larger than the critical field value by at least half an order of magnitude, $\varphi \gtrsim 10^{0.5}\varphi_c$. We do so because the effective Polonyi potential may receive nonperturbative corrections around $\varphi \sim \varphi_c$ that we do not have under control. In fact, around the critical field value, the inflaton-dependent mass $M(\varphi)$ drops below the dynamical scale Λ [see Eq. (41)]. This triggers the IYIT sector to transition from the high-energy quark-gluon regime into the low-energy meson regime. During this mesonic phase transition, the IYIT quarks become confined in the composite mesons and the strongly coupled $SU(2)$ gauge group becomes completely broken by the nonzero squark VEVs $\langle\Psi^1\Psi^2\rangle$ and $\langle\Psi^3\Psi^4\rangle$ that contribute to the meson VEV $\langle X^0\rangle$ [see Eqs. (15) and (21)],[8]

$$\langle\Psi^1\Psi^2\rangle = \langle\Psi^3\Psi^4\rangle \neq 0 \Rightarrow SU(2) \to \mathbb{1}. \tag{48}$$

Thanks to the fact that the $SU(2)$ symmetry is spontaneously broken down to "nothing," no topological defects are formed during the confining phase transition [82]. We emphasize that this is an important phenomenological feature of the IYIT model based on $SU(2)$, as it allows for the particular breaking pattern $SU(2) \to \mathbb{1}$. In summary, our scenario of Polonyi inflation, therefore, crucially differs from ordinary F-term hybrid inflation in the following respect: While ordinary F-term hybrid inflation ends in a waterfall transition—in the course of which the inflationary vacuum energy density is "eaten up" by the FHI waterfall fields and which potentially leads to the production of troublesome topological defects—our scenario of Polonyi inflation undergoes a confining quark-meson phase transition that *conserves* the inflationary vacuum energy density and that does *not* lead to the production of topological defects.

D. Embedding into supergravity and choice of the Kähler potential

In supergravity, the flatness of the tree-level Polonyi potential in global supersymmetry, $V(\varphi) = V_0 = \mu^4$, is not only lifted by the radiative corrections in Eq. (45), but also by gravitational corrections (see our discussion at the end of Sec. I C). In our case, these SUGRA corrections turn out to be rather mild for basically two reasons: (i) Since the superpotential in Eq. (27) only contains terms linear in Φ, the tree-level SUGRA mass of the Polonyi field accidentally cancels, as long as we assume a canonical Kähler potential. As far as the embedding into supergravity is concerned, this represents an important advantage of F-term hybrid inflation (and of our model) over alternative models that do feature higher powers of the inflaton field in the superpotential. (ii) Since we intend to realize Polonyi inflation in the logarithmic part of the effective potential, we will consistently work with sub-Planckian field values. Our scenario of Polonyi inflation will, hence, turn out to be a *small-field* model of inflation. Accordingly, the SUGRA corrections in our model are bound to be less significant than in alternative *large-field* models of inflation.

[8]Here, $\langle\Psi^1\Psi^2\rangle = \langle\Psi^3\Psi^4\rangle$ follows from the fact that the linear combination X^5 vanishes in the vacuum, $\langle X^5\rangle = 0$.

The total scalar potential for the complex Polonyi field $\tilde{\phi} \subset \Phi$ in supergravity now reads,[9]

$$V(\tilde{\phi}) = V_F(\tilde{\phi}) + V_{1\text{-loop}}(|\tilde{\phi}|), \tag{49}$$

with $V_{1\text{-loop}}$ being given in Eq. (45) and where V_F denotes the tree-level F-term potential in supergravity,

$$V_F = |F|^2 - 3\exp\left[\frac{K}{M_{\rm Pl}^2}\right]\frac{|W|^2}{M_{\rm Pl}^2}. \tag{50}$$

Here, $|F|$ denotes the norm of the generalized F-term vector $\mathcal{F}^i$ induced by the Kähler metric $K_i^{\bar{j}}$,

$$|F| = (\mathcal{F}\cdot\mathcal{F}^*)^{1/2}, \qquad \mathcal{F}\cdot\mathcal{F}^* = \mathcal{F}^i K_i^{\bar{j}}\mathcal{F}_{\bar{j}}^*,$$
$$K_i^{\bar{j}} = \frac{\partial^2 K}{\partial\phi^i\partial\phi_{\bar{j}}^*}. \tag{51}$$

The individual components of the F-term vector $\mathcal{F}^i$ are proportional to the conjugate of Kähler-covariant derivatives of the superpotential, $(D_jW)^*$, multiplied by the inverse of the Kähler metric, $K_{\bar{j}}^i$,

$$\mathcal{F}^i = -K^i{}_{\bar{j}}\exp\left[\frac{K}{2M_{\rm Pl}^2}\right](D_jW)^*, \qquad K^i{}_{\bar{j}} = (K^{-1})^i{}_{\bar{j}},$$
$$D_iW = \frac{\partial W}{\partial\phi^i} + \frac{W}{M_{\rm Pl}^2}\frac{\partial K}{\partial\phi^i}. \tag{52}$$

After integrating out the heavy fields Ξ^0, X^n, Σ, and S_n in Eq. (27), the effective superpotential of the IYIT sector reduces to the SUSY-breaking tadpole term for the Polonyi field [see Eq. (6)],

$$W_{\rm eff} \simeq \mu^2\Phi. \tag{53}$$

Meanwhile, the tree-level Kähler potential K is not unambiguously defined. At first sight, there are several well-motivated choices for K that come into question. For instance, we might think that an approximately shift-symmetric Kähler potential could help in keeping the SUGRA corrections small,[10]

$$K = \pm\frac{1}{2}(\Phi \pm \Phi^\dagger)^2 \mp \frac{\epsilon}{2}(\Phi \mp \Phi^\dagger)^2 + \mathcal{O}(\epsilon^2, M_{\rm Pl}^{-2}),$$
$$\epsilon \ll 1. \tag{54}$$

Depending on the sign choice, the real-valued (not canonically normalized) inflaton field $\tilde{\varphi}$ is then identified either as the real ($-$) or the imaginary ($+$) part of the complex scalar

[9]The field $\tilde{\phi}$ is not necessarily canonically normalized, which we indicate by placing a tilde on top of the symbol ϕ.

[10]A shift symmetry in the direction of the Polonyi field can never be an exact symmetry, as it is always explicitly broken by the superpotential, $W_{\rm eff} \simeq \mu^2\Phi$, and the radiative corrections in the scalar potential [see Eq. (45)]. For this reason, we expect that, also in the tree-level Kähler potential, the shift symmetry should be explicitly broken.

$\tilde{\phi} \subset \Phi$. Likewise, the real-valued (not canonically normalized) scalar inflaton partner $\tilde{\sigma}$ (i.e., the "sinflaton") corresponds to the imaginary ($-$) or the real ($+$) part of the complex scalar $\tilde{\phi} \subset \Phi$. As long as we do not assume an exact shift symmetry in the Kähler potential (i.e., $\epsilon \equiv 0$), $\tilde{\varphi}$ and $\tilde{\sigma}$ need to be canonically normalized,

$$\varphi = (1+\epsilon)^{1/2}\tilde{\varphi}, \qquad \sigma = (1+\epsilon)^{1/2}\tilde{\sigma}. \tag{55}$$

Unlike in the case of a canonical Kähler potential, the SUGRA F-term potential V_F now contains nonzero tree-level masses for the inflaton and its scalar partner. At the origin in field space, we find

$$m_\varphi^2 = -\frac{1-\epsilon}{(1+\epsilon)^2}\frac{3\mu^4}{M_{\rm Pl}^2}, \qquad m_\sigma^2 = +\frac{1-\epsilon}{(1+\epsilon)^2}\frac{3\mu^4}{M_{\rm Pl}^2}. \tag{56}$$

For $\epsilon \ll 1$, the inflaton direction is, hence, tachyonic. If we forget about the radiative corrections $V_{1-\rm loop}$ for a moment, this immediately implies that Polonyi inflation based on an approximately shift-symmetric Kähler potential is impossible—which completes our argument in Sec. II A (see footnote 5). On the other hand, taking the radiative corrections into account, we may hope that the effective inflaton mass $m_{\rm eff}$ around the origin could possibly compensate the tachyonic tree-level mass m_φ and, thus, stabilize the vacuum at $\varphi = 0$. It is clear that this will not be easy, since m_φ in Eq. (56) is typically very large,

$$|m_\varphi| \simeq 8\times 10^{15}\ \text{GeV}\left(\frac{\lambda}{0.2}\right)\left(\frac{\Lambda}{M_{\rm Pl}}\right)^2 + \mathcal{O}(\epsilon,\lambda^3). \tag{57}$$

In order to sufficiently suppress it, we need to go to smaller values of the dynamical scale Λ. Requiring that the total inflaton mass at the origin be nontachyonic, $m_\varphi^2 + m_{\rm eff}^2 > 0$, we find [see Eq. (46)],

$$\Lambda \lesssim 3\times 10^{15}\ \text{GeV}\left(\frac{\lambda}{0.2}\right)^2. \tag{58}$$

For such small values of the dynamical scale, the critical field value φ_c is clearly sub-Planckian [see Eq. (41)]. As discussed in the previous section, this eliminates the possibility of realizing inflaton in the quadratic part of the potential close to the origin. Our only remaining hope, therefore, is to realize inflation at field values $\varphi \gtrsim \varphi_c$. To this end, it is important that the total scalar potential does not exhibit a local minimum above the critical field value—otherwise the inflaton will never reach the origin and inflation would never end. But as it turns out, exactly such a local minimum is always present.

The point is that the quadratic part of the effective potential can only compensate the tachyonic tree-level mass up to field values of $\mathcal{O}(\varphi_c)$. At the same time, the gradient of the tree-level potential only turns positive at very large field values, $\varphi \geq \varphi_{\rm min} \propto M_{\rm Pl}/\epsilon$, where $\varphi_{\rm min}$ denotes the position of a global AdS vacuum. As long as $\varphi_c \lesssim \varphi_{\rm min}$, the gradient of the total inflaton potential is, therefore, guaranteed to turn negative in between φ_c and $\varphi_{\rm min}$—which indicates the presence of a local minimum. To evaluate the relation between φ_c and $\varphi_{\rm min}$ more precisely, consider the tree-level inflaton potential for $\sigma = 0$,

$$V_F(\varphi) = \exp\left[\frac{\epsilon}{1+\epsilon}\frac{\varphi^2}{M_{\rm Pl}^2}\right]\left(v_0 + \frac{1}{2}M_\varphi^2\varphi^2 + \frac{\lambda_\varphi}{4!}\varphi^4\right), \tag{59}$$

with the three parameters v_0, M_φ, and λ_φ being given as follows:

$$v_0 = \frac{\mu^4}{1+\epsilon}, \qquad M_\varphi^2 = -\frac{3-\epsilon}{1+\epsilon}\frac{v_0}{M_{\rm Pl}^2},$$
$$\lambda_\varphi = 24\left(\frac{\epsilon}{1+\epsilon}\right)^2\frac{v_0}{M_{\rm Pl}^4}. \tag{60}$$

The local minimum in the SUGRA tree-level potential, $\varphi_{\rm min}$, is then located at

$$\varphi_{\rm min} = \left(\frac{1+\epsilon}{\epsilon}\right)^{1/2}[(1+\text{sgn}(\epsilon))a - 1]^{1/2}M_{\rm Pl},$$
$$a = -\frac{\epsilon}{1+\epsilon}\frac{6}{\lambda_\varphi}\frac{M_\varphi^2}{M_{\rm Pl}^2}, \tag{61}$$

so that the AdS vacuum energy density takes the following value:

$$V_{\rm min} = -\exp\left[\frac{\epsilon}{1+\epsilon}\frac{\varphi_{\rm min}^2}{M_{\rm Pl}^2}\right]\left(\frac{3-\epsilon}{2|\epsilon|} - 2\right)v_0. \tag{62}$$

Given the above expression for $\varphi_{\rm min}$, we find that, in the entire parameter space of interest ($\lambda/\eta \lesssim 10^{-0.5}$ and $|\epsilon| \lesssim 10^{-0.5}$), φ_c and $\varphi_{\rm min}$ are always separated by at least one order of magnitude, $\varphi_c/\varphi_{\rm min} \lesssim 0.1$. In the shift-symmetric limit, $\epsilon \to 0$, the AdS vacuum disappears, in particular, altogether, $\varphi_{\rm min} \to \infty$, so that the total inflaton potential becomes unbounded from below. In summary, this shows that the effective potential manages to compensate the tachyonic tree-level mass only for too small a field range. It only "bends around" the tachyonic part of the tree-level potential up to $\varphi \sim \varphi_c$, while it should actually do this up to field values that are at least ten times as large, $\varphi \sim \varphi_{\rm min} \gtrsim 10\varphi_c$. This finally completes our argument that Polonyi inflation is incompatible with an approximate shift symmetry in the Kähler potential, independent of whether we account for the presence of radiative corrections or not.[11]

[11]Note that this conclusion does not hold for ordinary F-term hybrid inflation. As has been shown in [83], F-term hybrid inflation can, in fact, be successfully embedded into supergravity based on a shift-symmetric Kähler potential—provided that supersymmetry is broken at a high scale and broken in a different sector. That is, once the inflaton is *not* identified as the SUSY-breaking Polonyi field, the inflaton may receive a large soft mass that compensates its tachyonic tree-level mass.

This result leaves a near-canonical tree-level Kähler potential as the best possible choice,

$$K = \Phi^\dagger\Phi + \frac{\epsilon}{(2!)^2}\left(\frac{\Phi^\dagger\Phi}{M_{\rm Pl}}\right)^2 + \mathcal{O}(\epsilon^2, M_{\rm Pl}^{-4}). \qquad (63)$$

Of course, simply assuming a canonical Kähler potential, $K = \Phi^\dagger\Phi$, would represent an even more minimal choice. In that case, the Polonyi field would not "equip itself" with a Hubble-induced mass (see our discussion at the beginning of this section) and we would not have to fear any dangerous SUGRA corrections. But such a scenario is rather unlikely. Terms featuring higher powers of $\Phi^\dagger\Phi$ are not forbidden by any symmetry and, hence, we cannot exclude their presence in the Kähler potential. Quite the contrary, we rather expect that they are the unavoidable consequence of radiative corrections in quantum gravity that manifest themselves as Planck-suppressed operators in the low-energy effective theory. The only open question regarding the higher-dimensional terms in the Kähler potential pertains to the size of their coefficients. The leading correction to the canonical Kähler potential, e.g., results in an inflaton mass squared of $m_\varphi^2 = -3\epsilon H_{\rm inf}^2$, which shifts the slow-roll parameter η by $\Delta\eta = -\epsilon$. Therefore, if ϵ is of $\mathcal{O}(1)$, as one might naively expect, we would encounter the η problem [57]. This tells us that the coefficient ϵ needs to be suppressed, $\epsilon \lesssim 0.1$, which introduces a parameter fine-tuning at the level of $1\ldots10\%$. In the following, we will not speculate what the physical reason for this fine-tuning might be and simply assume that ϵ is sufficiently small. We merely remark that there is a number of models in the literature that attempt to explain why ϵ ends up being suppressed. For instance, a small coefficient ϵ may be the result of a sub-Planckian cutoff scale, $M_* \lesssim 0.1 M_{\rm Pl}$, in the sector giving rise to the higher-dimensional term in Eq. (63) [84] (see [85] for a realization of this idea in string theory), or ϵ may be small because it is, for one reason or another, suppressed by NDA factors of $\mathcal{O}(4\pi)$ [52].

Having defined the Kähler potential, we now need to canonically normalize the complex Polonyi field $\tilde\phi = \tilde\varphi/\sqrt{2}e^{i\tilde\theta}$. Here, we will restrict ourselves to normalizing the radial component $\tilde\varphi$ only, as the complex phase $\tilde\theta$ will be irrelevant during inflation. The canonically normalized field φ is then given as

$$\begin{aligned}\varphi(\tilde\varphi) &= \int d\tilde\varphi\left(\frac{\partial^2 K}{\partial\tilde\phi\partial\tilde\phi^*}\right)^{1/2}\\ &= \tilde\varphi\left[1 + \frac{\epsilon}{12}\left(\frac{\tilde\varphi}{M_{\rm Pl}}\right)^2 + \mathcal{O}(\epsilon^2)\right]. \end{aligned}\qquad (64)$$

This relation is readily inverted, which provides us with an expression for $\tilde\varphi$ in terms of φ,

$$\tilde\varphi(\varphi) = \varphi\left[1 - \frac{\epsilon}{12}\left(\frac{\varphi}{M_{\rm Pl}}\right)^2 + \mathcal{O}(\epsilon^2)\right]. \qquad (65)$$

We are now in the position to write down the full tree-level SUGRA potential V_F for φ and $\tilde\theta$. In doing so, we shall first use the full Polonyi superpotential in Eq. (6) for illustrative purposes (i.e., we will add a constant w to the effective superpotential, although $w = 0$ in the IYIT model). Equation (50) then yields

$$\begin{aligned}V_F(\varphi,\tilde\theta) = c_0 + c_1\varphi\cos\tilde\theta + \frac{c_2}{2}\varphi^2 + \frac{c_3}{3!}\varphi^3\cos\tilde\theta\\ + \frac{c_4}{4!}\varphi^4 + \mathcal{O}(\varphi^5),\end{aligned}\qquad (66)$$

where the coefficients c_0, c_1, c_2, c_3, and c_4 are given as follows:[12]

$$\begin{aligned}&c_0 = V_0 - 3\frac{w^2}{M_{\rm Pl}^2}, \qquad c_1 = -2\sqrt{2}\frac{w}{M_{\rm Pl}^2}V_0^{1/2},\\ &c_2 = -\epsilon\frac{V_0}{M_{\rm Pl}^2} - 2\frac{w^2}{M_{\rm Pl}^4}, \qquad c_3 = -3\sqrt{2}\left(1+\frac{\epsilon}{6}\right)\frac{w}{M_{\rm Pl}^4}V_0^{1/2},\\ &c_4 = 3\left(1 - \frac{7\epsilon}{2} + \frac{8\epsilon^2}{3}\right)\frac{V_0}{M_{\rm Pl}^4} - 3\left(1+\frac{\epsilon}{6}\right)\frac{w^2}{M_{\rm Pl}^6}.\end{aligned}\qquad (67)$$

This form of the scalar potential illustrates the impact of the constant term in the superpotential, $W \supset w$. Not only does w induce a dependence on the phase $\tilde\theta$ of the complex inflaton field (thereby introducing odd powers of the radial component φ in the scalar potential), it also results in a large tachyonic tree-level mass. In Appendix B 1, we show that this potential is always too steep for slow-roll inflation. Moreover, we note that the potential in Eq. (66) is similar to the tree-level potential of F-term hybrid inflation. Equation (66) yields the SUGRA FHI tree-level potential up to $\mathcal{O}(\varphi^4)$, if we make the following replacements:

$$\begin{aligned}&c_0 \to c_0 + 3\frac{w^2}{M_{\rm Pl}^2}, \qquad c_1 \to c_1, \qquad c_2 \to c_2 + 3\frac{w^2}{M_{\rm Pl}^4},\\ &c_3 \to c_3, \qquad c_4 \to c_4 + 9\left(1-\frac{\epsilon}{6}\right)\frac{w^2}{M_{\rm Pl}^6},\end{aligned}\qquad (68)$$

where the constant w is to be understood as a measure for the gravitino mass, $w = m_{3/2}M_{\rm Pl}^2$. In F-term hybrid inflation, the functional dependence of the scalar potential on the phase of the complex inflaton field is, therefore, exactly the same as in the full Polonyi model (see [65] for the first analysis that properly treats F-term hybrid inflation as a

[12] A dimension-6 operator in the Kähler potential, $K \supset \epsilon'/(3!)^2(\Phi^\dagger\Phi)^3/M_{\rm Pl}^4$, only contributes to the coefficient of the quartic inflaton term, $c_4 \to c_4 - 3/2\epsilon' V_0/M_{\rm Pl}^4$. This effect is negligibly small, which is why we will set $\epsilon' = 0$ in the following.

two-field model in the complex plane). By contrast, the soft inflaton mass is no longer tachyonic in F-term hybrid inflation, but identical to the gravitino mass. Meanwhile, the change in the coefficient c_4 is suppressed by $m_{3/2}^2/M_{\rm Pl}^2$ and, thus, irrelevant. Furthermore, it is worthwhile pointing out that F-term hybrid inflation is *not* compatible with the idea of a gravitino mass as large as the inflationary Hubble rate. Instead, $m_{3/2}$ and $H_{\rm inf}$ must always maintain a hierarchy of at least three orders of magnitude, $m_{3/2} \lesssim 10^{-3} H_{\rm inf}$, since otherwise slow-roll inflation becomes spoiled by too large SUGRA corrections [65]. This is in contrast to our scenario of Polonyi inflation, where $m_{3/2}$ and $H_{\rm inf}$ are related to each other at a fundamental level and, in fact, equal to each other [see Eq. (3)].

Let us now turn to the scalar potential of *our* inflationary scenario. For our purposes, the constant w needs to be zero during inflation. The total inflaton potential for Polonyi inflation, therefore, follows from the sum of the logarithmic one-loop potential in Eq. (45) and V_F in Eq. (66) evaluated for $w = 0$,

$$V(\varphi) = V_0 - \frac{\epsilon}{2}\frac{V_0}{M_{\rm Pl}^2}\varphi^2 + \frac{1}{8}\left(1 - \frac{7\epsilon}{2} + \frac{8\epsilon^2}{3}\right)\frac{V_0}{M_{\rm Pl}^4}\varphi^4 + \Lambda_{\rm HE}^4 \ln\left(\frac{\varphi}{\varphi_c}\right) + \mathcal{O}(x^{-4}, \varphi^6). \tag{69}$$

This inflaton potential is identical to the potential of F-term hybrid inflation (including corrections from supergravity and a noncanonical Kähler potential) in the limit of a vanishing gravitino mass [64].

III. SPONTANEOUS *R* SYMMETRY BREAKING AFTER THE END OF INFLATION

A. Implications of broken *R* symmetry for *R* parity and the gravitino mass

In the previous section, we have seen what it takes to identify the SUSY-breaking Polonyi field Φ as the chiral inflaton field. We considered three different types of radiative corrections to the scalar Polonyi potential (none, quadratic corrections, logarithmic corrections), two different shapes of the tree-level Kähler potential (approximately shift-symmetric and near-canonical), and two different choices for the constant in the superpotential ($w = 0$ and $w \neq 0$)—and we found only one possible combination of "model ingredients" that could not be ruled out immediately: the Polonyi model supplemented by logarithmic radiative corrections in combination with a near-canonical Kähler potential and $w = 0$ during inflation.

In the following, we will focus on this scenario and present a dynamical model that illustrates how the constant in the superpotential, $W \supset w$, may be generated at the end of inflation. As we will see, the constant w then represents (not the only one, but) the dominant contribution to the VEV of the superpotential in the low-energy vacuum, $\langle W \rangle \simeq w$. Consequently, w turns out to be the main source of spontaneous R symmetry breaking at low energies. This is the reason why, throughout the entire paper, we refer to the generation of the constant term in the superpotential after inflation also as the spontaneous breaking of R symmetry at late times. Technically speaking, this is not quite correct, as, globally, R symmetry is already spontaneously broken during inflation by the VEV of the Polonyi field,

$$\langle \Phi \rangle \neq 0 \Rightarrow Z_4^R \to Z_2^R. \tag{70}$$

This is because Φ carries R charge 2. What we actually mean by speaking of late-time R symmetry breaking is, therefore, not the spontaneous breaking of R symmetry from a global perspective, but only the spontaneous breaking of R symmetry in the sector responsible for the generation of the constant term w. From the perspective of the low-energy effective theory, the process of late-time R symmetry breaking then amounts to the (explicit) breaking of R symmetry at the level of the Lagrangian by means of a constant in the superpotential, as opposed to the (spontaneous) breaking of R symmetry at the level of a scalar field VEV during inflation. Here, it is important to note that the latter kind of R symmetry breaking mostly vanishes in the true vacuum after inflation, whereas the former kind of R symmetry also remains at low energies. Likewise, we shall refer to the sector responsible for the generation of the constant w (i.e., the sector responsible for R symmetry breaking at low energies) as the R-symmetry-breaking sector for short. We also point out that the temporary breaking of R symmetry by the VEV of the inflaton field is, in fact, mandatory from a phenomenological point of view. Otherwise, the spontaneous breaking of R symmetry during the generation of the constant w would result in the production of domain walls at the end of inflation, which would render our scenario phenomenologically unviable [86].[13]

Similar to the Polonyi VEV, the constant w also breaks the Z_4^R symmetry down to a Z_2^R parity,

$$w \neq 0 \Rightarrow Z_4^R \to Z_2^R, \tag{71}$$

so that, during inflation as well as in the low-energy vacuum after inflation, the only discrete symmetry that remains globally unbroken corresponds to an exact Z_2^R. Interestingly enough, this parity is suited to be identified as

[13]If the inflationary dynamics did not break R symmetry (in the context of some *other* model), R symmetry would need to be broken in a separate sector way before the end of inflation, to make sure that all domain walls are sufficiently diluted. Such an alternative scenario provides a different explanation for why supersymmetry is necessarily broken at a high scale [87].

the R parity of the MSSM, P_R [88]. Here, note that R parity is not a proper R symmetry in the actual sense, as it is equivalent to matter parity, P_M, which is a non-R symmetry [31]. Our model, therefore, automatically offers a possible explanation for the origin of R parity in the MSSM. In contrast to other models, it does not depend on extra continuous symmetries, such as a global or local $U(1)_{B-L}$, that would leave behind $P_M \cong P_R$ as a discrete remnant subgroup after spontaneous symmetry breaking.

Given the above sources of R symmetry breaking during and after inflation, the physical gravitino mass $m_{3/2}$ varies as a function of time. In general, $m_{3/2}$ is controlled by the VEV of the superpotential,

$$m_{3/2} = \exp\left[\frac{\langle K\rangle}{2M_{\rm Pl}^2}\right]\frac{\langle W\rangle}{M_{\rm Pl}^2}, \qquad W = W_{\rm eff} + w. \tag{72}$$

During inflation, the constant w vanishes and, thus, $m_{3/2}$ turns out to be roughly proportional to the Polonyi field value. Meanwhile, after inflation, $m_{3/2}$ is dominated by the contribution from the constant w. In order to tune the CC to zero, we have to require that $m_{3/2}$ eventually takes the following value:

$$m_{3/2} = \frac{\Lambda_{\rm SUSY}^2}{\sqrt{3}M_{\rm Pl}}, \qquad \Lambda_{\rm SUSY} = \langle |F|\rangle^{1/2}. \tag{73}$$

This condition can always be satisfied by fine-tuning the constant w to the particular value w_0,

$$w_0 = \exp\left[-\frac{\langle K\rangle}{2M_{\rm Pl}^2}\right]\frac{M_{\rm Pl}}{\sqrt{3}}\langle |F|\rangle_{w=w_0} - \langle W_{\rm eff}\rangle. \tag{74}$$

Here, it is important to note that the SUSY-breaking F-term $|F|$ also depends on the constant w via the covariant derivative of the superpotential [see Eq. (52)]. For this reason, the above relation is, in fact, only an implicit definition of w_0, which, in principle, still needs to be solved for w_0. In practice, the dependence of $|F|$ on w is, however, negligible in most cases, so that the right-hand side of Eq. (74) readily yields the required value for w. In the following, we shall now show how the constant w may be generated after inflation (see Secs. III B and III C) and discuss under which conditions it may be successfully matched with the desired value w_0, so that the CC vanishes at low energies (see Sec. III D).

B. Gaugino condensation in a mass-deformed strongly coupled hidden sector

The simplest way to break R symmetry via strong dynamics is to make use of dynamical gaugino condensation in a pure SYM theory [67]. For instance, in a strongly coupled pure SYM theory based on $SU(N_c)$ and associated with a dynamical scale Λ', gaugino condensation results in an R-symmetry-breaking constant w of $\mathcal{O}(\Lambda'^3)$. Let us now illustrate how we can use this property for our purposes.

Our starting point is supersymmetric quantum chromodynamics (SQCD): a strongly coupled $SU(N_c)$ gauge theory with N_c colors and $N_f < 3N_c$ flavors, where each flavor consists of a quark/antiquark pair $\{Q^i, \bar{Q}^i\}$. We will constrain the viable values for N_c and N_f shortly. For the moment, however, let us stay as general as possible and leave the concrete values of N_c and N_f unspecified. We assume that all quark/antiquark pairs are *a priori* massless, i.e., we do *not* introduce any bare mass terms of the form $m_i Q^i \bar{Q}^i$ in the tree-level superpotential. Instead, we suppose that all of the N_f flavors share a Yukawa coupling with some $SU(N_c)$ singlet field P, which, thus, results in field-dependent masses for all flavors,

$$W^R = c_i P Q^i \bar{Q}^i. \tag{75}$$

Here, the coefficients c_i denote dimensionless Yukawa couplings of $\mathcal{O}(1)$. By varying the VEV of the singlet field P, we are then able to control the quark mass matrix, M_Q, in our SQCD theory,

$$M_Q = {\rm diag}\{M_{Q^i}\} = {\rm diag}\{c_i\langle P\rangle\}. \tag{76}$$

As long as P is stabilized at the origin, $\langle P\rangle = 0$, all quarks are massless and the SQCD sector remains what it is: an $SU(N_c)$ gauge theory with N_f dynamical flavors, described by a quantum moduli space at low energies [60].[14] On the other hand, once P acquires a large VEV, all flavors become heavy and can be integrated out. This mass deformation transforms the SQCD sector into a pure SYM theory, in which R symmetry is spontaneously broken: First, R symmetry is broken down to $Z_{2N_c}^R$ due to $SU(N_c)$ instantons. Then, the discrete $Z_{2N_c}^R$ symmetry is broken further down to Z_2^R via gaugino condensation. This results in the eagerly anticipated constant term w in the effective superpotential,

$$W_{\rm eff}^R = w = \frac{1}{\tilde{\eta}^2} N_c \tilde{\Lambda}_{\rm eff}^3. \tag{77}$$

Here, $\tilde{\eta}$ is a dimensionless "fudge factor" that encompasses all numerical factors entering into the expression for w except for N_c and $\tilde{\Lambda}_{\rm eff}^3$. Based on NDA, we again expect $\tilde{\eta}$ to be of $\mathcal{O}(4\pi)$, which is why we will set $\tilde{\eta} = 4\pi$ in the following, for definiteness. Meanwhile, $\tilde{\Lambda}_{\rm eff}$ denotes the effective dynamical scale of the SYM theory after integrating out all heavy flavors. We obtain an expression for $\tilde{\Lambda}_{\rm eff}$

[14] During inflation, the scalar quark/antiquark fields acquire Hubble-induced masses. In the special case of $N_f < N_c$, these stabilize the runaway directions on the quantum moduli space induced by the nonperturbative Affleck-Dine-Seiberg (ADS) superpotential [50,89].

by matching the running of the $SU(N_c)$ gauge coupling constant $\tilde{g}$ at the respective heavy-quark mass thresholds,

$$\tilde{\Lambda}_{\rm eff}^{3N_c} = \bar{M}_Q^{N_f}\tilde{\Lambda}^{3N_c-N_f}, \tag{78}$$

where $\tilde{\Lambda}$ denotes the dynamical scale of the original high-energy theory and where $\bar{M}_Q$ represents the effective quark mass scale, i.e., the geometric mean of all quark mass eigenvalues M_{Q^i},

$$\bar{M}_Q = \left(\prod_i M_{Q^i}\right)^{1/N_f} = (\det M_Q)^{1/N_f} = \bar{c}\langle P\rangle,$$
$$\bar{c} = \left(\prod_i c_i\right)^{1/N_f}. \tag{79}$$

In the following, we will set $\bar{c}$ (the geometric mean of all Yukawa couplings c_i) to $\bar{c}=1$, for simplicity. For given values of N_c and N_f, the constant w in Eq. (77) then ends up being a function of $\langle P\rangle$ and $\tilde{\Lambda}$,

$$w = \frac{N_c}{16\pi^2}\langle P\rangle^{N_f/N_c}\tilde{\Lambda}^{3-N_f/N_c}. \tag{80}$$

Here, $\langle P\rangle = \bar{M}_Q$ must be larger than $\tilde{\Lambda}$ to ensure that the heavy quark flavors can be integrated out perturbatively in the high-energy theory. Our result for w in Eq. (80) is, therefore, bounded from above,

$$\tilde{\Lambda} \lesssim \langle P\rangle \Rightarrow w \lesssim w_{\rm max} = \frac{N_c}{16\pi^2}\langle P\rangle^3. \tag{81}$$

The dynamical scale of the high-energy theory, $\tilde{\Lambda}$, is generated via dimensional transmutation and, hence, solely depends on the value of the $SU(N_c)$ gauge coupling constant $\tilde{g}$ at the Planck scale,

$$\tilde{\Lambda} = M_{\rm Pl}\exp\left[-\frac{8\pi^2}{b}\frac{1}{\tilde{g}^2(M_{\rm Pl})}\right], \qquad b = 3N_c - N_f. \tag{82}$$

We assume the $SU(N_c)$ beta function coefficient b to be positive, which means that $\tilde{\Lambda}$ can basically take any desired value below the Planck scale. For a given VEV $\langle P\rangle$, the constant w can then be tuned simply by varying the value of the gauge coupling constant $\tilde{g}$ at the UV cutoff scale. From this perspective, the fine-tuning of the CC in the true vacuum turns into an issue that pertains to the UV boundary conditions of the low-energy effective theory. It is no longer a problem within the low-energy effective theory itself.

How can we now apply these results in the context of Polonyi inflation? The crucial idea is to relate the value of the VEV $\langle P\rangle$ to the inflationary dynamics. We must make sure that the singlet P is stabilized at $\langle P\rangle = 0$ during inflation and that it becomes destabilized only towards the end of inflation. In other words, we need to achieve a situation in which small inflaton field values trigger a mass deformation of the SQCD sector, $\bar{M}_Q = 0 \to \bar{M}_Q \neq 0$, so that this sector turns into a pure SYM theory and spontaneously breaks R symmetry via gaugino condensation. This approach shares some similarities with the models presented in [90], in which small inflaton field values trigger the mass generation in a separate Intriligator-Seiberg-Shih (ISS) sector [43], which then spontaneously breaks supersymmetry. Moreover, our scenario of R symmetry breaking should be compared with the approaches in [74,91] and [92]. In [74,91], inflation ends in a supersymmetric, but R-symmetry-breaking ground state, while in [92], R symmetry is broken by a "field-dependent constant" term in the superpotential that only becomes large at the end of inflation.

C. Triggering a late-time mass deformation by a small inflaton field value

We can realize the scenario described in the previous section by introducing the following superpotential:

$$W^P = \alpha Y\left(v_P^2 - \frac{1}{2}P^2\right) + \frac{\beta}{3!}Y^3. \tag{83}$$

Here, Y is a singlet field that carries R charge 2, α and β are dimensionless coupling constants, and v_P is a mass scale. It is tempting to speculate that also the scale v_P might be dynamically generated in a strongly coupled hidden sector, similar to the scale v in Eq. (35). However, for our purposes, it will not be necessary to specify its origin. Instead, we will treat it as a hard dimensionful input scale, the only one in our model. Before we explain in more detail how W^P allows us to give a VEV to the field P at the end of inflation, let us examine for which values of N_c and N_f the superpotential in Eq. (83) is consistent with the (approximate) global symmetries of the theory in the first place. Promoting these global symmetries to gauge symmetries (and/or assuming that they are at least sufficiently protected by other gauge symmetries), we will then be able to narrow down the viable choices for N_c and N_f.

The superpotential in Eq. (75) is invariant under an anomaly-free global $U(1)_B \times Z_{2N_f} \times U(1)_R$ flavor symmetry, under which the fields Q^i, $\bar{Q}^i$, P, and Y are charged as follows:

	Q^i	$\bar{Q}^i$	P	Y
$U(1)_R$	$1-N_c/N_f$	$1-N_c/N_f$	$2N_c/N_f$	2
$U(1)_B$	1/3	−1/3	0	0
Z_{2N_f}	1	1	−2	0

As mentioned earlier, in the pure SYM theory after integrating out the heavy quarks, the continuous $U(1)_R$ is broken to a discrete $Z_{2N_c}^R$ symmetry by $SU(N_c)$ instantons. This immediately suggests to set the number of colors

N_c to 2, so that the global $Z_{2N_c}^R$ symmetry of the pure SYM theory can be identified with our gauged Z_4^R symmetry.[15] According to the above table, the operator YP^2 in Eq. (83) then only complies with the exact $Z_{2N_c}^R \equiv Z_4^R$ symmetry, if we choose the number of flavors N_f to be either 1 or 2,

$$[YP^2]_R = 2 + \frac{4N_c}{N_f} = 2 \mod 2N_c \Rightarrow N_f = 1, 2. \quad (84)$$

On second sight, $N_f = 1$ is, however, not a viable option. In that case, the R charge of the field P vanishes, $[P]_R = 0$, so that we cannot prevent the unwanted operator YP from popping up in the superpotential. Also the approximate discrete $Z_{2N_f} \equiv Z_2$ flavor symmetry cannot help us in this situation, since the singlet field P is required to transform even, $[P]_{Z_2} = 0$, under this parity. At the same time, the singlet field Y must also transform even, $[Y]_{Z_2} = 0$, to allow for the presence of the Y tadpole term in Eq. (83). The discrete Z_2 flavor symmetry, therefore, does not allow us to eliminate the operator YP, as long as we want to keep the Y tadpole term. This leaves us with the following unique choice for N_c and N_f:

$$N_c = 2, \qquad N_f = 2, \quad (85)$$

which results in exactly the same matter and gauge field content as in the IYIT sector.

For $N_c = N_f = 2$, the singlet field P carries R charge 2. From the perspective of R symmetry, the fields P and Y are, hence, indistinguishable from each other, so that, next to the three wanted operators Y, YP^2, and Y^3, the three unwanted operators P, PY^2, and P^3 are also allowed to appear in the superpotential. We are, therefore, led to impose the $Z_{2N_f} \equiv Z_4$ flavor symmetry as an approximate symmetry—which we assume to be protected by some exact gauge symmetry in the UV.[16] This approximate Z_4 acts on the fields P and Y again as a Z_2 parity—the crucial difference to the $N_f = 1$ case being that, now, P transforms odd, $[P]_{Z_2} = 1$, while Y still transforms even, $[Y]_{Z_2} = 0$. By virtue of this Z_2 parity, the coefficients of the unwanted operators P, PY^2, and P^3 then end up being suppressed. Moreover, we emphasize the importance of the fact that the Z_4 symmetry is only an approximate symmetry. If it was an exact symmetry, its spontaneous breaking at the end of inflation would result in the formation of dangerous domain walls. As the Z_4 symmetry is, however, bound to be explicitly broken (e.g., by the suppressed Z_2-odd operators P, PY^2, and P^3), we are safe from running into this problem.

Given the superpotential in Eq. (83), let us now demonstrate how it enables us to trigger a mass deformation of the R-symmetry-breaking sector towards the end of inflation. The first thing to note is that the superpotential in Eq. (83) results in two tachyonic mass eigenstates in global supersymmetry,

$$\text{Global SUSY}: m_{p^\pm}^2 = \pm\alpha^2 v_P^2, \quad m_{y^\pm}^2 = \pm\alpha\beta v_P^2, \quad (86)$$

where $p^\pm$ and $y^\pm$ denote the real scalar degrees of freedom (DOFs) contained in P and Y, respectively. These masses, however, receive corrections in supergravity that may render them nontachyonic,

$$\begin{aligned} m_{p^\pm}^2(\varphi) &= \pm\alpha^2 v_P^2 + \left(\frac{V(\varphi)}{M_{\rm Pl}^2} + \frac{\Delta V_P}{M_{\rm Pl}^2}\right) + \Delta m_{p^\pm}^2(\varphi), \\ \Delta V_P &= \alpha^2 v_P^4, \\ m_{y^\pm}^2(\varphi) &= \pm\alpha\beta v_P^2 + \left(\frac{V(\varphi)}{M_{\rm Pl}^2} + \mathcal{O}(M_{\rm Pl}^{-4})\right) + \Delta m_{y^\pm}^2(\varphi). \end{aligned} \quad (87)$$

Here, the corrections proportional to $V(\varphi)$ are nothing but the usual Hubble-induced masses that scalar fields typically acquire during inflation. Apart from that, we collect all further corrections that explicitly depend on φ in the field-dependent mass contributions $\Delta m_{p^\pm}^2(\varphi)$ and $\Delta m_{y^\pm}^2(\varphi)$. Just like the Hubble-induced masses, these corrections result from the F-term tree-level potential in supergravity.

As an important aside, we point out that the Hubble-induced masses in Eq. (87) are, in fact, sensitive to higher-dimensional operators in the Kähler potential between the fields P and Y and the Polonyi field,

$$\begin{aligned} K_{\rm mix} &= F_P(\Phi/M_*, \Phi^\dagger/M_*)|P|^2 \\ &\quad + F_Y(\Phi/M_*', \Phi^\dagger/M_*')|Y|^2 \end{aligned} \quad (88)$$

where F_P and F_Y denote two unknown functions of Φ and $\Phi^\dagger$ and where M_* and M_*' represent appropriate cutoff scales. Note that such couplings between P, Y, and Φ in the

[15]Here, we suppose some kind of Kähler interaction between the two hidden sectors, such as, e.g., $K \supset \epsilon' \Phi P^\dagger + \text{H.c.}$, that relates R symmetry transformations in the one sector to R symmetry transformations in the other sector. Besides that, we assume the coefficients of these mixing operators to be negligibly small for all practical purposes, $\epsilon' \lll 1$. This might, e.g., be achieved if the cross terms in the Kähler potential are suppressed by (several powers of) a large cutoff scale.

[16]If we did *not* identify the $Z_{2N_c}^R$ of the pure SYM theory with our anomaly-free Z_4^R, i.e., if we allowed for $N_c > 2$ (while keeping $N_f = 2$ fixed), we would not be forced to impose any *additional* approximate flavor symmetry. In that case, the $Z_{2N_c}^R$ alone would suffice to restrict the set of allowed operators to Y and YP^2 (meaning that the Y^3 term would be suppressed). Such a scenario would, however, not present an advantage over the $N_c = 2$ case. Either way, we have to assume a sufficiently intact global symmetry. While for $N_c = 2$ the Z_4 needs to be imposed as an approximate symmetry, the $Z_{2N_c}^R$ itself would need to be play the role of an approximate symmetry for $N_c > 2$. We will, therefore, ignore this possibility in the following.

Kähler potential are the only relevant means of communication between the SUSY-breaking sector and the R-symmetry-breaking sector in our model. For instance, we could imagine that both sectors only communicate with each other via gravitational interactions. In that case, the scales M_* and M'_* should be identified as the Planck scale. Alternatively, we may assume that the R-symmetry-breaking sector couples to the SUSY-breaking sector via the exchange of heavy messenger particles of mass $M_{\rm mess}$. Then, $M_*^{(\prime)}$ would have to be identified as a typical mass scale of the messenger sector, $M_*^{(\prime)} \sim M_{\rm mess}$. To illustrate the effect of the Kähler potential in Eq. (88), let us suppose that the leading-order contributions to F_P and F_Y take the following form:

$$K_{\rm mix} = \epsilon_P \frac{|\Phi|^2|P|^2}{M_*^2} + \epsilon_Y \frac{|\Phi|^2|Y|^2}{M_*^2} + \dots. \qquad (89)$$

Such "mixing terms" shift the Hubble-induced masses in Eq. (87). Consider, e.g., the mass of p^-,

$$m_{p^-}^2(\varphi) \subset \frac{V(\varphi)}{M_{\rm Pl}^2} \to \left(1 - \epsilon_P \frac{M_{\rm Pl}^2}{M_*^2} + \frac{\epsilon_P^2}{2}\frac{M_{\rm Pl}^4}{M_*^4}\frac{\varphi^2}{M_{\rm Pl}^2} + \dots\right)\frac{V(\varphi)}{M_{\rm Pl}^2}, \qquad (90)$$

where the ellipsis stands for further, subdominant corrections. $m_{p^+}^2$ and $m_{y^\pm}^2$ are modified in a similar way. For negative coefficients ϵ_P and ϵ_Y, these corrections may help establish a strong hierarchy between the masses $m_{p^\pm}$ and $m_{y^\pm}$ on the one hand and the inflationary Hubble rate $H_{\rm inf}$ on the other hand,

$$-\epsilon_{P,Y}\frac{M_{\rm Pl}^2}{M_*^2} \gg 1 \Rightarrow |m_{p^\pm}|, |m_{y^\pm}| \gg H_{\rm inf}, \qquad (91)$$

so as to make sure that, during inflation, all of the scalar fields $p^\pm$ and $y^\pm$ are sufficiently stabilized. At the same time, it is clear that the Kähler potential in Eq. (88) significantly complicates the analysis of the R-symmetry-breaking sector. For presentational purposes, we will, therefore, neglect the effect of $K_{\rm mix}$ in the following. We have checked explicitly that, accounting for the exemplary Kähler potential in Eq. (89), all results remain quantitatively the same—the only difference being the hierarchy in Eq. (91). In this sense, all results in this and the following section should be regarded as approximate results that convey the general flavor of our mechanism of late-time R symmetry breaking. An exact numerical study including a more systematic treatment of the functions F_P and F_Y in Eq. (88) is left for future work. For our purposes, it is merely important to remember that the masses of the fields $p^\pm$ and $y^\pm$ during inflation may be further increased by adding higher-dimensional terms to the Kähler potential.

The idea now is to choose the coupling constants α and β such that the scalar DOF p^- is stabilized during inflation, turning tachyonically unstable only at a certain (small) inflaton field value φ_0,

$$\varphi \geq \varphi_0 \Rightarrow m_{p^+}^2(\varphi) > 0, \quad m_{p^-}^2(\varphi) \geq 0, \quad m_{y^\pm}^2(\varphi) > 0,$$
$$\varphi < \varphi_0 \Rightarrow m_{p^+}^2(\varphi) > 0, \quad m_{p^-}^2(\varphi) < 0, \quad m_{y^\pm}^2(\varphi) > 0. \quad (92)$$

To determine the required values for α and β in order to implement such a scenario, we first note that the field-dependent mass $m_{p^-}^2$ in Eq. (87) vanishes at $\varphi = 0$ for the following choice of α:

$$m_{p^-}^2(0) = -\alpha_0^2 v_P^2 + \frac{V_0}{M_{\rm Pl}^2} + \frac{\alpha_0^2 v_P^4}{M_{\rm Pl}^2} = 0$$
$$\Rightarrow \alpha_0 = \frac{\mu^2}{v_P(M_{\rm Pl}^2 - v_P^2)^{1/2}}. \qquad (93)$$

In other words, for $\alpha = \alpha_0$, the Hubble-induced mass compensates the tachyonic mass, $-\alpha^2 v_P^2$, for all field values during inflation, i.e., as long as $\varphi > 0$, and only at the origin the scalar p^- becomes massless. At the same time, we need to make sure that the scalar DOFs contained in Y are stabilized at all times. This is achieved by requiring that their respective masses in global supersymmetry, $\pm\alpha\beta v_P^2$, are always outweighed by the Hubble-induced mass, even at the origin, where the Hubble-induced mass is smallest,

$$|\alpha\beta v_P^2| < \frac{V_0}{M_{\rm Pl}^2} \Rightarrow |\beta| < \beta_0 = \frac{V_0}{\alpha v_P^2 M_{\rm Pl}^2}. \qquad (94)$$

As long as this condition is satisfied, the field Y is stabilized at a nonzero VEV,[17]

$$\langle Y\rangle = \frac{\alpha v_P^2 \mu^2}{m_{y^+}^2 M_{\rm Pl}^2}\frac{\varphi}{\sqrt{2}} + \mathcal{O}(\varphi^5), \qquad (95)$$

which induces a positive Majorana mass for the field P via the waterfall superpotential in Eq. (83),

$$m_{p^\pm}(\varphi) \supset \alpha\langle Y\rangle. \qquad (96)$$

This field-dependent mass is contained in $\Delta m_{p^\pm}^2(\varphi)$ in Eq. (87) and, similarly as in ordinary F-term hybrid inflation, it helps to stabilize the field P during inflation.

[17] In global supersymmetry, the VEV $\langle Y\rangle$ vanishes during and after inflation. In supergravity, the scalar potential, however, contains a tadpole term linear in y^+ (due to the nonzero F-term of the field Y) that displaces $\langle Y\rangle$ from the origin.

Only at the end of inflation, $\Delta m^2_{p^\pm}(\varphi)$ and, in particular, the mass term in Eq. (96) vanish, so that P has a chance of becoming unstable.

Let us now study the inflaton-dependent mass corrections in Eq. (87) more closely. In doing so, it turns out to be convenient to use α_0 and β_0 as reference values for α and β. To this end, we define

$$a = \left(\frac{\alpha}{\alpha_0}\right)^2 = \frac{\alpha^2 v_P^2 - \alpha^2 v_P^4/M_{\rm Pl}^2}{V_0/M_{\rm Pl}^2}, \qquad b = \frac{\beta}{\beta_0} = \frac{\alpha\beta v_P^2}{V_0/M_{\rm Pl}^2}. \tag{97}$$

The field-dependent shift $\Delta m^2_{p^-}(\varphi)$ in the mass of the scalar DOF p^- can then be written as follows:

$$\Delta m^2_{p^-}(\varphi) = \left[1 - \frac{2a}{(1+b)^2}\left(1 - \frac{a}{2} + \frac{7b}{4} + \frac{b^2}{2}\right) + \mathcal{O}(M_{\rm Pl}^{-2})\right] \times \frac{V_0}{M_{\rm Pl}^2}\frac{\varphi^2}{2M_{\rm Pl}^2} + \mathcal{O}(\varphi^4). \tag{98}$$

The parameter a is supposed to be chosen, such that the total mass $m^2_{p^-}$ vanishes at $\varphi = \varphi_0$,

$$m^2_{p^-}(\varphi_0) = (1 - a(\varphi_0))\frac{V_0}{M_{\rm Pl}^2} + \frac{1}{2}m_{\rm eff}^2\left(\frac{\varphi_0}{M_{\rm Pl}}\right)^2 + \Delta m^2_{p^-}(\varphi_0) = 0. \tag{99}$$

For $\varphi_0 \ll M_{\rm Pl}$, this condition yields a quadratic equation for $a(\varphi_0)$, which has two independent solutions,

$$a_1(\varphi_0) = 1 + \frac{1}{2}\frac{m_{\rm eff}^2\varphi_0^2}{V_0} - \frac{3b}{4(1+b)^2}\left(\frac{\varphi_0}{M_{\rm Pl}}\right)^2 + \mathcal{O}(\varphi_0^4), \qquad a_2(\varphi_0) = 2(1+b)^2\left(\frac{M_{\rm Pl}}{\varphi_0}\right)^2 + \mathcal{O}(\varphi_0^0). \tag{100}$$

Here, the first solution is close to unity, $a_1(\varphi_0) \approx 1$, so that $\alpha(\varphi_0) = a^{1/2}(\varphi_0)\alpha_0 \approx \alpha_0$. For $a = a_1(\varphi_0)$, the Hubble-induced mass $V/M_{\rm Pl}^2$ cancels the tachyonic mass, $-\alpha^2 v_P^2$, to good precision and the field-dependent correction $\Delta m^2_{p^-}(\varphi_0)$ nearly vanishes on its own. The second solution is, by contrast, much larger, $a_2(\varphi_0) \gg 1$. In that case, the individual contributions to $m^2_{p^-}(\varphi_0)$ do not vanish independently. Instead, all contributions are large, have different signs, and simply happen to cancel each other at $\varphi = \varphi_0$. As we will see further below, the solution a_2 will turn out to be unviable for phenomenological reasons. For now, we will, however, remain impartial and treat both solutions on the same footing.

Combining our results in Eqs. (93), (94), (97), and (100), we eventually obtain our final expressions for the parameter α. Again, we find two solutions,

$$\alpha_1(\varphi_0) = \frac{1}{v_P}\frac{V_{\rm eff}^{1/2}(\varphi_0)}{(M_{\rm Pl}^2 - v_P^2)^{1/2}} - \frac{3\beta\varphi_0^2}{8(M_{\rm Pl}^2 - v_P^2)} + \mathcal{O}(\beta^2, \varphi_0^4),$$
$$V_{\rm eff}(\varphi_0) = V_0 + \frac{1}{2}m_{\rm eff}^2\varphi_0^2,$$
$$\alpha_2(\varphi_0) = \frac{2V_0}{\sqrt{2}v_P\mu^2(M_{\rm Pl}^2 - v_P^2)^{1/2}\varphi_0/M_{\rm Pl} - 2\beta v_P^2 M_{\rm Pl}^2} + \mathcal{O}(\varphi_0^0). \tag{101}$$

These expressions mark one of the main results in this paper: Provided that the parameter α in the superpotential W^P in Eq. (83) takes either of these two values, the scalar field p^- becomes unstable, once the inflaton field φ reaches the value φ_0 at the end of inflation or after inflation. This demonstrates how varying the value of the parameter α puts us in the position to control the time at which R symmetry is spontaneously broken via gaugino condensation at the end of inflation. For small values of β and v_P (such that $\beta \ll \beta_0$ and $v_P \ll M_{\rm Pl}$), the solutions for α in Eq. (101) reduce to more compact expressions,

$$\alpha_1(\varphi_0) \approx \frac{V_{\rm eff}^{1/2}(\varphi_0)}{v_P M_{\rm Pl}}, \qquad \alpha_2(\varphi_0) \approx \frac{\sqrt{2}\mu^2}{v_P\varphi_0}. \tag{102}$$

As evident from Eq. (101), the first solution, $\alpha = \alpha_1(\varphi_0)$, depends only weakly on the field value φ_0. In fact, for given values of all other free parameters (μ, λ, v_P, and β), matching α with $\alpha_1(\varphi_0)$ for a reasonable value of φ_0 requires some amount of fine-tuning. To first approximation, all values $\alpha_1(\varphi_0)$ correspond to the reference value α_0, their variation with φ_0 being a subdominant effect,

$$\alpha_1(\varphi_0) \approx \alpha_0 \simeq \frac{\mu^2}{v_P M_{\rm Pl}} \simeq 4\times 10^{-6}\left(\frac{\mu}{10^{15}\ \mathrm{GeV}}\right)^2\left(\frac{10^{17}\ \mathrm{GeV}}{v_P}\right). \tag{103}$$

In the context of our model, this is, however, not a problem. Recall that the ultimate purpose of the superpotential W^P in Eq. (83) is to help us in addressing the CC problem. From that point of view, it is expected that some (if not all) of the parameters involved in determining the final value of the CC end up being subject to some amount of fine-tuning. In addition, we emphasize that the value of α may have an important impact on the time when inflation ends. In regions of space (or the "string landscape") where α takes too small a value, gaugino condensation never occurs, i.e., inflation never ends. On the other hand, in regions of space where α takes too large a value, the constant in the superpotential, $W \supset w$, is already generated at early times. Inflation then only lasts for a few e-folds or does not take place at all. Universes in which α matches $\alpha_1(\varphi_0)$ for a reasonable value of φ_0, i.e., universes in which inflation lasts for a long but finite time interval, therefore, eventually occupy the largest spatial volume among all postinflationary universes. This might explain why, in habitable universes, α appears to be fine-tuned.

Finally, we evaluate the total p^- mass $m_{p^-}^2$ in Eq. (87) for the two solutions for α in Eq. (101),

$$\alpha \to \alpha_1 \Rightarrow m_{p^-}^2(\varphi) = \left[\frac{1}{2}m_{\rm eff}^2 - \frac{3b}{4(1+b)^2}\frac{V_0}{M_{\rm Pl}^2} + \mathcal{O}(M_{\rm Pl}^{-4})\right] \times \frac{\varphi^2-\varphi_0^2}{M_{\rm Pl}^2} + \mathcal{O}(\varphi^4),$$
$$\alpha \to \alpha_2 \Rightarrow m_{p^-}^2(\varphi) = \left[2(1+b)^2\frac{V_0}{\varphi_0^2} + \mathcal{O}(M_{\rm Pl}^{-2})\right] \times \frac{\varphi^2-\varphi_0^2}{\varphi_0^2} + \mathcal{O}(\varphi^4), \qquad (104)$$

which nicely illustrates how the scalar DOF p^- turns tachyonic, once the inflaton field φ reaches the value φ_0. In this sense, Polonyi inflation does end in a "waterfall transition," after all. But, instead of taking place in the inflaton sector itself, this transition occurs in separate hidden sector. The inflaton and waterfall sectors are, therefore, separated in our model, which allows us to bring together two phenomena that would otherwise mutually exclude each other: (i) Inflation in a scalar potential equivalent to that of ordinary F-term hybrid inflation and (ii) supersymmetry breaking at a very high scale. For a discussion of the tension between these two phenomena and a possible resolution in the context of string theory, see [93]. We note that, in contrast to [93], our scenario represents a fully field-theoretic construction.

In closing, we also mention that, adding the particular Kähler potential in Eq. (89), the expressions for m_{p^-} in Eq. (104) remain quantitatively the same. The only effect of $K_{\rm mix}$ is that it increases the overall magnitude of the terms in square brackets in Eq. (104). We have explicitly checked numerically that, if we choose $\epsilon_{P,Y} \sim -1$, $M_* \sim 10^{17}$ GeV, and $M_*{}' \sim M_{\rm Pl}$ in Eq. (89), our first solution for m_{p^-} exceeds the inflationary Hubble rate for all times during inflation, $m_{p^-} \gg H_{\rm inf}$. At the same time, for field values slightly below $\varphi \simeq \varphi_0$, the absolute value of m_{p^-} is again much larger than the Hubble rate, $m_{p^-} \ll -H_{\rm inf}$. This ensures that the scalar field p^- quickly reaches the minimum of the scalar potential. Meanwhile, all other scalars always have masses that are large compared to $H_{\rm inf}$, so that they are safely stabilized at the origin. A possible interpretation of these numerical findings is to suppose the existence of an other strongly coupled hidden gauge sector around the scale $M_* \sim 10^{17}$ GeV. If the singlet P is the only field that is strongly coupled to this new dynamical sector (meaning that Y and Φ are only weakly coupled), simple arguments from NDA automatically result in a Kähler potential of just the desired form.[18] In particular, such new strong dynamics do not induce a dangerous $|\Phi|^4/M_*^2$ term in the Kähler potential (which would otherwise spoil slow-roll inflation), as long as the Polonyi field is only weakly coupled to the new sector. However, we emphasize that this picture of a new gauge sector around $M_* \sim 10^{17}$ GeV is just a simple example to illustrate what additional dynamics could possibly lead to a cutoff scale of $\mathcal{O}(10^{17})$ GeV in the $|\Phi|^2|P|^2$ operator in the noncanonical Kähler potential $K_{\rm mix}$. From the perspective of our model, the main consequence of $K_{\rm mix}$ in Eq. (89) is that it provides us with the ability to increase the absolute value as well as the gradient of m_{p^-} as a function of φ. Apart from that, our simplified analysis based on $K_{\rm mix} = 0$ captures all relevant aspects of our mechanism of R symmetry breaking. For this reason, we leave a more detailed study of $K_{\rm mix}$ and its UV origin for future work.

D. Backreaction on inflation and low-energy ground state after inflation

The "waterfall superpotential" W^P in Eq. (83) not only triggers a mass deformation in the R-symmetry-breaking sector, it also results in Planck-suppressed corrections to the inflaton potential in Eq. (69). During inflation, i.e., as long as $w = 0$, the scalar Polonyi potential now reads

$$V(\varphi) = C_0 + \frac{C_2}{2}\varphi^2 + \frac{C_4}{4!}\varphi^4 + \mathcal{O}(\varphi^{-6}), \qquad (105)$$

where the coefficients C_0, C_2, and C_4 correspond to the coefficients c_0, c_2, and c_4 in Eq. (67) evaluated at $w = 0$ and multiplied by correction factors that stem from the R-symmetry-breaking sector,

$$C_0 = (1+\nu)c_0|_{w=0}, \qquad C_2 = \left(1 - \frac{b}{1+b}\frac{\nu}{\epsilon}\right)c_2|_{w=0},$$
$$C_4 = \left[1 + \left(3 + \frac{32b}{3} + \frac{79b^2}{6} + 6b^3 + b^4 + \mathcal{O}(\epsilon,\nu)\right) \times \frac{\nu}{(1+b)^4}\right] c_4|_{w=0}. \qquad (106)$$

Here, we have introduced the parameter ν to measure the relative importance of these new corrections,

[18]For $\epsilon_P < 0$, the kinetic term for the field P becomes singular at $\varphi \simeq M_*$, as long as we only assume the Kähler potential in Eq. (89). For this reason, we expect additional higher-dimensional operators in F_P to become important around $\varphi \simeq M_*$, which regulate the kinetic term for P [see Eq. (88)]. Here, one simple solution is to assume that F_P as a function of φ simply saturates at a maximal negative value, $F_P \to -F_0 > -1$, around field values of $\mathcal{O}(M_*)$, so that the prefactor of the kinetic term for P always remains positive, $K \supset (1+F_P)P^\dagger P \to (1-F_0)P^\dagger P$. At large field values, $\varphi \gtrsim M_*$, the field P is then stabilized by the field-dependent mass in Eq. (96), while for $\varphi_0 \lesssim \varphi \lesssim M_*$, it is stabilized by the terms in Eq. (90).

$$\nu = \frac{\Delta V_P}{V_0} = \frac{\alpha^2 v_P^4}{\mu^4}, \tag{107}$$

which is nothing but the ratio of the Y and Φ F-terms during inflation, $\nu = |F_Y/F_\Phi|$. To make sure that the R-symmetry-breaking sector does not disturb the inflationary dynamics, we have to require that ν is sufficiently small. In view of the expression for C_2 in Eq. (106), we have to demand in particular that

$$\nu \ll \epsilon, \tag{108}$$

since otherwise the inflaton field will receive too large a Hubble-induced mass, $m_\varphi^2 \sim 3\nu H_{\rm inf}^2$. This constraint then translates into upper bounds on the mass scale v_P in the R-symmetry-breaking sector,

$$\begin{aligned}
\alpha = \alpha_1(\varphi_0) &\Rightarrow v_P \ll v_P^{\max,1} = \epsilon^{1/2} M_{\rm Pl}, \\
\alpha = \alpha_2(\varphi_0) &\Rightarrow v_P \ll v_P^{\max,2} = \frac{\epsilon^{1/2}\varphi_0}{\sqrt{2}(1+b)}.
\end{aligned} \tag{109}$$

The scale v_P is a free parameter, which suggests that we should actually always be able to satisfy these bounds. At the end of this section, we will, however, see that the upper bound $v_P^{\max,2}$ turns out to be too restrictive, so that only our first solution, $\alpha = \alpha_1(\varphi_0)$, remains phenomenologically viable.

To get there, we first need to examine the true vacuum after the end of inflation as well as the tuning of the CC in more detail. In global supersymmetry, the VEVs of the fields Φ, Y, and P are found to be

$$\text{Global SUSY}: \langle\Phi\rangle = 0, \quad \langle Y\rangle = 0, \quad \langle P\rangle = \sqrt{2} v_P. \tag{110}$$

The fluctuations around the true vacuum are, therefore, described by the following superpotential:

$$W = \mu^2\Phi - \sqrt{2}\alpha v_P Y P' - \frac{\alpha}{2} Y P'^2 + \frac{\beta}{3!} Y^3 + w, \tag{111}$$

where P' denotes the singlet field P shifted by its VEV in global supersymmetry, $P \to \sqrt{2} v_P + P'$. Equation (111) tells us that, from the perspective of global supersymmetry, the F-term of the field Y completely vanishes in the true vacuum, so that the Polonyi field Φ remains as the only SUSY-breaking field. In addition, we find that the fields Y and P now share a common supersymmetric Dirac mass. In global supersymmetry, the scalar mass eigenvalues at low energies are consequently given as follows:

$$\text{Global SUSY}: m_\phi = m_{\rm eff}, \quad m_y = m_p = \sqrt{2}\alpha v_P, \tag{112}$$

where ϕ, y, and p stand for the complex scalars contained in the chiral fields Φ, Y, and P', respectively.

In order to find the SUGRA corrections to the VEVs in Eq. (110) and the masses in Eq. (112), we expand the full SUGRA scalar potential up to second order in the fluctuations ϕ, y, and p,

$$\begin{aligned}
V(\phi, y, p) = c_0 &+ c_\phi(\phi + \phi^*) + c_y(y + y^*) \\
&+ c_p(p + p^*) + m_\phi^2|\phi|^2 + m_y^2|y|^2 + m_p^2|p|^2 \\
&+ m_{\phi y}^2(\phi y^* + \phi^* y) + m_{\phi p}^2(\phi p^* \\
&+ \phi^* p) + m_{yp}^2(y p^* + y^* p) + \ldots.
\end{aligned} \tag{113}$$

Here, the coefficients of the linear terms (c_ϕ, c_y, c_p), the masses around the true vacuum (m_ϕ, m_y, m_p), and the mass mixing parameters ($m_{\phi y}$, $m_{\phi p}$, m_{yp}) take, to leading order, the following form:

$$\begin{aligned}
&c_\phi \simeq -2\frac{\mu^2 w}{M_{\rm Pl}^2}, \quad c_y \simeq -2\frac{\alpha v_P^2 w}{M_{\rm Pl}^2}, \\
&c_p \simeq \sqrt{2}\frac{v_P V_0}{M_{\rm Pl}^2} - 2\sqrt{2}\frac{v_P w^2}{M_{\rm Pl}^4}, \quad m_\phi^2 \simeq m_{\rm eff}^2 - \epsilon\frac{V_0}{M_{\rm Pl}^2} - 2\frac{w^2}{M_{\rm Pl}^4}, \\
&m_y^2 \simeq 2\alpha^2 v_P^2 + \frac{V_0}{M_{\rm Pl}^2} - 2\frac{w^2}{M_{\rm Pl}^4}, \quad m_p^2 \simeq 2\alpha^2 v_P^2 + \frac{V_0}{M_{\rm Pl}^2} - 2\frac{w^2}{M_{\rm Pl}^4}, \\
&m_{\phi y}^2 \simeq -2\frac{\alpha v_P^2 \mu^2}{M_{\rm Pl}^2}, \quad m_{\phi p}^2 \simeq -2\sqrt{2}\frac{v_P \mu^2 w}{M_{\rm Pl}^4}, \\
&m_{yp}^2 \simeq -6\sqrt{2}\frac{\alpha v_P^3 w}{M_{\rm Pl}^4}.
\end{aligned} \tag{114}$$

Given that $\nu \ll 1$ [see Eqs. (107) and (108)], the mass mixing among ϕ, y, and p turns out be negligible. Setting the off-diagonal masses $m_{\phi y}$, $m_{\phi p}$, and m_{yp} in Eq. (113) to zero, we then arrive at the following simple expressions for the VEVs of the fields Φ, Y, and P in supergravity:

$$\langle\Phi\rangle \simeq -\frac{c_\phi}{m_\phi^2}, \quad \langle Y\rangle \simeq -\frac{c_y}{m_y^2}, \quad \langle P\rangle \simeq \sqrt{2} v_P - \frac{c_p}{m_p^2}, \tag{115}$$

which results in the following vacuum energy density:

$$\begin{aligned}
&\langle V\rangle \simeq c_0 + c_\phi\langle\Phi\rangle + c_y\langle Y\rangle + c_p(\langle P\rangle - \sqrt{2}v_P), \\
&c_0 = V_0 - 3\frac{w^2}{M_{\rm Pl}^2},
\end{aligned} \tag{116}$$

where the inflaton term, $c_\phi\langle\Phi\rangle$, clearly constitutes the largest correction to the leading c_0 term. In order to make the CC $\langle V\rangle$ vanish, we have to fine-tune the constant w, so that it takes the following value:

$$w_0 \simeq \frac{1}{\sqrt{3}}(1-\gamma)\mu^2 M_{\rm Pl}, \quad \gamma = \frac{2m_{3/2}'^2}{m_{\rm eff}^2 - \epsilon V_0/M_{\rm Pl}^2 + 2m_{3/2}'^2},$$
$$m_{3/2}' = \frac{\mu^2}{\sqrt{3}M_{\rm Pl}^2}. \tag{117}$$

Here, $m_{3/2}'$ denotes the "asymptotic gravitino mass" in the limit $\langle\Phi\rangle \to 0$. The physical gravitino mass, on the other hand, follows from plugging our results for w_0 as well as for the field VEVs into Eq. (72),

$$m_{3/2} \simeq \exp\left[\frac{\langle\Phi\rangle^2}{2M_{\rm Pl}^2}\right]\left(\frac{1-\gamma}{\sqrt{3}} + \frac{\langle\Phi\rangle}{M_{\rm Pl}}\right)\frac{\mu^2}{M_{\rm Pl}},$$
$$\langle\Phi\rangle \simeq \frac{\sqrt{3}\gamma}{1-\gamma+\gamma^2}M_{\rm Pl} \approx \frac{4}{3}\sqrt{3}\gamma M_{\rm Pl}. \tag{118}$$

As evident from Eq. (118), the parameter γ controls the size of the Polonyi VEV in the true vacuum. In fact, it is a convenient measure for the relation between the individual contributions to the "asymptotic inflaton mass" m_ϕ' in the limit $\langle\Phi\rangle \to 0$. In that limit, we have $m_{3/2} \to m_{3/2}'$, $\gamma \to 0$, and $m_\phi \to m_\phi'$, with the asymptotic inflaton mass m_ϕ' being given as follows:

$$m_\phi'^2 = m_{\phi,0}^2 + m_{\phi,w}^2, \qquad m_{\phi,0}^2 = m_{\rm eff}^2 - \epsilon\frac{V_0}{M_{\rm Pl}^2},$$
$$m_{\phi,w}^2 = -2m_{3/2}'^2. \tag{119}$$

Here, $m_{\phi,0}$ denotes the effective inflaton mass close to the origin for $w=0$, whereas $m_{\phi,w}$ represents the correction to m_ϕ' in the true vacuum, i.e., the correction appearing after the generation of the constant w. Making use of these definitions, we recognize that γ parametrizes nothing else than the following ratio:

$$\gamma = \frac{|m_{\phi,w}|^2}{|m_{\phi,0}|^2 + |m_{\phi,w}|^2}. \tag{120}$$

In order to stabilize the Polonyi field at a sub-Planckian field value after the end of inflation, $\langle\Phi\rangle \ll M_{\rm Pl}$, we have to require that the additional SUGRA correction generated in the course of late-time R symmetry breaking, $m_{\phi,w}$, always remains smaller than the effective one-loop mass $m_{\rm eff}$,

$$\langle\Phi\rangle \ll M_{\rm Pl} \Leftrightarrow \gamma \ll 1 \Leftrightarrow m_{3/2}' \ll m_{\rm eff}. \tag{121}$$

We remark that this requirement is a consequence of the fact that $m_{\rm eff}$ and $m_{3/2}'$ are controlled by the same dynamics, i.e., by the dynamical scale Λ and the Yukawa coupling λ. These two masses are, therefore, potentially of the same order of magnitude, in which case the SUGRA correction $m_{\phi,w}$ threatens to destabilize the minimum of the effective potential in global supersymmetry. For this reason, we have to explicitly impose the requirement $m_{3/2}' \ll m_{\rm eff}$ as an extra condition. Once this condition is satisfied, our results in Eqs. (114), (117), and (118) simplify considerably. For small values of γ, we find

$$w_0 \simeq \frac{\mu^2 M_{\rm Pl}}{\sqrt{3}}, \qquad m_{3/2} \simeq m_{3/2}', \qquad m_\phi^2 \simeq m_{\rm eff}^2 - 2m_{3/2}^2,$$
$$m_y^2 \simeq 2\alpha^2 v_P^2 + m_{3/2}^2, \qquad m_p^2 \simeq m_y^2. \tag{122}$$

In combination with Eqs. (114) and (115), these results lead to simple expressions for $\langle\Phi\rangle$, $\langle Y\rangle$, and $\langle P\rangle$ that are valid for $\gamma \ll 1$. Neglecting all effects of $\mathcal{O}(\gamma)$ and $\mathcal{O}(\epsilon)$ wherever possible, we obtain

$$\langle\Phi\rangle \simeq \frac{2\sqrt{3}m_{3/2}^2}{m_{\rm eff}^2 - 2m_{3/2}^2}M_{\rm Pl}, \qquad \langle Y\rangle \simeq \frac{2\alpha v_P m_{3/2}}{2\alpha^2 v_P^2 + m_{3/2}^2}v_P,$$
$$\langle P\rangle \simeq \left(1 - \frac{m_{3/2}^2}{2\alpha^2 v_P^2 + m_{3/2}^2}\right)\sqrt{2}v_P. \tag{123}$$

These expressions are the SUGRA counterpart to the VEVs in Eqs. (110) and represent our final results for $\langle\Phi\rangle$, $\langle Y\rangle$, and $\langle P\rangle$. In this vacuum, also the fields Y and P have nonzero F-terms. These F-terms are suppressed compared to the Polonyi F-term by a factor of $\mathcal{O}(v_P/M_{\rm Pl})$ and are, hence, negligible.

The masses $m_{3/2}'$ and $m_{\rm eff}$ scale as follows with the Yukawa coupling λ and the dynamical scale Λ,

$$m_{\rm eff} \simeq 4\times 10^{10}\ \mathrm{GeV}\left(\frac{\lambda}{0.2}\right)^3\left(\frac{\Lambda}{10^{16}\ \mathrm{GeV}}\right),$$
$$m_{3/2}' \simeq 4\times 10^{10}\ \mathrm{GeV}\left(\frac{\lambda}{0.2}\right)\left(\frac{\Lambda}{10^{16}\ \mathrm{GeV}}\right)^2. \tag{124}$$

For a fixed value of Λ, the requirement that $m_{\rm eff}$ must exceed $m_{3/2}'$ then results in a lower bound on λ,

$$m_{3/2}' \lesssim m_{\rm eff} \Rightarrow \lambda \gtrsim 0.2\left(\frac{\Lambda}{10^{16}\ \mathrm{GeV}}\right)^{1/2}. \tag{125}$$

But this is not the end of the story. We must also make sure that the final Polonyi VEV, $\langle\varphi\rangle$, lies below the critical field value φ_c, so as to stay in the quadratic part of the effective potential close to the origin,

$$\langle\varphi\rangle \lesssim \varphi_c. \tag{126}$$

Otherwise, the mass term in the effective potential disappears altogether and the Polonyi field rolls back to field values of the order of the Planck scale. This requirement yields an even stronger bound on λ,

$$\langle\varphi\rangle \lesssim \varphi_c \Rightarrow \lambda \gtrsim 1.0\left(\frac{\Lambda}{10^{16}\ \mathrm{GeV}}\right)^{1/3}. \qquad (127)$$

Together with the requirement of perturbativity, this limits the range of viable λ values to $1 \lesssim \lambda \lesssim 4$.

Having determined the position of the true vacuum after inflation in field space, we are now finally in the position to link the R-symmetry-breaking sector to the IYIT sector. In order to tune the CC in the true vacuum to zero, the constant w generated in the R-symmetry-breaking sector [see Eq. (80)] needs to be matched with the value w_0 dictated by the IYIT sector [see Eq. (122)]. For $N_c = N_f = 2$, we have

$$w \simeq m_{3/2} M_{\mathrm{Pl}}^2 \simeq \frac{1}{8\pi^2}\langle P\rangle\tilde{\Lambda}^2, \qquad w_0 \simeq \frac{\mu^2 M_{\mathrm{Pl}}}{\sqrt{3}}, \qquad (128)$$

which results in a condition on the dynamical scale $\tilde{\Lambda}$ in the R-symmetry-breaking sector,

$$w = w_0 \Rightarrow \tilde{\Lambda} \simeq \left(\frac{8\pi^2}{\sqrt{3}}\frac{\mu^2 M_{\mathrm{Pl}}}{\langle P\rangle}\right)^{1/2}. \qquad (129)$$

This solution for $\tilde{\Lambda}$ is only consistent as long as it is smaller than the heavy-quark mass scale $\bar{M}_Q = \langle P\rangle$, i.e., as long as the gaugino condensation scale is smaller than heavy-quark mass scale, $w \lesssim w_{\max}$ [see Eq. (81)]. This requirement can be reformulated as an upper bound on the SUSY breaking scale μ,

$$\tilde{\Lambda} \lesssim \langle P\rangle \Rightarrow \mu \lesssim \left(\frac{\sqrt{3}}{8\pi^2}\frac{\langle P\rangle^3}{M_{\mathrm{Pl}}}\right)^{1/2}. \qquad (130)$$

In supergravity, $\langle P\rangle$ follows from Eq. (123). Depending on our choice for α, we find two solutions,

$$\langle P\rangle \simeq \left(1 - \frac{1}{1+6a}\right)\sqrt{2}v_P \simeq \sqrt{2}v_P \times \begin{cases} 6/7 & ; a = a_1(\varphi_0) \simeq 1 \\ 1 & ; a = a_2(\varphi_0) \gg 1 \end{cases}. \qquad (131)$$

Together with the limits on the scale v_P in Eq. (109), these relations imply upper bounds on Λ,

$$\alpha = \alpha_1(\varphi_0) \Rightarrow \Lambda \lesssim 2\times 10^{18}\ \mathrm{GeV}\left(\frac{\epsilon}{0.2}\right)^{3/4}\left(\frac{1}{\lambda}\right)^{1/2}\left(\frac{v_P}{v_P^{\max,1}}\right)^{3/2},$$

$$\alpha = \alpha_2(\varphi_0) \Rightarrow \Lambda \lesssim 3\times 10^{14}\ \mathrm{GeV}\left(\frac{\epsilon}{0.2}\right)^{3/4}\left(\frac{1}{\lambda}\right)^{1/2}\left(\frac{v_P}{v_P^{\max,2}}\right)^{3/2} \times\left(\frac{1}{1+b}\right)^{3/2}\left(\frac{\varphi_0}{10^{16}\ \mathrm{GeV}}\right)^{3/2}. \qquad (132)$$

Once this condition is satisfied, there always exists a sufficiently small value of $\tilde{\Lambda}$ that allows us to tune the CC to zero (so that $w = w_0$, while at the same time $\tilde{\Lambda} \lesssim \langle P\rangle$). We, however, note that the upper bound on Λ in the case of the second solution is too severe. For such a small dynamical scale, the vacuum energy density during inflation does not suffice to account for the scalar spectral amplitude A_s. The requirement of successful Polonyi inflation, therefore, singles out the fine-tuned first solution $\alpha = \alpha_1(\varphi_0) \approx \alpha_0$.

IV. PHENOMENOLOGICAL IMPLICATIONS

A. Properties of the scalar potential driving inflation

We now have everything at our disposal that we need to study the phenomenological implications of our model. As argued in Sec. II D, the scalar potential driving inflation takes the following form [see Eq. (69)]:

$$V(\varphi) = V_0 + \frac{1}{2}m_\varphi^2\varphi^2 + \frac{\lambda_\varphi}{4!}\varphi^4 + \Lambda_{\mathrm{HE}}^4 \ln\frac{\varphi}{\varphi_c}, \qquad (133)$$

where V_0 and Λ_{HE}^4 are given in Eqs. (8) and (44), respectively, and where m_φ^2 and λ_φ follow from Eq. (67),

$$V_0 = \mu^4, \qquad m_\varphi^2 = -\epsilon\frac{V_0}{M_{\mathrm{Pl}}^2} = -3\epsilon H_{\mathrm{inf}},$$

$$\lambda_\varphi = 3\left(1 - \frac{7\epsilon}{2} + \frac{8\epsilon}{3}\right)\frac{V_0}{M_{\mathrm{Pl}}^4}, \qquad \Lambda_{\mathrm{HE}}^4 = N_X\frac{m^4}{16\pi^2}. \qquad (134)$$

This scalar potential is identical to the inflaton potential of F-term hybrid inflation (including corrections from supergravity and a noncanonical Kähler potential) in the limit of a vanishing gravitino mass. Its implications for the inflationary CMB observables have been studied for the first time in [64]. In the following, we shall review the aspects of this inflationary scenario that are most relevant for our purposes and summarize the predictions for the inflationary CMB observables in our notation.

First of all, we note that the potential in Eq. (133) has an inflection point at the following field value:

$$\varphi_{\text{flex}} = \frac{1}{\lambda_\varphi^{1/2}}[-m_\varphi^2 + (m_\varphi^4 + 2\lambda_\varphi \Lambda_{\text{HE}}^4)^{1/2}]^{1/2},$$
$$V''(\varphi_{\text{flex}}) = 0. \tag{135}$$

A priori, the sign of the potential gradient at $\varphi = \varphi_{\text{flex}}$ is undetermined. In particular, the inflection point at $\varphi = \varphi_{\text{flex}}$ turns into a saddle point, $V'(\varphi_{\text{flex}}) = 0$, once the following relation is satisfied:

$$3m_\varphi^4 - 2\lambda_\varphi \Lambda_{\text{HE}}^4 = 0. \tag{136}$$

This condition can be fulfilled by setting the coefficient ϵ to a particular critical value ϵ_0,

$$\epsilon_0 = \frac{12\rho}{21\rho + (57\rho^2 + 72\rho)^{1/2}} = \frac{12\lambda}{21\lambda + (45\lambda^2 + 384\pi^2)^{1/2}},$$
$$\rho = \frac{\Lambda_{\text{HE}}^4}{V_0}, \tag{137}$$

which solely depends on the size of the Yukawa coupling λ. For λ values of $\mathcal{O}(1)$, the critical coefficient ϵ_0 takes values of $\mathcal{O}(0.1)$. Once we have $\epsilon > \epsilon_0$, the potential gradient at $\varphi = \varphi_{\text{flex}}$ is negative, $V'(\varphi_{\text{flex}}) < 0$, whereas, for $\epsilon < \epsilon_0$, the potential gradient at $\varphi = \varphi_{\text{flex}}$ is positive, $V'(\varphi_{\text{flex}}) > 0$. For $\epsilon > \epsilon_0$, the scalar potential, therefore, exhibits a local minimum (+) as well as a local maximum (−) in the vicinity of φ_{flex},

$$\epsilon \geq \epsilon_0 \Rightarrow \varphi_{\text{min,max}} = \frac{\sqrt{3}}{\lambda_\varphi^{1/2}}\left[-m_\varphi^2 \pm \frac{1}{\sqrt{3}}(3m_\varphi^4 - 2\lambda_\varphi \Lambda_{\text{HE}}^4)^{1/2}\right]^{1/2},$$
$$\varphi_{\text{max}} \leq \varphi_{\text{flex}} \leq \varphi_{\text{min}}. \tag{138}$$

Note that these two field values become identical, once the condition in Eq. (136) is satisfied,

$$\epsilon = \epsilon_0 \Rightarrow \varphi_{\text{min}} = \varphi_{\text{flex}} = \varphi_{\text{max}} = \left(\frac{-3m_\varphi^2}{\lambda_\varphi}\right)^{1/2} \simeq \epsilon_0^{1/2} M_{\text{Pl}}. \tag{139}$$

The presence of two local extrema in the scalar potential may be regarded as a disadvantage, as it requires the initial inflaton field value, φ_{ini}, to be smaller than φ_{max}. For one thing, this requirement may necessitate some amount of fine-tuning, $\varphi_{\text{ini}} \simeq \varphi_{\text{max}}$, so as to achieve a sufficient number of e-folds of inflation. For another thing, the constraint $\varphi_{\text{ini}} \leq \varphi_{\text{max}}$ is not compatible with the notion that inflation is expected to descend from a "Planck epoch," during which the inflaton field, the potential energy density, and all other relevant quantities take values of the order of the Planck scale. For these reasons, we consider the option $\epsilon > \epsilon_0$ less likely than the alternative possibility $\epsilon < \epsilon_0$. We can take this argument even one step further by making the explicit assumption that, for one reason or another, φ_{ini} is necessarily larger than the Planck scale, $\varphi_{\text{ini}} \gtrsim M_{\text{Pl}}$. Under

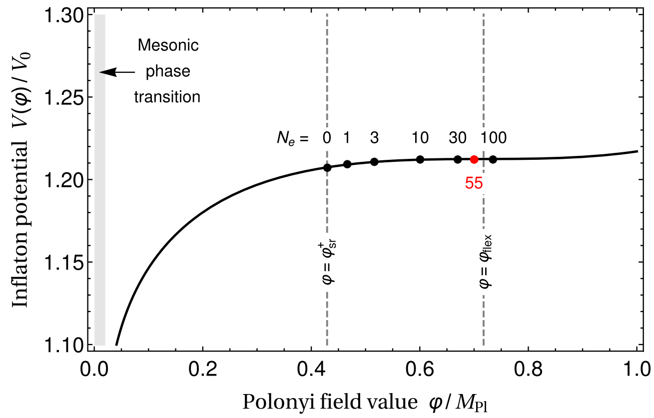

FIG. 2. Inflaton potential in the inflection-point regime ($\epsilon < \epsilon_0$). Here, the values for the free parameters in the inflaton sector have been chosen as in the benchmark scenario discussed in Sec. IV D: $\Lambda \simeq 1.26 \times 10^{16}$ GeV, $\lambda \simeq 1.66$, and $\epsilon \simeq 0.204$.

this assumption, ϵ *must* be smaller than ϵ_0, as the inflaton would otherwise get trapped in the false vacuum at $\varphi = \varphi_{\text{min}}$. From this perspective, the critical coefficient ϵ_0 plays the role of an upper bound on ϵ, which decides whether the inflaton field has a chance of reaching the true vacuum or not. In Fig. 2, we plot the scalar potential in Eq. (133) for an ϵ value that is slightly smaller than the critical value ϵ_0. In this case, the scalar potential features a flat plateau around $\varphi \sim \varphi_{\text{flex}}$, which gives rise to successful inflation in accord with the observational data.

Last but not least, we note that the scalar potential in Eq. (133) only allows for slow-roll inflation within a limited range of field values. The slow-roll conditions, $\varepsilon \ll 1$ and $\eta \ll 1$ [see Eq. (9)], are only satisfied in the interval $\varphi_{\text{sr}}^+ \leq \varphi \leq \varphi_{\text{sr}}^-$, where the two boundary values $\varphi_{\text{sr}}^\pm$ are given as follows[19]:

$$\varphi_{\text{sr}}^\pm = \frac{1}{\lambda_\varphi^{1/2}}[-M_\pm^2 + (M_\pm^4 + 2\lambda_\varphi \Lambda_{\text{HE}}^4)^{1/2}]^{1/2},$$
$$M_\pm = \left(m_\varphi^2 \pm |\eta_{\text{max}}| \frac{V(\varphi_{\text{sr}}^\pm)}{M_{\text{Pl}}^2}\right)^{1/2}. \tag{140}$$

Here, $|\eta_{\text{max}}|$ denotes the maximally allowed absolute value of the slow-roll parameter η. The slow-roll parameter ε is, by comparison, always subdominant during inflation, $\varepsilon \ll |\eta|$. For definiteness, we will use $|\eta_{\text{max}}| = 10^{-0.5}$ in the following. Finally, we also mention that, if we formally take the limit $|\eta_{\text{max}}| \to 0$ in Eq. (140), the expression for $\varphi_{\text{sr}}^\pm$ in Eq. (140) reduces to our result for φ_{flex} in Eq. (135).

For the parameter region of interest, $1 \lesssim \lambda \lesssim 4$ [see Eq. (127)], the field value φ_{sr}^+ is always larger than the critical value field φ_c associated with the quark-meson

[19]In principle, Eq. (140) represents an implicit definition of $\varphi_{\text{sr}}^\pm$, as $M_\pm$ depends on $V(\varphi_{\text{sr}}^\pm)$. In practice, this dependence is, however, very weak, $V(\varphi_{\text{sr}}^\pm) \approx V_0$, so that the right-hand side of Eq. (140) readily yields the desired values of $\varphi_{\text{sr}}^\pm$.

phase transition in the IYIT model. Inflation, therefore, always ends, once φ reaches $\varphi_{\rm sr}^+$. That is, inflation ends because the slow-roll conditions become violated, and not because of a sudden phase transition (such as the one in the R-symmetry-breaking sector) triggered by a small inflaton field value. Correspondingly, the number of e-folds of slow-roll inflation, N_e, in dependence of the inflaton field value is always given by the following integral:

$$N_e(\varphi) = \int_{\varphi_{\rm sr}^+}^{\varphi} \frac{d\varphi'}{M_{\rm Pl}} \frac{1}{\sqrt{2\varepsilon(\varphi')}}, \qquad (141)$$

where N_e is defined such that it counts the remaining number of e-folds until the end of inflation. In the following, we will work in the approximation that the number of e-folds in between the end of slow-roll inflation and the onset of reheating is negligible. This is to say that we assume all scalar fields to settle at their respective VEVs sufficiently fast, once slow-roll inflation has ended. Solving the integral in Eq. (141) and inverting the relation $N_e = N_e(\varphi)$ then provides us with the inflaton field value as a function of the number of e-folds, $\varphi = \varphi(N_e)$. For the purposes of this paper, we perform these steps numerically.

B. Inflationary CMB observables

The function $\varphi(N_e)$ is exactly what we need to determine our predictions for the CMB observables (the scalar spectral amplitude A_s, the scalar spectral index n_s, as well as the tensor-to-scalar ratio r),

$$A_s = \frac{1}{24\pi^2\varepsilon_*}\frac{V_*}{M_{\rm Pl}^4}, \quad n_s = 1 + 2\eta_* - 6\varepsilon_*, \quad r = 16\varepsilon_*. \qquad (142)$$

Here, the asterisk indicates that all quantities are to be evaluated at $\varphi_* = \varphi(N_e^*)$, where $N_e^* \simeq 55$ denotes the required number of e-folds during slow-roll inflation. According to the latest results of the Planck Collaboration, the 95% confidence intervals for these observables are given as follows [58]:

$$A_s^{\rm obs} = e^{3.062\pm0.029} \times 10^{-10} \simeq 2 \times 10^{-9},$$
$$n_s^{\rm obs} = 0.9677 \pm 0.0060, \qquad r < 0.11. \qquad (143)$$

The scalar potential in Eq. (133) is basically controlled by three free parameters: the dynamical scale Λ, the Yukawa coupling λ, and the coefficient in the higher-dimensional Kähler potential, ϵ. This parametric freedom always allows us to choose Λ, such that we succeed in reproducing the correct normalization of the scalar power spectrum, $A_s \simeq A_s^{\rm obs}$. After fixing Λ in this way, we then have to deal with two free parameters—λ and ϵ (see also the discussion at the end of Sec. I C)—which leaves us with the task of studying the predictions for the inflationary CMB observables as functions of λ and ϵ.

The outcome of our analysis is shown in Fig. 3. Both panels of Fig. 3 indicate where in parameter space our prediction for n_s falls into the observed 2σ range.

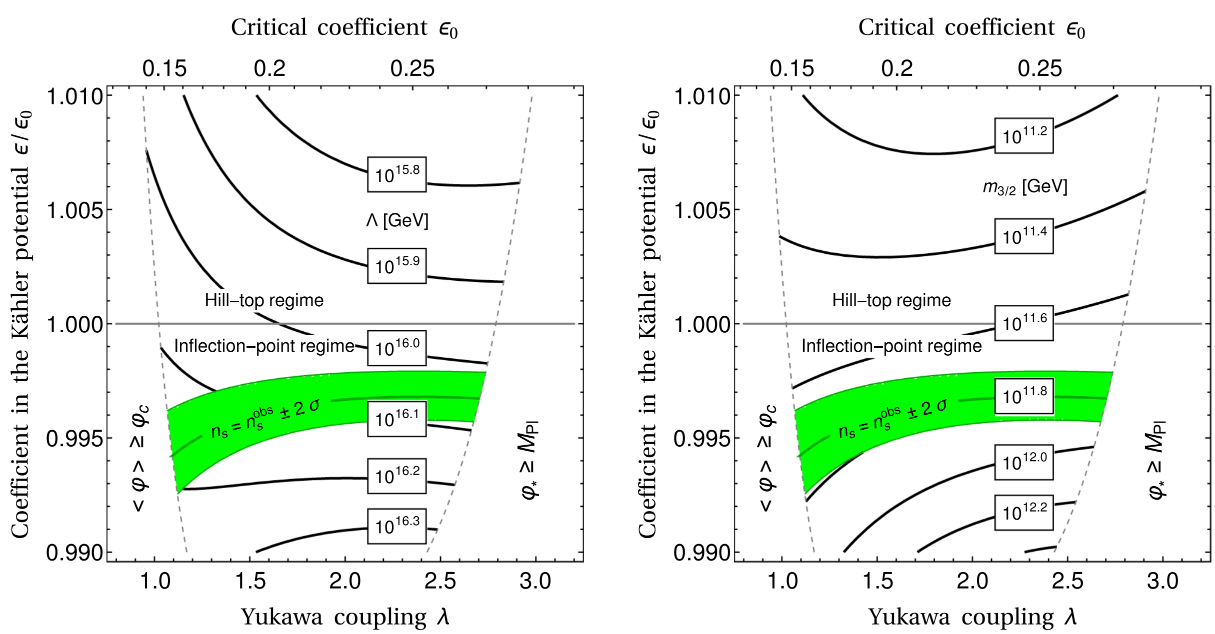

FIG. 3. Viable region in the parameter space spanned by the Yukawa coupling λ and the coefficient in the noncanonical Kähler potential, ϵ. Here, the ϵ values on the vertical axis are normalized by the critical values ϵ_0, see Eq. (137). For $\lambda \lesssim 1$, the VEV of the Polonyi field begins to exceed the critical field value φ_c, see Eq. (127), whereas for $\lambda \gtrsim 3$, the Polonyi field begins to take super-Planckian values during inflation [this follows from Eq. (141)]. The black contours respectively indicate the values of Λ (left panel) and $m_{3/2}$ (right panel) required to obtain the correct scalar spectral amplitude, $A_s \simeq A_s^{\rm obs}$.

Interestingly enough, we are able to reproduce the observed value $n_s^{\rm obs}$ only in the "inflection-point regime" ($\epsilon < \epsilon_0$), in which the scalar potential does not exhibit any local extrema close to the inflection point at $\varphi = \varphi_{\rm flex}$. In the "hilltop regime" ($\epsilon > \epsilon_0$), on the other hand, the scalar spectral index always turns out to be small, $n_s \lesssim 0.94$. If we adopt the notion that the critical coefficient ϵ_0 should, in fact, be regarded as an upper bound on ϵ (because the inflaton would otherwise get trapped in a false vacuum), this result entails the following interesting observation: Maybe the coefficient ϵ has, for one reason or another, the tendency to saturate its upper bound from below, $\epsilon \simeq \epsilon_0$. The coefficient ϵ might, e.g., be intrinsically of $\mathcal{O}(1)$, and it only ends up being slightly suppressed, $\epsilon \simeq 0.2$, because it is bounded from above. Alternatively, ϵ might tend to be close to ϵ_0, because such a parameter choice leads to a particularly flat potential and, hence, to a particularly long period of inflation (i.e., to a large spatial volume in which $\epsilon \simeq \epsilon_0$). But anyhow, no matter what the origin of this fine-tuning is, the point is this: once we suppose that there is a physical reason that singles out ϵ values close to ϵ_0, we are automatically led to values of n_s close to the observed value. Or, put differently, it is an intriguing coincidence that the observed value of n_s is reproduced just in the vicinity of the only special point on the ϵ axis. One could have expected that successful inflation would require $\epsilon \simeq \epsilon_0$. But then it would not have been guaranteed that the resulting values of n_s would end up being close to $n_s^{\rm obs}$.

Another interesting coincidence pertains to the required value of the dynamical scale Λ. As evident from the left panel of Fig. 3, the scale Λ needs to take values very close to the scale of grand unification,

$$\Lambda \sim 10^{16}\ \mathrm{GeV}. \tag{144}$$

This suggest the fascinating possibility that the scales of supersymmetry breaking, inflation, and grand unification might, in fact, all be unified,

$$\Lambda_{\rm SUSY} \equiv \Lambda_{\rm inf} \equiv \Lambda_{\rm GUT}. \tag{145}$$

A trivial condition for this kind of "triple unification" is that the energy scales of inflation and supersymmetry breaking are identical to each other. The unification of the three scales $\Lambda_{\rm SUSY}$, $\Lambda_{\rm inf}$, and $\Lambda_{\rm GUT}$ can, thus, only be realized in the context of Polonyi inflation and in no other inflation model. In this paper, we have not made any attempt to embed our model into a larger GUT framework. More work on the possible connection between Polonyi inflation and grand unification is, therefore, certainly needed. As evident from the right panel of Fig. 3, the gravitino turns out to be superheavy in our model,

$$m_{3/2} \sim 10^{12}\ \mathrm{GeV}. \tag{146}$$

As already pointed out in the Introduction [see Eq. (5)], the gravitino mass is directly related to our prediction for the tensor-to-scalar ratio r (this relation readily follows from $m_{3/2} \simeq H_{\rm inf}$ in our model),

$$r \simeq 2 \times 10^{-5} \left(\frac{m_{3/2}}{10^{12}\ \mathrm{GeV}}\right)^2. \tag{147}$$

The experimental confirmation of this relation is certainly challenging—but *if* it should be accomplished, it would represent an unequivocal smoking-gun signal for Polonyi inflation, i.e., for a close link between the dynamics of supersymmetry breaking and inflation.

C. Neutralino dark matter and thermal leptogenesis

In the minimal framework of pure gravity mediation, gravitino masses of $\mathcal{O}(10^{12})$ GeV typically result in MSSM gaugino masses $M_{\tilde{g},\tilde{w},\tilde{b}}$ of $\mathcal{O}(10^{10})$ GeV,

$$M_{\tilde{g},\tilde{w},\tilde{b}} \sim \frac{m_{3/2}}{16\pi^2} \sim 10^{10}\ \mathrm{GeV}, \tag{148}$$

where the suppression by the loop factor $1/(16\pi^2)$ is a consequence of anomaly mediation [11,94]. Here, the wino $\tilde{w}$ typically ends up being the lightest gaugino. At tree level, the gaugino masses vanish in our model. To see this, recall that the Polonyi field Φ carries R charge 2. Couplings between the Polonyi field and the SM gauge fields of the form $\Phi \mathcal{W}_\alpha \mathcal{W}^\alpha$ are, thus, forbidden in the superpotential. After inflation, the Polonyi field predominantly decays into gravitinos (via the $|\Phi|^4$ term in the effective Kähler potential [95]), so that inflation is followed by a phase of gravitino domination in our scenario [71]. The non-thermal wino production in gravitino decays then easily overcloses the universe [96]. To achieve an acceptable present-day wino abundance, we, therefore, have to assume that the wino mass is somehow suppressed. Fortunately, this is possible in the context of pure gravity mediation, where the gaugino masses in Eq. (148) also receive threshold corrections from Higgsino loops [13,14]. These loop corrections are potentially of the same order of magnitude as the anomaly-mediated masses. The different contributions to the wino mass can, in particular, cancel, so that the mass $M_{\tilde{w}}$ is reduced down to[20]

[20]Alternatively, the wino mass may be reduced via gauge mediation [71]. We also mention that, in split-SUSY spectra [97], anomaly-mediated gaugino masses are absent altogether. In this case, gaugino masses may be as small as $m_{3/2}^3/M_{\rm Pl}^2 \sim 1$ GeV, which would even require some extra mass contributions. A split-SUSY spectrum may, e.g., be achieved, if supersymmetry breaking is a pure SUGRA effect [98]. It is an interesting open question how to realize a split-SUSY spectrum in the context of DSB models *à la* IYIT. We leave this question for future work. For now, we simply assume that the contribution to $M_{\tilde{w}}$ from anomaly mediation ends up being canceled by another mass contribution of equal magnitude and opposite sign.

$$M_{\tilde{w}} \sim 3\ \mathrm{TeV}. \tag{149}$$

We shall assume that this is the case for anthropic reasons. In case the wino mass is in the TeV range, the wino reaches thermal equilibrium after reheating and eventually represents an excellent candidate for dark matter in the form of weakly interacting massive particles (WIMPs). Larger wino masses, on the other hand, result in an overabundance of dark matter, thereby making our Universe hostile for life.

The phase of gravitino domination after Polonyi inflation ends once the gravitino begins to decay into the massless fields of the MSSM. Here, the corresponding gravitino decay rate is given as [71]

$$\Gamma_{3/2} \simeq \frac{193}{384\pi}\frac{m_{3/2}^3}{M_{\mathrm{Pl}}^2} \simeq 30\ \mathrm{MeV}\left(\frac{m_{3/2}}{10^{12}\ \mathrm{GeV}}\right)^3. \tag{150}$$

The reheating temperature reached during gravitino decay, T_{rh}, then takes the following value:

$$T_{\mathrm{rh}} \simeq \left(\frac{90}{\pi^2 g_*}\right)^{1/4}\sqrt{\Gamma_{3/2}M_{\mathrm{Pl}}} \simeq 1\times 10^8\ \mathrm{GeV}\left(\frac{m_{3/2}}{10^{12}\ \mathrm{GeV}}\right)^{3/2}, \tag{151}$$

where $g_* = 915/4$ denotes the effective number of relativistic DOFs in the MSSM at high temperatures. Contrary to many other inflation models, the reheating temperature after Polonyi inflation is, therefore, not a free parameter, which is controlled by the unknown strength of the inflaton couplings to matter. T_{rh} rather follows directly from the universal decay rate of the gravitino, which only depends on the gravitino mass. As the gravitino mass is, in turn, more or less fixed by the amplitude of the scalar power spectrum in Polonyi inflation (see Fig. 3), we arrive at the interesting result that Polonyi inflation makes a definite prediction for the reheating temperature: $T_{\mathrm{rh}} \sim 10^8$ GeV. Remarkably enough, this is not far away from the temperature that is needed for the successful realization of thermal leptogenesis [72]. In its simplest form, thermal leptogenesis requires the reheating temperature to be at least of $\mathcal{O}(10^9)$ GeV. But resonance effects in the case of mildly degenerate heavy-neutrino mass spectrum easily increase the efficiency of thermal leptogenesis also at lower temperatures [73]. The generation of the baryon asymmetry of the Universe after Polonyi inflation may, thus, very well be accounted for by thermal leptogenesis.

D. Benchmark scenario

Finally, we shall illustrate our findings by means of a concrete example. To this end, we will now consider the following point in parameter space, which may be regarded as a representative benchmark scenario:

$$\Lambda \simeq 1.26\times 10^{16}\ \mathrm{GeV}, \quad \lambda \simeq 1.66, \quad \epsilon \simeq 0.204. \tag{152}$$

For this choice of λ, the critical coefficient ϵ_0 is given by $\epsilon_0 \simeq 0.205$ [see Eq. (137)]. Our choice of ϵ is, therefore, fine-tuned in the sense that it deviates from ϵ_0 only by half a percent or so. For the parameter values in Eq. (152), we then obtain the following predictions for the inflationary CMB observables:

$$A_s \simeq 2.14\times 10^{-9}, \quad n_s \simeq 0.968, \quad r \simeq 6.01\times 10^{-6}, \tag{153}$$

where the values for A_s and n_s coincide with the experimental best-fit values by construction. Next to the parameters Λ, λ, and ϵ (which represent the free parameters in the SUSY-breaking sector), we also need to specify the parameters in the R-symmetry-breaking sector, i.e., the coefficients α and β as well as the energy scale v_P [see Eq. (83)]. Here, the scale v_P is required to take a value within a finite interval,

$$5.3\times 10^{16}\ \mathrm{GeV} \lesssim v_P \lesssim 1.1\times 10^{18}\ \mathrm{GeV}. \tag{154}$$

We recall that the lower bound on v_P is a consequence of the requirement that the heavy-quark mass scale in the R-symmetry-breaking sector should always exceed the dynamical scale, $\langle P\rangle \gtrsim \tilde{\Lambda}$ [see Eq. (130)], while the upper bound on v_P follows from the requirement that the dynamics of the R-symmetry-breaking sector should have no noticeable effect on inflation, $v_P \ll \epsilon^{1/2}M_{\mathrm{Pl}}$ [see Eq. (109)]. In the following, we shall use the geometric mean of the two boundary values in Eq. (154) as a characteristic value for v_P,

$$\bar{v}_P \simeq 2.41\times 10^{17}\ \mathrm{GeV}. \tag{155}$$

For the VEVs $\langle P\rangle$ and $\langle Y\rangle$ as well as for the dynamical scale $\tilde{\Lambda}$, we then obtain [see Eqs. (82) and (123)]

$$\begin{aligned} \langle P\rangle &\simeq 2.9\times 10^{17}\ \mathrm{GeV}\left(\frac{v_P}{\bar{v}_P}\right), \\ \langle Y\rangle &\simeq 1.2\times 10^{17}\ \mathrm{GeV}\left(\frac{v_P}{\bar{v}_P}\right), \\ \tilde{\Lambda} &\simeq 3.0\times 10^{16}\ \mathrm{GeV}\left(\frac{\bar{v}_P}{v_P}\right)^{1/2}. \end{aligned} \tag{156}$$

Similarly, we find the following relations for the two reference values α_0 and β_0 [see Eqs. (93) and (94)]:

$$\alpha_0 \simeq 4.0\times 10^{-6}\left(\frac{\bar{v}_P}{v_P}\right), \quad \beta_0 \simeq 4.0\times 10^{-6}\left(\frac{\bar{v}_P}{v_P}\right). \tag{157}$$

TABLE I. Numerical results for a number of important energy scales, masses and field values in the context of the benchmark scenario discussed in Sec. IV D. The input parameter values used in this scenario are listed in Eqs. (152), (155), and (157). Quantities labeled with a star (*) are evaluated for $v_P = \bar{v}_P$. Their scaling with v_P is indicated in Eqs. (156) and (157).

Quantity	Symbol	Value (GeV)	Reference
Energy and mass scales in the SUSY-breaking sector			
Dynamical scale	Λ	1.3×10^{16}	Eq. (10)
Effective supersymmetric mass scale	$\bar{m}$	1.3×10^{16}	Eq. (37)
Soft SUSY-breaking mass scale	m	1.7×10^{15}	Eq. (29)
SUSY breaking scale	μ	1.5×10^{15}	Eq. (28)
Energy scale of inflation	$V_*^{1/4}$	1.5×10^{15}	Eq. (142)
Effective potential energy scale at large field values	$\Lambda_{\rm HE}$	7.4×10^{14}	Eq. (44)
Effective potential energy scale at small field values	$\Lambda_{\rm LE}$	5.9×10^{14}	Eq. (44)
Effective Polonyi mass around the origin	$m_{\rm eff}$	3.2×10^{13}	Eq. (43)
Gravitino mass	$m_{3/2}$	5.6×10^{11}	Eq. (72)
Inflationary Hubble rate	$H_{\rm inf}$	5.6×10^{11}	Eq. (2)
Energy and mass scales in the R-symmetry-breaking sector			
VEV of the field P in the true vacuum*	$\langle P \rangle$	2.9×10^{17}	Eq. (123)
Characteristic value for the scale v_P	$\bar{v}_P$	2.4×10^{17}	Eq. (155)
Gaugino condensation scale	$\Lambda_{\rm eff}$	6.4×10^{16}	Eq. (78)
VEV of the field Y in the true vacuum*	$\langle Y \rangle$	1.2×10^{17}	Eq. (123)
Dynamical scale*	Λ	3.0×10^{16}	Eq. (82)
Constant term in the superpotential	$w^{1/3}$	1.5×10^{16}	Eq. (77)
Mass of the field P in the true vacuum*	m_p	1.5×10^{12}	Eq. (122)
Mass of the field Y in the true vacuum*	m_y	1.5×10^{12}	Eq. (122)
Important Polonyi field values during and after inflation			
Beginning of slow-roll inflation	$\varphi_{\rm sr}^-$	$1.04 M_{\rm Pl}$	Eq. (140)
Location of the inflection point	$\varphi_{\rm flex}$	$0.72 M_{\rm Pl}$	Eq. (135)
Field value $N_e^* = 55$ e-folds before the end of inflation	φ_*	$0.70 M_{\rm Pl}$	Eq. (141)
End of slow-roll inflation	$\varphi_{\rm sr}^+$	$0.43 M_{\rm Pl}$	Eq. (140)
Critical field value	φ_c	1.5×10^{16}	Eq. (41)
Polonyi VEV in the true vacuum	$\langle \varphi \rangle$	3.6×10^{15}	Eq. (123)

The actual value of α is given by $a_1^{1/2}(\varphi_0)\alpha_0$ [see Eq. (100)], where $a_1(\varphi_0)$ varies as a function of φ_0, i.e., the Polonyi field value at the onset of gaugino condensation in the R-symmetry-breaking sector. Allowing for φ_0 values in the range $0 \lesssim \varphi_0 \lesssim \varphi_c$, the function $a_1(\varphi_0)$ roughly takes values between 1.00 and 1.02. Up to deviations of the order of one percent or so, the coefficient α, therefore, coincides with the reference value α_0. Meanwhile, the coefficient β can be freely chosen, as long as we respect the constraint $|\beta| \lesssim \beta_0$.

The parameter values in Eqs. (152), (155), and (157) now completely fix the numerical properties of our benchmark scenario. In Table I, we give an overview of all the resulting numerical values for the various energy scales and field values of interest in our model. Table I, thus, provides an example for a possible realization of our model, which illustrates how to achieve successful Polonyi inflation as well as spontaneous R symmetry breaking at late times in the context of strongly coupled gauge theories.

V. CONCLUSIONS AND OUTLOOK

The large value of the SM Higgs boson mass as well as the null result of SUSY searches at the LHC call for a paradigm shift in our expectations towards the role of supersymmetry in nature. From the perspective of string theory, dark matter, and grand unification, supersymmetry still represents a well-motivated extension of the standard model. But in view of the recent experimental data, we now begin to realize that supersymmetry's main purpose may actually not lie in stabilizing the electroweak scale. Instead, it now appears more likely that supersymmetry is, in fact, spontaneously broken at a scale much higher than the electroweak scale. This opens up a whole range of new phenomenological possibilities.

As we have demonstrated in this paper, the spontaneous breaking of supersymmetry at a high scale, $\Lambda \sim 10^{16}$ GeV, may, e.g., offer a dynamical explanation for the occurrence of cosmic inflation in the early universe. We have dubbed the ensuing inflationary scenario *Polonyi inflation*, as it identifies cosmic inflation as a natural by-product of spontaneous supersymmetry breaking in an effective Polonyi model. Generally speaking, the key idea of our proposal is to obtain inflation from the Polonyi superpotential,

$$W = \mu^2 \Phi + w, \tag{158}$$

with the Polonyi superfield Φ breaking supersymmetry via its nonvanishing F-term, $\langle |F_\Phi| \rangle = \mu^2$, and acting as the chiral inflaton field at the same time. In the context of Polonyi inflation, spontaneous supersymmetry breaking and cosmic inflation are nothing but two sides of the same coin: Inflation is driven by the vacuum energy density associated with the spontaneous breaking of supersymmetry and the scalar inflaton is identified as the complex Polonyi field, i.e., the pseudoflat direction in the scalar potential of the SUSY-breaking sector. This connection between supersymmetry breaking and inflation results in several characteristic parameter relations that cast a new light on well-known quantities. In Polonyi inflation, the Hubble rate during inflation, e.g., is equal to the gravitino mass at low energies,

$$H_{\rm inf} \simeq m_{3/2} \sim 10^{12}\ {\rm GeV}. \tag{159}$$

This implies that the amplitude of the primordial density fluctuations, $\delta\rho/\rho$, scales with $m_{3/2}$,

$$\begin{aligned} \frac{\delta\rho}{\rho} &= A_s^{1/2} = \frac{\sqrt{2}}{\pi} \frac{1}{r^{1/2}} \frac{m_{3/2}}{M_{\rm Pl}} \\ &\sim 10^{-5} \left(\frac{10^{-4}}{r}\right)^{1/2} \left(\frac{m_{3/2}}{10^{12}\ {\rm GeV}}\right). \end{aligned} \tag{160}$$

There are indications that the observed value of $\delta\rho/\rho$ might be the result of anthropic selection [70]. Given typical values of the tensor-to-scalar ratio r in slow-roll inflation, the above relation then points towards very large values of the gravitino mass. In Polonyi inflation, the ultimate reason why supersymmetry is broken at a high scale, therefore, consists in the necessity of reproducing the anthropic value of $\delta\rho/\rho$. The general concept of Polonyi inflation, thus, not only yields an answer to the question "Why inflation?," it also explains why supersymmetry is necessarily broken at a very high scale. Moreover, we point out that $m_{3/2} \sim 10^{12}$ GeV is an interesting result in view of the stability of the SM Higgs potential. Recall that, solely within the standard model, the Higgs quartic coupling (most likely) turns negative at energies around $\mathcal{O}(10^{11\ldots12})$ GeV [99]. This instability can be remedied by supersymmetry—as long as the soft sparticle masses are at most of $\mathcal{O}(10^{12})$ GeV. In this sense, our result for $m_{3/2}$ turns out to saturate the upper bound on the gravitino mass implied by vacuum stability. This may or may not be a coincidence.

In this paper, we have constructed a minimal model that illustrates how the idea of Polonyi inflation may be realized in the context of strongly coupled supersymmetric gauge theories. Our construction consists of two separate hidden sectors that respectively account for the dynamical origin of the parameters μ and w in Eq. (158). Here, one hidden sector is identical to the simplest version of the IYIT model of dynamical supersymmetry breaking. This sector contains the Polonyi field Φ and is responsible for the generation of the SUSY-breaking parameter μ via strong dynamics. The other hidden sector features the same matter content as the SUSY-breaking sector, but contains fewer singlet fields. The masses of the matter fields in this sector are controlled by the VEV of a singlet field P. At early times, $\langle P \rangle$ vanishes. At this time, the constant in the superpotential, $W \supset w$, is zero. As we are able to show, the effective scalar potential for the Polonyi field at this stage then takes the same form as the inflaton potential of ordinary F-term hybrid inflation in the limit of a vanishing gravitino mass. It is this potential that gives rise to the actual period of Polonyi inflation.[21] At small Polonyi field values after the end of slow-roll inflation, the Hubble-induced mass for the field P decreases—until, at a certain field value, P becomes tachyonically unstable. This instability results in a waterfall transition in the R-symmetry-breaking sector, such that $\langle P \rangle \neq 0$. The second hidden sector then turns into a pure SYM theory, in which R symmetry is broken via gaugino condensation. In consequence of that, the constant term w appears in the superpotential. Here, the final value of w needs to be fine-tuned, such that the CC in the true vacuum vanishes.

To sum up, we conclude that our effective Polonyi model results in an inflationary scenario similar to ordinary F-term hybrid inflation, apart from a few key differences: (i) All corrections to the scalar potential proportional to $m_{3/2}$ are missing in our model. (ii) This means, in particular, that the scalar potential does not contain any odd powers of the inflaton field. Because of that, the inflaton potential does not depend on the complex phase of the inflaton field, as it is usually the case in F-term hybrid inflation. (iii) Moreover, in our scenario, the low-energy value of the gravitino mass, $m_{3/2}$, is not bounded from above by the inflationary Hubble rate. In Polonyi inflation, these two scales are, in fact, equal to each other, $m_{3/2} \simeq H_{\rm inf}$, whereas in F-term hybrid inflation we have to demand that $m_{3/2} \lesssim 10^{-3} H_{\rm inf}$, so as to ensure that the slow-roll conditions do not get violated. (iv) In Polonyi inflation, the inflaton sector does not undergo a waterfall phase transition at the end of inflation. Instead, the inflationary vacuum energy density continues to act as the vacuum energy density associated with the spontaneous breaking of supersymmetry after inflation. (v) Meanwhile, small inflaton field values trigger a (harmless) waterfall phase transition in a separate hidden sector, in the course of which only an

[21] It is amusing to note that the scalar potential of Polonyi inflation is identical to the potential of F-term hybrid inflation in the special limit $m_{3/2} \to 0$—albeit, in the true vacuum at low energies, $m_{3/2}$ is exceptionally large in our model.

approximate global Z_2 symmetry becomes broken. In contrast to F-term hybrid inflation, Polonyi inflation, therefore, does not suffer from the production of dangerous topological defects at the end of inflation.

In the present paper, we have only touched on the phenomenological implications of our model and more work is certainly needed. For one thing, the study of the inflationary phase may still be further refined. As it turns out, inflation takes place at field values only shortly below the Planck scale in our model. Inflaton terms of $\mathcal{O}(\varphi^6)$ in the scalar potential may, therefore, have a noticeable effect on the predictions for the inflationary CMB observables. We do not expect these terms to change our conclusions qualitatively. But quantitatively, they may be relevant. For another thing, the end of inflation and the subsequent reheating phase require a closer examination. Here, it would be interesting to study the implications of different interactions between the two hidden sectors in the Kähler potential more comprehensively. Also, the actual dynamics of reheating deserve further investigation. Our model does not suffer from the usual Polonyi problem, as the decay of the Polonyi field is identified with the first stage of the reheating process itself. But besides that, a dedicated study of reheating tracking the expansion of the universe during the transition to the radiation-dominated era as well as the oscillatory motion of all scalar fields is needed. Here, particular attention should be paid to the study of all possible decay modes of the Polonyi field as well as of the gravitino. A better understanding of reheating would then enable us to better assess the prospects of successful thermal leptogenesis. Finally, the phenomenological implications of our model at low energies need to be explored in more detail. As examples of observational signatures at low energies, we merely mention two interesting possibilities: (i) In our scenario, a wino LSP with an anthropically selected mass around 3 TeV accounts for dark matter. This wino can be searched for in direct- and indirect-detection experiments. In addition, the almost mass-degenerate chargino may be seen in the form of macroscopic charged tracks in collider experiments. All in all, our model predicts that the neutral and the charged wino are the only sparticles that should show up at low energies. All other sparticles have masses of at least $\mathcal{O}(10^{10})$ GeV and are, thus, expected to be decoupled from low-energy phenomenology for the most part. (ii) Moreover, the discrete Z_4^R symmetry that we use to forbid the constant term in the superpotential at early times predicts the existence of vectorlike matter fields charged under the SM gauge group, which cancel the SM contributions to the Z_4^R gauge anomalies. These vectorlike matter fields may have masses within the reach of collider experiments, which would allow to probe our assumption of an anomaly-free R symmetry at accessible energies.

Last but not least, one should attempt to implement our idea of Polonyi inflation into other dynamical models. It would be interesting to assess which other DSB models apart from the IYIT model might serve as a UV completion of the inflationary dynamics. In addition to that, a more systematic study of the Polonyi Kähler potential—beyond near-canonical and approximately shift-symmetric Kähler potentials—would be desirable. Maybe there are particular (exotic) choices for the Kähler potential that render it unnecessary to supplement the tree-level Polonyi model with radiative corrections, as we have done it in this paper. Likewise, our model of late-time R symmetry breaking may be extended or supplemented by alternative means to generate the constant term in the superpotential at the end of inflation. As an alternative to our scenario, which employs the waterfall field P, one might consider a direct coupling of the Polonyi field to the gauge fields of a strongly coupled pure SYM theory via the gauge-kinetic function,

$$\mathcal{L} \supset \int d^2\theta \frac{1}{4}\left(\frac{1}{g_0^2} + \frac{\Phi^n}{\phi_0^n}\right)\mathcal{W}_\alpha \mathcal{W}^\alpha + \text{H.c.}, \quad n = 1, 2, \ldots. \tag{161}$$

For appropriately chosen values of the parameters g_0 and ϕ_0 and the integer n, the SYM theory is weakly coupled during inflation, returning to the strongly coupled regime only at the end of inflation. This results in gaugino condensation and, hence, late-time R symmetry breaking at $|\phi| \lesssim \phi_0$. We will give a more careful discussion of this mechanism in the context of Polonyi inflation elsewhere. Furthermore, gaugino condensation in a pure SYM theory is not the only way to break R symmetry and other possibilities should be considered as well. Eventually, the model presented in this paper may only be regarded as a first step towards a new understanding of the close relation between inflation and spontaneous supersymmetry breaking. The concept of Polonyi inflation leads us into uncharted territory and we are convinced that, exploring these new avenues, we will encounter both big surprises and valuable insights.

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APPENDIX A: EFFECTIVE SCALAR POTENTIAL FOR THE POLONYI/INFLATON FIELD

In this appendix, we give the details of our derivation of the effective one-loop potential, $V_{1-\text{loop}}$, for the Polonyi field $\phi = \varphi/\sqrt{2}e^{i\theta}$, i.e., the pseudoflat scalar direction in the IYIT DSB model, see Eqs. (42), (43), (44), and (45) in Sec. II C. The starting point of our computation is the Coleman-Weinberg formula for the effective one-loop potential [80], which takes the following form in supersymmetric theories:

$$V_{1-\mathrm{loop}}(\varphi) = \frac{1}{64\pi^2}\,\mathrm{STr}\left[\mathcal{M}^4(\varphi)\left(\ln\left(\frac{\mathcal{M}^2(\varphi)}{Q^2}\right) + c\right)\right]. \tag{A1}$$

Here, Q is the renormalization scale and c denotes a dimensionless constant that is introduced for notational convenience. We note that Q and c are only defined up to the following transformation:

$$c \to c + \Delta c, \qquad Q \to e^{\Delta c/2} Q, \tag{A2}$$

so that the constant c can always be removed by a finite renormalization. Equation (A1) now tells us that the potential $V_{1-\mathrm{loop}}$ is given as the supertrace (STr) of a specific function of the mass matrix $\mathcal{M}$, which encompasses all bosonic and fermionic mass eigenvalues in the presence of a nonvanishing Polonyi field background. In a first step, we, therefore, need to determine the full mass spectrum of the IYIT sector.

1. Mass spectrum of the low-energy effective theory

Apart from the Polonyi field Φ (which is massless at tree level in global supersymmetry), the IYIT sector of our model consists of six meson flavors, Ξ^0 and X^n, and six singlet fields, Σ and S_n. The scalar and fermionic masses for these chiral multiplets follow from the effective tree-level superpotential in Eq. (27). For an arbitrary background value of the Polonyi field φ, we find the following scalar mass eigenvalues:

$$M_a^2(\varphi;p,q) = m_a^2 + \frac{1}{2}(M^2(\varphi) + qm^2) + \frac{p}{2}[(M^2(\varphi) + qm^2)^2 + (2m_a M(\varphi))^2]^{1/2}, \tag{A3}$$

where $a = 0$ refers to the zeroth flavor, i.e., to the fields Ξ^0 and Σ, while $a = 1, \ldots, 5$ refers to the n other flavors, i.e., to the fields X^n and S_n. Moreover, the discrete parameter $p = \pm 1$ distinguishes between the two different types of particles involved, i.e., between mesons ($p = +1$) and singlets ($p = -1$), while the parameter $q = \pm 1$ accounts for the soft mass splitting within each complex scalar, $\pm m^2$, in consequence of spontaneous supersymmetry breaking. In this respect, Eq. (A3) serves as an illustration of our statement in Sec. II B, where we note that the mass parameter m plays the role of the soft SUSY-breaking mass scale in the IYIT model, see the discussion below Eq. (29). By setting the soft mass m to zero, we then readily obtain the mass eigenvalues of the fermionic components in the meson and singlet multiplets,

$$m = 0 \Rightarrow \tilde{M}_a^2(\varphi;p) = m_a^2 + \frac{1}{2}M^2(\varphi) + \frac{p}{2}[M^4(\varphi) + (2m_a M(\varphi))^2]^{1/2}. \tag{A4}$$

For $\varphi \neq 0$, all scalar and fermionic masses are different from each other (as long as none of the Dirac masses m_a are identical). Away from the origin, the meson and singlet fields, therefore, give rise to, in total, 24 real scalars and 12 Majorana fermions. At the origin, $M_a^2(\varphi;p,q)$ and $\tilde{M}_a^2(\varphi;p)$ reduce to

$$M(\varphi) = 0 \Rightarrow M_a^2(0;p,q) = m_a^2 + \frac{1}{2}(p+q)m^2, \qquad M_a^2(0;p,q) = m_a^2, \tag{A5}$$

which illustrates that, for $\varphi = 0$, half of the real scalars and all of the Majorana fermions pair up to form complex scalars and Dirac fermions, respectively. At $\varphi = 0$, we then end up with 12 real and 6 complex scalars as well as with 6 Dirac fermions. Here, the complex scalars and the Dirac fermions are, in particular, composed half of mesonic and half of singlet DOFs.

Furthermore, we note that none of the masses $M_a^2(\varphi;p,q)$ and $\tilde{M}_a^2(\varphi;p)$ ever turn tachyonic,[22] which means that the meson and singlet fields always remain stabilized at the origin. This should be compared with the situation in ordinary F-term hybrid inflation, where the masses of the FHI waterfall fields correspond to (the nonzero values of) $M_a^2(\varphi;p,q)$ and $\tilde{M}_a^2(\varphi;p)$ in the limit of zero Dirac masses,

$$m_a = 0 \Rightarrow$$
$$M_a^2(\varphi;p,q) = \frac{1}{2}(M^2(\varphi) + qm^2) + \frac{p}{2}|M^2(\varphi) + qm^2| \to M^2(\varphi) + qm^2,$$
$$\tilde{M}_a^2(\varphi;p) = \frac{1}{2}(1+p)M^2(\varphi) \to M^2(\varphi). \tag{A6}$$

In this case, the scalar mass eigenstates corresponding to $q = -1$ are tachyonically unstable, once $M(\varphi) < m$. This instability is absent in our "waterfall-transition-free" version of F-term hybrid inflation.

Finally, before moving on and presenting our results for the effective Polonyi potential, let us state approximate expressions for $M_a^2(\varphi;p,q)$ and $\tilde{M}_a^2(\varphi;p)$ that are valid at small and large values of the Polonyi field, respectively, and which will become useful later on. In the low-energy

[22] A necessary and sufficient condition for $M_a^2(\varphi;p,q) \geq 0$ is that the soft SUSY-breaking mass scale m does not exceed any of the supersymmetric Dirac masses: $m_a \geq m$ for all flavors a, which is always satisfied in our model, see Eq. (30).

regime close to the origin in field space, i.e., at small values of the order parameter, $x(\varphi) = \varphi/\varphi_c \lesssim 1$ [see Eq. (37)], we find

$$x(\varphi) \lesssim 1 \Rightarrow$$

$$M_a^2(\varphi; p, q) = m_a^2\left(1 + p\frac{M^2(\varphi)}{m^2}\right) + \frac{q}{2}(p+q)(M^2(\varphi) + qm^2) + \mathcal{O}(x^4), \quad \text{(A7)}$$

$$\tilde{M}_a^2(\varphi; p) = m_a^2 + \frac{1}{2}M^2(\varphi) + p m_a M(\varphi) + \mathcal{O}(x^3),$$

which reduces to the expressions in Eq. (A5) in the limit $\varphi \to 0$. Conversely, in the high-energy regime, at large values of $x(\varphi)$, the masses $M_a^2(\varphi; p, q)$ and $\tilde{M}_a^2(\varphi; p)$ may be approximated as follows:

$$x(\varphi) \gtrsim 1 \Rightarrow$$

$$M_a^2(\varphi; p, q) = \frac{1}{2}(1+p)(M^2(\varphi) + 2m_a^2 + qm^2) - p\frac{m_a^2(m_a^2 + qm^2)}{M^2(\varphi)} + \mathcal{O}(x^{-4}), \quad \text{(A8)}$$

$$\tilde{M}_a^2(\varphi; p) = \frac{1}{2}(1+p)(M^2(\varphi) + 2m_a^2) - p\frac{m_a^4}{M^2(\varphi)} + \mathcal{O}(x^{-4}).$$

In this regime, the meson masses ($p = +1$) asymptotically approach the usual expressions from F-term hybrid inflation, see Eq. (A6), while the singlet masses ($p = -1$) only acquire suppressed masses of $\mathcal{O}(m_a^2/M)$. Here, note that the small singlet masses in the large-field regime are the result of a seesaw mechanism of a sort: For a large Majorana mass $M(\varphi)$, the Dirac masses m_a become suppressed by a factor of $\mathcal{O}(m_a/M)$. This is exactly what happens in the usual seesaw mechanism, which explains the small masses of the SM neutrinos as the outcome of large Majorana and small Dirac neutrino masses [100].

2. Coleman-Weinberg one-loop effective potential

We are now in the position to calculate the effective one-loop potential for the scalar Polonyi field. The full result for $V_{1-\text{loop}}$ follows from Eqs. (A1), (A3), and (A4) and takes a rather complicated form. For this reason, we refrain from explicitly writing down the full expression for $V_{1-\text{loop}}$, and merely refer to Fig. 1, where we plot the exact result for $V_{1-\text{loop}}$ as a function of the order parameter $x(\varphi)$. Instead, let us now evaluate $V_{1-\text{loop}}$ for specific field values of interest. For instance, at the origin, we find

$$V_{1-\text{loop}}(0) = \frac{m^4}{32\pi^2}\left[N_X\left(\ln\left(\frac{\bar{m}^2}{Q^2}\right) + c\right) + \sum_a L(r_a)\right], \quad \text{(A9)}$$

where $N_X = 6$ counts the number of meson flavors in the IYIT model, $\bar{m} = \Lambda$ denotes the effective supersymmetric mass scale in the IYIT sector [see Eqs. (37) and (38)], and where L is a loop function that needs to be evaluated at the respective mass ratios r_a [see Eq. (29)],

$$L(r_a) = \frac{1}{2}\left(1 + \frac{1}{r_a^2}\right)^2 \ln(1 + r_a^2) + \frac{1}{2}\left(1 - \frac{1}{r_a^2}\right)^2 \ln(1 - r_a^2) = \frac{3}{2} + \mathcal{O}(r_a^4). \quad \text{(A10)}$$

Note that this result for $V_{1-\text{loop}}(0)$ can also be obtained by plugging the masses in Eq. (A5) into the CW formula. The function L is nearly constant over its entire domain, $2\ln 2 = L(1) \le L(r_a) \le L(0) = 3/2$, and, thus, well approximated by $L(r_a) \approx 3/2$. This motivates the following choices for Q and c:

$$Q = \bar{m}, \qquad c = -\frac{3}{2}. \quad \text{(A11)}$$

For these values of the renormalization scale Q and the constant c, the tree-level vacuum energy density of the IYIT model, $V_0 = \langle |F_\Phi| \rangle^2 = \mu^4$, receives a small negative shift ΔV_0 of $\mathcal{O}(m^8/m_a^4)$,

$$\Delta V_0 = V_{1-\text{loop}}(0) = -N_X\frac{m^4}{32\pi^2}\left(\frac{3}{2} - \frac{1}{N_X}\sum_a L(r_a)\right). \quad \text{(A12)}$$

This correction to the vacuum energy density and, hence, the ratio $|\Delta V_0|/V_0$, are bounded from above,

$$|\Delta V_0| \le N_X\frac{m^4}{32\pi^2}\left(\frac{3}{2} - 2\ln 2\right), \qquad \frac{|\Delta V_0|}{V_0} \lesssim 10^{-3}(1 + r_0^2)\lambda^2. \quad \text{(A13)}$$

Here, we have used that L takes its smallest value in the limit $r_a \to 1$, where $L \to 2\ln 2$. For perturbative values of the Yukawa coupling, $\lambda \lesssim 4$, we can, therefore, safely neglect the radiative correction ΔV_0.

Having fixed Q and c, let us now study the effective potential for extreme values of the order parameter, i.e., for $x(\varphi) \ll 1$ and $x(\varphi) \gg 1$. Below the critical field value, the inflaton-induced mass $M(\varphi)$ is smaller than the supersymmetric Dirac masses m_a. At energies below m_a, the "heavy" meson flavors can then be integrated out, which results in a quadratic potential for φ around the origin,

$$V_{1-\text{loop}}^{\text{LE}}(\varphi) = \frac{1}{2}m_{\text{eff}}^2\varphi^2 + \mathcal{O}(x^4), \quad \text{(A14)}$$

where $m_{\rm eff}$ denotes the effective Polonyi mass at one-loop level [see also Eq. (43) and footnote 7],

$$m_{\rm eff}^2 = (2\ln 2 - 1) N_X^{\rm eff}(r_a) \frac{\kappa_\Phi^2}{16\pi^2} m^2, \quad N_X^{\rm eff}(r_a) = \sum_a \omega(r_a). \tag{A15}$$

Here, ω represents a loop function that acts as a normalized weight, $0 \leq \omega \leq 1$, for each meson flavor,

$$\omega(r_a) = \frac{\ell(r_a)}{2\ln 2 - 1}, \quad \ell(r_a) = \frac{1}{2}\left(1 + \frac{1}{r_a^2}\right)^2 \ln(1 + r_a^2) - \frac{1}{2}\left(1 - \frac{1}{r_a^2}\right)^2 \ln(1 - r_a^2) - \frac{1}{r_a^2}. \tag{A16}$$

Evaluated at the mass ratio r_a, the function ω quantifies the contribution from the respective meson flavor to $m_{\rm eff}$. Correspondingly, the sum over all factors $\omega(r_a)$ yields the effective number of flavors, $N_X^{\rm eff}$, that contribute to the effective Polonyi mass. Moreover, we note that, over its entire domain, $0 \leq r_a \leq 1$, the function ω is well approximated by $\omega(r_a) \approx r_a^2$. This allows us to rewrite $V_{1-\rm loop}$ in the following way:

$$\begin{aligned} V_{1-\rm loop}^{\rm LE}(\varphi) &= \frac{m^2}{16\pi^2} M^2(\varphi) \sum_a \ell(r_a) + \mathcal{O}(x^4) \\ &\approx (2\ln 2 - 1) \frac{m^2}{16\pi^2} M^2(\varphi) \sum_a r_a^2 + \mathcal{O}(x^4) \\ &= (2\ln 2 - 1) \frac{m^4}{16\pi^2} \sum_a R_a^2(\varphi) + \mathcal{O}(x^4), \end{aligned} \tag{A17}$$

which provides us with a couple of alternative expressions for $V_{1-\rm loop}$. For instance, the first line of Eq. (A17) makes explicit the dependence of the effective potential on the two mass scales m and $M(\varphi)$ at small field values, $V_{1-\rm loop}(\varphi) \propto m^2 M^2(\varphi)$. It is interesting to note that this result for the effective potential can also be obtained by plugging the approximate expressions for the scalar and fermionic masses in Eq. (A7) into the CW formula in Eq. (A1). Meanwhile, the third line of Eq. (A17) illustrates how $V_{1-\rm loop}$ may be expressed as a function of the ratios $R_a(\varphi) = M(\varphi)/m_a$ [see Eq. (36)]. In this case, the effective potential turns out to be given by the fourth power of the soft SUSY-breaking mass, m^4, times a function of the SUSY-preserving mass parameters $M(\varphi)$ and m_a, $V_{1-\rm loop}(\varphi) \propto m^4 f(R_a)$. The expressions for $V_{1-\rm loop}$ in Eq. (A17), therefore, comply with the expectation that, in supersymmetric theories, the effective potential ought to be proportional to the soft SUSY-breaking mass scale m [25]. Moreover, we can write $V_{1-\rm loop}$ in the small-field regime directly as a function of the order parameter,

$$V_{1-\rm loop}^{\rm LE}(\varphi) = \Lambda_{\rm LE}^4 x^2(\varphi) + \mathcal{O}(x^4), \tag{A18}$$

where we have introduced $\Lambda_{\rm LE}$ as the effective potential energy scale at small field values,

$$\begin{aligned} \Lambda_{\rm LE}^4 &= \frac{m^2 \bar{m}^2}{16\pi^2} \sum_a \ell(r_a) \\ &\approx (2\ln 2 - 1) \frac{m^4}{16\pi^2} \sum_a \left(\frac{\bar{m}}{m_a}\right)^2. \end{aligned} \tag{A19}$$

Finally, let us examine the effective Polonyi potential at small field values in the limit of small supersymmetry breaking, i.e., for a small soft SUSY-breaking mass m in comparison to large SUSY-preserving Dirac masses, $m \ll m_a$. This limit is best quantified in terms of the geometric mean of all mass ratios r_a, which we will refer to as the hierarchy parameter y in the following [see Eq. (29)]:

$$y = \left(\prod_a r_a\right)^{1/N_X} = \frac{m}{\bar{m}}, \tag{A20}$$

which can take values between 0 and 1 [see Eq. (30)]. Here, $y = 0$ corresponds to the SUSY-preserving limit, while $y = 1$ represents the case of maximal supersymmetry breaking. From Eqs. (29) and (38), it immediately follows that the hierarchy parameter y is, in fact, nothing but an alternative measure for the strength of the Yukawa coupling of the Polonyi field to the IYIT matter fields,

$$y = \frac{\lambda}{\eta}. \tag{A21}$$

That is, for perturbative values of the Yukawa coupling λ, supersymmetry is only mildly broken in the meson and singlet multiplets in the IYIT sector, while in the strongly coupled limit, i.e., for nonperturbative values of the Yukawa coupling λ, the effect of supersymmetry breaking is maximal. In passing, we mention that the parameter y also quantifies the ratio between the order parameter of our model, $x(\varphi)$, and the order parameter of the waterfall transition in ordinary F-term hybrid inflation, $R(\varphi)$,

$$x(\varphi) = y R(\varphi), \qquad R(\varphi) = \frac{M(\varphi)}{m} = \frac{|\phi|}{v} = \frac{\varphi}{\sqrt{2}v}. \tag{A22}$$

In the limit of maximal supersymmetry breaking, both order parameters, therefore, coincide with each other. In all other cases, we have $R(\varphi) > x(\varphi)$. This implies that, for $y < 1$, the mesonic phase transition in the IYIT model (see Fig. 1) takes place around a critical field value, $\varphi \sim \varphi_c$, that is parametrically larger than the critical field value associated with the waterfall transition in F-term hybrid inflation, $\varphi = \sqrt{2}v$. Let us now turn to the limit of small supersymmetry breaking. For small values of y, i.e., for small mass ratios r_a, the weight function ω in Eq. (A16) can be approximated as follows:

$$\omega(r_a) = \frac{r_a^2}{3(2\ln 2 - 1)} + \mathcal{O}(r_a^6). \tag{A23}$$

The effective potential $V_{1-\mathrm{loop}}$ and the potential energy scale Λ_{LE} at small field values then reduce to

$$V_{1-\mathrm{loop}}^{\mathrm{LE}}(\varphi) = \frac{m^4}{48\pi^2}\sum_a R_a^2(\varphi) + \mathcal{O}(x^4, y^8),$$
$$\Lambda_{\mathrm{LE}}^4 = \frac{m^4}{48\pi^2}\sum_a \left(\frac{\bar{m}}{m_a}\right)^2 + \mathcal{O}(y^8). \tag{A24}$$

We will come back to these results in Sec. A 3.

Next, we shall examine the effective Polonyi potential for large values of the order parameter $x(\varphi)$. In the large-field regime, $x(\varphi) \gtrsim 1$, the IYIT matter fields acquire a large field-dependent mass, such that their "bare" supersymmetric masses m_a become irrelevant, $M(\varphi) \gtrsim m_a$. Integrating out the matter fields then results in the usual logarithmic effective potential known from ordinary F-term hybrid inflation,

$$V_{1-\mathrm{loop}}^{\mathrm{HE}}(\varphi) = N_X\frac{m^4}{16\pi^2}\ln x(\varphi) + \mathcal{O}(x^{-4}),$$
$$x(\varphi) = \frac{\varphi}{\varphi_c} = \frac{M(\varphi)}{\bar{m}}. \tag{A25}$$

At large field values, the dependence of the effective potential on the mass scales m and $M(\varphi)$ is, therefore, of the form $V_{1-\mathrm{loop}}(\varphi) \propto m^4 \ln M(\varphi)$. Up to corrections of $\mathcal{O}(x^{-2})$, we are able to reproduce this result by inserting the approximate masses in Eq. (A8) into the CW formula in Eq. (A1). Moreover, we can read off the effective potential energy scale in the large-field regime, Λ_{HE}, from Eq. (A25),

$$V_{1-\mathrm{loop}}^{\mathrm{HE}}(\varphi) = \Lambda_{\mathrm{HE}}^4 \ln x(\varphi) + \mathcal{O}(x^{-4}),$$
$$\Lambda_{\mathrm{HE}}^4 = N_X\frac{m^4}{16\pi^2}, \tag{A26}$$

which is again of $\mathcal{O}(m^4)$. Last but not least, we can rewrite the effective potential as a function of the mass ratios $R_a(\varphi)$, which provides us with a large-field counterpart to the small-field result in Eq. (A24),

$$V_{1-\mathrm{loop}}^{\mathrm{HE}}(\varphi) = \frac{m^4}{16\pi^2}\sum_a \ln R_a(\varphi) + \mathcal{O}(x^{-4}). \tag{A27}$$

3. Reformulation in terms of an effective Kähler potential

At tree level, the F-term scalar potential is determined by two input functions, the superpotential W and the Kähler potential K. Here, the superpotential does not receive any radiative corrections in perturbation theory according to the SUSY nonrenormalization theorem [23]. The Kähler potential, on the other hand, *is* renormalized, which allows one to rewrite parts of the effective action in the form of an effective Kähler potential. We shall now demonstrate how this applies to our results in the previous section. First, we note that, schematically, the full result for $V_{1-\mathrm{loop}}$ in Eq. (A1) is of the following form:

$$V_{1-\mathrm{loop}}(\varphi) = \frac{m^4}{16\pi^2}\sum_{n=0}^{\infty}\frac{f_n(R_a)}{x^{4n}(\varphi)}y^{4n}$$
$$= \frac{m^4}{16\pi^2}f_0(R_a) + \mathcal{O}(y^8), \tag{A28}$$

where the order parameter $x(\varphi) = M(\varphi)/\bar{m}$ and the mass ratios $R_a = M(\varphi)/m_a$ denote the ratios of supersymmetric mass scales, while the hierarchy parameter $y = m/\bar{m}$ quantifies the ratio between the soft SUSY-breaking mass scale m and the effective supersymmetric mass scale $\bar{m}$. The above expression for $V_{1-\mathrm{loop}}$, thus, illustrates that the effective potential represents a power series in y^4, the coefficients of which are functions of SUSY-preserving mass parameters. From our results in Eqs. (A24) and (A27), we can read off the coefficient function $f_0(R_a)$ belonging to the leading-order contribution to $V_{1-\mathrm{loop}}$,

$$f_0(R_a) = \begin{cases} \sum_a \frac{1}{3}R_a^2(\varphi) + \mathcal{O}(x^4) & ; \quad x \ll 1 \\ \sum_a \ln R_a(\varphi) + \mathcal{O}(x^{-4}) & ; \quad x \gg 1 \end{cases}. \tag{A29}$$

As we are now going to show, this result for the function $f_0(R_a)$ can be reproduced by supplementing the tree-level Kähler potential, K_{tree}, by an effective one-loop Kähler potential, $K_{1-\mathrm{loop}}$,

$$K_{\mathrm{eff}} = K_{\mathrm{tree}} + K_{1-\mathrm{loop}} \tag{A30}$$

where K_{tree} and $K_{1-\mathrm{loop}}$ are respectively given as follows [24,101]:

$$K_{\mathrm{tree}} = \Phi^\dagger\Phi, \tag{A31}$$
$$K_{1-\mathrm{loop}} = -\frac{1}{32\pi^2}\mathrm{Tr}\left[\tilde{\mathcal{M}}^\dagger\tilde{\mathcal{M}}\left(\ln\left(\frac{\tilde{\mathcal{M}}^\dagger\tilde{\mathcal{M}}}{Q^2}\right) - c'\right)\right].$$

Here, $\tilde{\mathcal{M}}$ denotes the fermionic mass matrix, which is identical to the Hessian of the superpotential, and c' is a numerical constant that we will determine shortly. In global supersymmetry, the effective Kähler potential in Eq. (A30) results in the following F-term scalar potential for the Polonyi field $\phi = \varphi/\sqrt{2}e^{i\theta}$,

$$V_{\rm eff}^{\rm approx}(\varphi) = \left(\frac{\partial^2 K_{\rm eff}(\phi,\phi^*)}{\partial\phi\partial\phi^*}\right)^{-1}\left|\frac{\partial W_{\rm eff}(\phi)}{\partial\phi}\right|^2, \qquad \text{(A32)}$$

where the effective superpotential consists of the Polonyi tadpole term, $W_{\rm eff} \simeq \mu^2\Phi$, such that

$$\left|\frac{\partial W_{\rm eff}(\phi)}{\partial\phi}\right|^2 = \mu^4 = V_0. \qquad \text{(A33)}$$

It is important to note that the scalar potential in Eq. (A32) is only an approximation of the full effective potential. To see this, let us expand the inverse Kähler metric in powers of the loop factor $\kappa_\Phi^2/(16\pi^2)$,

$$\left(\frac{\partial^2 K_{\rm eff}(\phi,\phi^*)}{\partial\phi\partial\phi^*}\right)^{-1} = 1 - \frac{\partial^2 K_{1-{\rm loop}}(\phi,\phi^*)}{\partial\phi\partial\phi^*} + \mathcal{O}\left(\left(\frac{\kappa_\Phi^2}{16\pi^2}\right)^2\right). \qquad \text{(A34)}$$

The leading contribution to $V_{\rm eff}^{\rm approx}$ then yields what we shall refer to as the truncated effective potential,

$$V_{\rm eff}^{\rm trunc}(\varphi) = \left(1 - \frac{\partial^2 K_{1-{\rm loop}}(\phi,\phi^*)}{\partial\phi\partial\phi^*}\right)\left|\frac{\partial W_{\rm eff}(\phi)}{\partial\phi}\right|^2. \qquad \text{(A35)}$$

This truncated effective potential coincides with the full effective potential up to corrections of $\mathcal{O}(y^8)$,

$$V_0 + V_{1-{\rm loop}}(\varphi) = V_{\rm eff}^{\rm trunc}(\varphi) + \mathcal{O}(y^8). \qquad \text{(A36)}$$

An explicit proof of this relation is given by Intriligator, Seiberg, and Shih in Appendix A.5 of [43]. For an earlier discussion of the effective action in supersymmetric theories, see [102]. From the relation in Eq. (A36), it is evident that the truncated potential $V_{\rm eff}^{\rm trunc}$ is only a good approximation of the actual (true) effective potential, as long as the effect of spontaneous supersymmetry breaking is small, i.e., as long as $y \ll 1$. Otherwise, the $\mathcal{O}(y^8)$ corrections in Eq. (A36) become important. These radiative corrections—while present in the full expression for $V_{1-{\rm loop}}$ in Eq. (A1)—cannot be captured by the effective Kähler potential. One example for such a correction is, e.g., the shift in the tree-level vacuum energy density, ΔV_0 [see Eq. (A12)], which is proportional to $m^4 r_a^4 = m^8/m_a^4 \sim m_a^4 y^8$ and which cannot be explained in terms of an effective Kähler potential. Instead, the shift in the vacuum energy density ΔV_0 and all other $\mathcal{O}(y^8)$ corrections follow from an effective potential for the auxiliary component of the Polonyi multiplet, $\Phi|_{\theta^2} = F_\Phi$. This effective auxiliary field potential contains higher-order terms in F_Φ and is only negligible, once the effect of supersymmetry breaking is nothing but a small correction to the otherwise supersymmetric dynamics [102]. Fortunately, this is exactly the case in our model, since the soft SUSY-breaking mass scale m is always the smallest mass parameter in the IYIT sector [see Eq. (30)].

Let us now derive an explicit expression for the one-loop Kähler potential $K_{1-{\rm loop}}$. To do so, we first need to promote the fermionic masses in Eq. (A4) to functions of the chiral Polonyi superfield Φ,

$$\tilde{M}_a^2(\varphi;p) \to \frac{1}{2}(\mathcal{A}_a + p\mathcal{B}_a), \qquad \mathcal{A}_a = \kappa_\Phi^2\Phi^\dagger\Phi + 2m_a^2,$$
$$\mathcal{B}_a = [\kappa_\Phi^2\Phi^\dagger\Phi(\kappa_\Phi^2\Phi^\dagger\Phi + 4m_a^2)]^{1/2}. \qquad \text{(A37)}$$

These masses then need to be inserted into the general expression for $K_{1-{\rm loop}}$ in Eq. (A31),

$$K_{1-{\rm loop}} = -\frac{1}{32\pi^2}\sum_a\left[\mathcal{A}_a\left(\ln\left(\frac{m_a^2}{Q^2}\right) + c'\right) + \mathcal{B}_a\left(\frac{1}{2}\ln\left(1+\frac{\mathcal{B}_a}{\mathcal{A}_a}\right) - \frac{1}{2}\ln\left(1-\frac{\mathcal{B}_a}{\mathcal{A}_a}\right)\right)\right]. \qquad \text{(A38)}$$

Note that this result for $K_{1-{\rm loop}}$ is independent of the SUSY-breaking mass scale m. Put differently, the effective Kähler potential "does not know anything about supersymmetry breaking." It is, thus, clear that $V_{\rm eff}^{\rm trunc}$ (or $V_{\rm eff}^{\rm approx}$ for that purpose) can approximate the full effective potential only up to $\mathcal{O}(y^4)$. We shall now study $K_{1-{\rm loop}}$ as a function of $\mathcal{X}(\Phi^\dagger\Phi)$, the superfield analog of the order parameter $x(\varphi)$,

$$\mathcal{X}(\Phi^\dagger\Phi) = \left(\frac{\Phi^\dagger\Phi}{\varphi_c^2/2}\right)^{1/2}. \qquad \text{(A39)}$$

Close to the origin, $\mathcal{X}(\Phi^\dagger\Phi) \ll 1$, the effective one-loop Kähler potential consists of the following terms:

$$K_{1-{\rm loop}} = \Delta K_0 - N_X\frac{\kappa_\Phi^2}{32\pi^2}\left(\ln\left(\frac{\bar{m}^2}{Q^2}\right) + c' + 2\right)\Phi^\dagger\Phi + \mathcal{O}(\mathcal{X}^2), \qquad \text{(A40)}$$

where the first term, ΔK_0, denotes a constant shift in the VEV of Kähler potential that does not have any consequences in global supersymmetry and that is negligible in supergravity, $\Delta K_0 \ll M_{\rm Pl}^2$,

$$\Delta K_0 = -\frac{1}{16\pi^2}\sum_a m_a^2\left(\ln\left(\frac{m_a^2}{\bar{m}^2}\right) - 2\right). \qquad \text{(A41)}$$

Meanwhile, the second term in Eq. (A40) represents a shift in the canonical kinetic term of the Polonyi field. In the following, we will choose the renormalization scale Q and the constant c' as follows,

$$Q = \bar{m}, \qquad c' = -2, \tag{A42}$$

so that this term exactly vanishes (i.e., so that the Polonyi field Φ remains canonically normalized).

The first relevant correction to $K_{\rm tree}$ in the small-field regime, $\mathcal{X}(\Phi^\dagger\Phi) \lesssim 1$, is then of $\mathcal{O}(\mathcal{X}^2)$,

$$K_{1-\rm loop}^{\rm LE} = -\frac{1}{64\pi^2}\sum_a \frac{1}{3}\left(\frac{\kappa_\Phi^2 \Phi^\dagger\Phi}{m_a}\right)^2 + \mathcal{O}(\mathcal{X}^4). \tag{A43}$$

This form of the Kähler potential nicely reflects the fact, at small field values, the radiative corrections to the Polonyi potential follow from integrating out "heavy" meson fields with masses m_a. In terms of the order parameter superfield $\mathcal{X}(\Phi^\dagger\Phi)$, the one-loop Kähler potential in Eq. (A43) reads as follows:

$$K_{1-\rm loop}^{\rm LE} = -\frac{1}{4V_0}\Lambda_{\rm LE}^4 \mathcal{X}^2(\Phi^\dagger\Phi)\Phi^\dagger\Phi + \mathcal{O}(\mathcal{X}^4),$$
$$\Lambda_{\rm LE}^4 = \frac{m^4}{48\pi^2}\sum_a\left(\frac{\bar{m}}{m_a}\right)^2, \tag{A44}$$

where the expression for the effective potential energy scale $\Lambda_{\rm LE}$ is the identical to the one that we obtained in Sec. A 2 in the limit of small supersymmetry breaking, see Eq. (A24). In this sense, $K_{1-\rm loop}^{\rm LE}$ indeed reproduces the effective one-loop potential in the small-field regime up to corrections of $\mathcal{O}(y^8)$,

$$V_{\rm eff}^{\rm trunc}(\varphi) = V_0 + \Lambda_{\rm LE}^4 x^2(\varphi) + \mathcal{O}(x^4). \tag{A45}$$

Meanwhile, for large values of the Polonyi field, $\mathcal{X}(\Phi^\dagger\Phi) \gtrsim 1$, the one-loop Kähler potential reduces to

$$K_{1-\rm loop}^{\rm HE} = -N_X\frac{\kappa_\Phi^2}{32\pi^2}\left(\ln\left(\frac{\kappa_\Phi^2\Phi^\dagger\Phi}{\bar{m}^2}\right) - 2\right)\Phi^\dagger\Phi + \Delta K_0' + \mathcal{O}(\mathcal{X}^{-4}), \tag{A46}$$

which represents a logarithmic renormalization of the kinetic term. Here, $\Delta K_0'$ is a correction to the VEV of the Kähler potential, which is meaningless in global supersymmetry and irrelevant in supergravity,

$$\Delta K_0' = -\frac{1}{16\pi^2}\left(\ln\left(\frac{\kappa_\Phi^2\Phi^\dagger\Phi}{\bar{m}^2}\right) - 1\right)\sum_a m_a^2. \tag{A47}$$

Again, we may rewrite the one-loop Kähler potential in terms of the order parameter superfield $\mathcal{X}(\Phi^\dagger\Phi)$,

$$K_{1-\rm loop}^{\rm HE} = -\frac{1}{V_0}\Lambda_{\rm HE}^4(\ln\mathcal{X}(\Phi^\dagger\Phi) - 1)\Phi^\dagger\Phi + \mathcal{O}(\mathcal{X}^{-4}),$$
$$\Lambda_{\rm HE}^4 = N_X\frac{m^4}{16\pi^2}, \tag{A48}$$

which reproduces $V_{1-\rm loop}$ and the potential energy scale $\Lambda_{\rm HE}$ in the large-field regime, see Eq. (A26),

$$V_{\rm eff}^{\rm trunc}(\varphi) = V_0 + \Lambda_{\rm HE}^4\ln x(\varphi) + \mathcal{O}(x^{-4}). \tag{A49}$$

We, therefore, conclude that our results for $K_{1-\rm loop}^{\rm LE}$ and $K_{1-\rm loop}^{\rm HE}$ in Eqs. (A43) and (A46) indeed suffice to describe the radiative corrections to the Polonyi potential, as long as the effect of supersymmetry breaking remains small, $y \ll 1$. In the limit of large supersymmetry breaking, $y \to 1$, the full information on the radiative corrections is, however, only contained in the effective one-loop potential in Eq. (A1).

APPENDIX B: VIABILITY OF INFLATION IN THE CASE OF EARLY *R* SYMMETRY BREAKING

In the main text, we study inflation in an effective Polonyi model that is based on the effective superpotential $W_{\rm eff} \simeq \mu^2\Phi$, a near-canonical Kähler potential and logarithmic radiative corrections. In particular, we assume that the constant term in the superpotential, $W \supset w$, which allows one to tune the CC in the true vacuum to zero, is generated only after inflation. In this appendix, we are now going to show that (for all Kähler potentials of interest) this assumption is, indeed, inevitable. That is, in contrast to the main text, we are now going to assume that the constant w takes its low-energy value already from the very beginning, $w = w_0$. We then study the prospects of realizing successful inflation in the Polonyi model for different choices of the Kähler potential (near-canonical and approximately shift-symmetric), finding that, for $w = w_0$ during inflation, Polonyi inflation is always bound to fail. Either the Polonyi potential turns out to be too steep to support slow-roll inflation or there does not even exist a global Minkowski vacuum. Of course, our analysis does not represent a general no-go theorem, as some intricate choice for the Kähler potential may render inflation in the Polonyi model viable, after all [76].[23] But for our purposes, as we intend to focus on *simple* forms of the Kähler potential,

[23] In single-field helical-phase inflation, the phase of the Polonyi field, $\theta = \arg\phi$, may play the role of the inflaton [103]. This scenario requires a particular tuning of the Kähler potential, in order to stabilize the radial component of the complex Polonyi field, $\varphi = \sqrt{2}|\phi|$, and eventually results in inflationary predictions that are equivalent to those of natural inflation [54].

these results suffice to convince us that successful Polonyi inflation is better off if R symmetry is broken at late times.

1. Canonical Kähler potential plus higher-dimensional corrections

We first consider the full Polonyi superpotential in combination with a near-canonical Kähler potential,

$$W = \mu^2 \Phi + w,$$
$$K = \Phi^\dagger \Phi + \frac{\epsilon}{(2!)^2}\left(\frac{\Phi^\dagger \Phi}{M_{\rm Pl}}\right)^2 + \mathcal{O}(\epsilon^2, M_{\rm Pl}^{-4}),$$
$$\epsilon \lesssim 1. \tag{B1}$$

Apart from the subdominant higher-dimensional correction (parametrized in terms of the parameter ϵ), this is nothing but the original Polonyi model [68]. For now, let us assume that this model does not receive any radiative corrections from Yukawa couplings between Φ and some (heavy) matter fields. Given the above Kähler potential, the complex scalar contained in the chiral Polonyi field, $\tilde{\phi} \subset \Phi$,

$$\tilde{\phi} = \frac{\tilde{\varphi}}{\sqrt{2}} e^{i\tilde{\theta}}, \tag{B2}$$

is not canonically normalized. The Polonyi phase, $\tilde{\theta} = \arg \tilde{\phi}$, is always stabilized at $\tilde{\theta} = 0$, so that we can neglect it in the following. Meanwhile, the properly normalized radial component φ is given as follows:

$$\varphi = \int d\tilde{\varphi} \left(\frac{\partial^2 K}{\partial\tilde{\phi}\partial\tilde{\phi}^*}\right)^{1/2}$$
$$= \tilde{\varphi}\left[1 + \frac{\epsilon}{12}\left(\frac{\tilde{\varphi}}{M_{\rm Pl}}\right)^2 + \mathcal{O}(\epsilon^2)\right]. \tag{B3}$$

The vacuum energy density in the global minimum vanishes, once the constant w is fine-tuned to

$$w_0 = (2 - \sqrt{3})(1 - \epsilon + \mathcal{O}(\epsilon^2))\mu^2 M_{\rm Pl}. \tag{B4}$$

For this value of w, the real Polonyi field φ takes the following value in the true vacuum:

$$\langle\varphi\rangle = \sqrt{2}\left[(\sqrt{3} - 1)(1 + \epsilon) - \frac{\epsilon}{6} + \mathcal{O}(\epsilon^2)\right] M_{\rm Pl}. \tag{B5}$$

Around this vacuum, the Polonyi field has a tree-level mass of the order of the gravitino mass,

$$m_\varphi^2 = 2\sqrt{3}[1 - (3 - \sqrt{3})\epsilon + \mathcal{O}(\epsilon^2)] m_{3/2}^2, \tag{B6}$$

where the gravitino mass follows from inserting Eqs. (B4) and (B5) into the general formula in Eq. (72),

$$m_{3/2} = e^{2-\sqrt{3}}\left[1 + \frac{1}{2}(3 - \sqrt{3})\epsilon + \mathcal{O}(\epsilon^2)\right]\frac{\mu^2}{M_{\rm Pl}}. \tag{B7}$$

As it turns out, the scalar potential around the true vacuum is too steep to support slow-roll inflation. This can be seen by inspecting the slow-roll parameters ε and η in the vicinity of $\varphi = \langle\varphi\rangle$,

$$\varepsilon(\varphi) = 2\left[\frac{1}{\delta(\varphi)} + \frac{2\sqrt{3} - 1}{2\sqrt{2}} - \frac{2\sqrt{3} - 3}{4\sqrt{2}}\epsilon\right]^2$$
$$+ \frac{\sqrt{3} + 1}{\sqrt{3}} - \frac{5}{3}\epsilon + \mathcal{O}(\epsilon^2, \delta),$$
$$\eta(\varphi) = 2\left[\frac{1}{\delta(\varphi)} + \frac{2\sqrt{3} - 1}{\sqrt{2}} - \frac{2\sqrt{3} - 3}{2\sqrt{2}}\epsilon\right]^2$$
$$+ \frac{17}{2\sqrt{3}} - 4 - \frac{2}{3}(6\sqrt{3} - 5)\epsilon + \mathcal{O}(\epsilon^2, \delta), \tag{B8}$$

where $\delta(\varphi)$ parametrizes the displacement from the true vacuum in units of the Planck scale,

$$\delta(\varphi) = \frac{\varphi - \langle\varphi\rangle}{M_{\rm Pl}}. \tag{B9}$$

The expressions in Eq. (B8) imply that we only have a chance of sufficiently suppressing ε and η, if we displace the inflaton field by some amount of $\mathcal{O}(M_{\rm Pl})$ from its VEV. In the case of such large field displacements, the scalar potential is, however, severely steepened by SUGRA corrections. In fact, scanning over the entire Polonyi field range and evaluating the slow-roll parameters ε and η numerically, we find that nowhere in field space the scalar potential is sufficiently flat for slow-roll inflation,

$$\min\max\{\varepsilon, |\eta|\} \sim 0.3. \tag{B10}$$

In view of this result, we then also have little hope that radiative corrections may improve the situation. Instead, we expect that radiative correction would, overall, rather increase the magnitude of the slow-roll parameters ε and η even further. We, thus, arrive at the same conclusion as the analysis in [30], namely that inflation in the original Polonyi model, with $w = w_0$ during inflation, does not work.

2. Approximate shift symmetry along the real axis

Next, we assume that the Kähler potential exhibits an approximate shift symmetry along the real axis,

$$W = \mu^2\Phi + w,$$
$$K = \frac{\epsilon}{2}(\Phi + \Phi^\dagger)^2 - \frac{1}{2}(\Phi - \Phi^\dagger)^2 + \mathcal{O}(\epsilon^2, M_{\rm Pl}^{-2}), \quad \epsilon \ll 1. \tag{B11}$$

The complex scalar contained in the chiral Polonyi field, $\tilde{\phi} \subset \Phi$, then decomposes as follows:

$$\tilde{\phi} = \frac{1}{\sqrt{2}}(\tilde{\varphi} + i\tilde{\sigma}), \tag{B12}$$

where $\tilde{\varphi}$ and $\tilde{\sigma}$ denote the (not canonically normalized) inflaton field and its scalar partner (i.e., the "sinflaton"), respectively. First, we note that, in the limit of an exact shift symmetry, this theory does not admit a Minkowski vacuum. For $\epsilon = 0$, the only stationary point of the scalar potential is located at

$$\langle\tilde{\varphi}\rangle = -\sqrt{2}\frac{w}{\mu^2}, \qquad \langle\tilde{\sigma}\rangle = 0, \tag{B13}$$

so that the Polonyi tadpole term "eats up" the constant w. At this stationary point, we, therefore, have $\langle W\rangle = 0$, which implies unbroken R symmetry and, hence, zero gravitino mass. But more importantly, this point in field space corresponds to a dS state, as the vacuum energy density does not vanish, $\langle V\rangle = \mu^4$. On top of that, it is not even stable, as the $\tilde{\varphi}$ direction acquires a tachyonic mass, $m_\varphi^2 = -3\mu^4/M_{\rm Pl}^2$.

We can remedy this situation and distort the scalar potential by slightly breaking the shift symmetry in the $\tilde{\varphi}$ direction, i.e., by allowing for small nonzero values of ϵ. This stabilizes the stationary point in Eq. (B13) and allows for the possibility of a global Minkowski vacuum. In addition, a slightly broken shift symmetry in the Kähler potential appears to be more natural, anyway, as the superpotential already breaks the shift symmetry explicitly. The canonically normalized fields φ and σ are then given as follows:

$$\varphi = (1+\epsilon)^{1/2}\tilde{\varphi}, \qquad \sigma = (1+\epsilon)^{1/2}\tilde{\sigma}. \tag{B14}$$

For $\epsilon \neq 0$, the model exhibits a global Minkowski vacuum at a super-Planckian value of the inflaton field,

$$\langle\varphi\rangle = \frac{1+\epsilon}{\epsilon}\left[\frac{\sqrt{3}}{\sqrt{2}} - \left(\frac{\epsilon}{1+\epsilon}\right)^{1/2}\right]M_{\rm Pl}, \quad \langle\sigma\rangle = 0. \tag{B15}$$

We note that this solution is only consistent as long as ϵ is positive. Moreover, it becomes unphysical in the limit of an exact shift symmetry, i.e., for $\epsilon \to 0$. This is also the reason why, in contrast to Sec. B 1, we now refrain from expanding our results in powers of the small parameter ϵ. In order to tune the vacuum energy density in the vacuum to zero, the constant w needs to take the following value:

$$w_0 = -\frac{1}{\sqrt{2}\epsilon}\left[\frac{\sqrt{3}}{\sqrt{2}}(1+\epsilon)^{1/2} - 2\epsilon^{1/2}\right]\mu^2 M_{\rm Pl}, \tag{B16}$$

which also diverges for $\epsilon \to 0$. As we are able to read off from Eqs. (B15) in (B16), the relation between $\langle\tilde{\varphi}\rangle$ and w in Eq. (B13) now no longer applies. Thanks to the slightly broken shift symmetry, we now have $\langle\tilde{\varphi}\rangle = (1+\epsilon)^{-1/2}\langle\varphi\rangle = -\sqrt{2}w/\mu^2 + \epsilon^{-1/2}$, which results in a nonzero gravitino mass,

$$m_{3/2} = \frac{1}{\sqrt{2}\epsilon^{1/2}}\exp\left[\frac{\epsilon}{2(1+\epsilon)}\frac{\langle\varphi\rangle^2}{M_{\rm Pl}^2}\right]\frac{\mu^2}{M_{\rm Pl}}. \tag{B17}$$

The field φ then acquires a nontachyonic mass of the order of the gravitino mass,

$$m_\varphi^2 = 2\sqrt{3}\left(\frac{2\epsilon}{1+\epsilon}\right)^{3/2} m_{3/2}^2. \tag{B18}$$

Again, the scalar potential around the true vacuum is too steep for slow-roll inflation. In the vicinity of $\varphi = \langle\varphi\rangle$, the slow-roll parameters ε and η [as functions of $\delta(\varphi) = (\varphi - \langle\varphi\rangle)/M_{\rm Pl}^2$] are now given as

$$\begin{aligned}\varepsilon(\varphi) &= 2\left[\frac{1}{\delta(\varphi)} + \frac{\sqrt{3}}{\sqrt{2}} - \frac{\epsilon^{1/2}}{2}\right]^2 + \mathcal{O}(\epsilon),\\ \eta(\varphi) &= 2\left[\frac{1}{\delta(\varphi)} + \sqrt{6} - \epsilon^{1/2}\right]^2\\ &\quad - \sqrt{6}(\sqrt{6} - 2\epsilon^{1/2}) + \mathcal{O}(\epsilon).\end{aligned} \tag{B19}$$

A numerical scan of the full expressions for ε and η over the entire field range reveals that both parameters are always of $\mathcal{O}(1)$ or larger and, hence, always too large for slow-roll inflation. Similarly as in the case of a near-canonical Kähler potential, we do not expect that radiative corrections could improve this situation. We, therefore, conclude that, assuming an approximate shift symmetry along the real axis in the Kähler potential, inflation based on the full Polonyi superpotential is unfortunately not an option.

3. Approximate shift symmetry along the imaginary axis

Finally, let us examine what happens if we replace the approximate shift symmetry along the real axis by an approximate shift symmetry along the imaginary axis. As we will see shortly, in this case, we find a stationary point at $\langle\varphi\rangle = 0$, similarly as in Sec. II D. For that reason, our analysis in this section will resemble the discussion in Sec. II D much more closely than our analysis in the previous section. But, first things first. Now, we supplement the Polonyi potential by a Kähler potential of the following form:

$$\begin{aligned}W &= \mu^2\Phi + w,\\ K &= \frac{1}{2}(\Phi + \Phi^\dagger)^2 - \frac{\epsilon}{2}(\Phi - \Phi^\dagger)^2 + \mathcal{O}(\epsilon^2, M_{\rm Pl}^{-2}).\end{aligned} \tag{B20}$$

The (not canonically normalized) inflaton field $\tilde{\varphi}$ and the (not canonically normalized) sinflaton field $\tilde{\sigma}$ then correspond to the imaginary and the real part of the complex scalar $\tilde{\phi} \subset \Phi$, respectively,

$$\tilde{\phi} = \frac{1}{\sqrt{2}}(\tilde{\sigma} + i\tilde{\varphi}), \tag{B21}$$

As in the previous section, the canonically normalized fields φ and σ are related to $\tilde{\varphi}$ and $\tilde{\sigma}$ as follows:

$$\varphi = (1+\epsilon)^{1/2}\tilde{\varphi}, \qquad \sigma = (1+\epsilon)^{1/2}\tilde{\sigma}. \tag{B22}$$

Now we find a stationary point at a Planckian value of the sinflaton field rather than the inflaton field,

$$\langle\varphi\rangle = 0, \quad \langle\sigma\rangle = (1+\epsilon)^{1/2}\left[1 - \frac{\sqrt{3}}{\sqrt{2}}(1+\epsilon)^{1/2}\right]M_{\rm Pl}. \tag{B23}$$

The vacuum energy density at this stationary point vanishes, once w takes the following value,

$$w_0 = -\left[\sqrt{2} - \frac{\sqrt{3}}{2}(1+\epsilon)^{1/2}\right]\mu^2 M_{\rm Pl}. \tag{B24}$$

We then obtain for the gravitino mass as well as for the inflaton mass around this stationary point,

$$m_{3/2} = \frac{1}{\sqrt{2}}\exp\left[\frac{1}{2(1+\epsilon)}\frac{\langle\sigma\rangle^2}{M_{\rm Pl}^2}\right]\frac{\mu^2}{M_{\rm Pl}},$$
$$m_\varphi^2 = \left[4 - 2\sqrt{3}\left(\frac{2}{1+\epsilon}\right)^{3/2}\right]m_{3/2}^2. \tag{B25}$$

None of the above quantities diverges in the limit of an exact shift symmetry. In contrast to the previous section, the limit $\epsilon \to 0$, therefore, does not take us out of the physical regime. At the same time, independent of whether ϵ is exactly zero or just small, $\epsilon \ll 1$, we find that the inflaton direction in field space turns out to be tachyonic, $m_\varphi^2 < 0$. The stationary point at $\varphi = 0$, therefore, does not represent the global minimum of the scalar potential. Instead, the scalar potential is either unbounded from below (for $\epsilon = 0$) or it exhibits a global AdS vacuum (for $\epsilon \neq 0$) at a large inflaton field value $\varphi_{\rm min} \propto M_{\rm Pl}/\epsilon$. It is obvious that the former case does not represent a viable setting for inflation. For this reason, we shall focus on the latter case in the following and argue that in this case, too, inflation cannot be successfully realized. Here, our argument will closely follow the discussion in Sec. II D. First, let us consider the scalar potential for the inflaton field φ with the sinflaton field σ being fixed at $\sigma = \langle\sigma\rangle$,

$$V(\varphi) = \exp\left[\frac{\epsilon}{1+\epsilon}\frac{\varphi^2}{M_{\rm Pl}^2}\right]\left(\frac{1}{2}m_\varphi^2\varphi^2 + \frac{\lambda_\varphi}{4!}\varphi^4\right),$$
$$\lambda_\varphi = \frac{48\epsilon^2}{(1+\epsilon)^3}\frac{m_{3/2}^2}{M_{\rm Pl}^2}. \tag{B26}$$

We assume that the shift symmetry the Kähler potential in Eq. (B20) is only slightly broken, $\epsilon \ll 6^{1/3} - 1$, so that $m_\varphi^2 < 0$, and we assume that the potential is not unbounded from below, $\epsilon \neq 0$, so that $\lambda_\varphi > 0$. Under these assumptions, the inflaton potential exhibits a global AdS vacuum at

$$\varphi_{\rm min} = \left(\frac{1+\epsilon}{\epsilon}\right)^{1/2}[(1+a^2)^{1/2} - (1-a)]^{1/2}M_{\rm Pl},$$
$$a = -\frac{\epsilon}{1+\epsilon}\frac{6}{\lambda_\varphi}\frac{m_\varphi^2}{M_{\rm Pl}^2}. \tag{B27}$$

At this field value, the scalar potential takes the following value:

$$V_{\rm min} = -\exp[(1+a^2)^{1/2} - (1-a)][(1+a^2)^{1/2} - 1]$$
$$\times \frac{4}{1+\epsilon}m_{3/2}^2 M_{\rm Pl}^2, \tag{B28}$$

which represents the vacuum energy density in the global AdS minimum. Similarly as in Sec. II D, we may hope that there might be a chance to lift this AdS vacuum by radiative corrections. But as it turns out, this attempt to stabilize the inflaton potential fails for the same reason as in Sec. II D. Within the parameter ranges of interest, $|\epsilon| \lesssim 10^{-0.5}$ and $\lambda/\eta \lesssim 10^{-0.5}$, we always find a hierarchy among the critical field value φ_c and the position of the AdS vacuum, $\varphi_c/\varphi_{\rm min} \lesssim 0.1$. The quadratic radiative corrections close to the origin, therefore, fail to stabilize the inflaton potential all the way up to $\varphi = \varphi_{\rm min}$. In consequence of that, we always end up with a local minimum in between φ_c and $\varphi_{\rm min}$. Inflation taking place in the logarithmic part of the effective potential then always "gets stuck" and the inflaton has no chance of reaching the Minkowski vacuum at the origin (see also our discussion in Sec. II D).

This completes our argument that inflation based on the Polonyi superpotential $W = \mu^2\Phi + w$—with the constant w being set to its final value $w = w_0$ already from the very beginning—works neither for a near-canonical nor for an approximately shift-symmetric Kähler potential. It is this conclusion that leads us to resort to studying inflation in combination with late-time R symmetry breaking in the main text. And indeed, successful Polonyi inflation turns out to be feasible, once w vanishes during inflation.

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