Polonyi-Starobinsky supergravity with inflaton in a massive vector multiplet with DBI and FI terms

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Abstract

We propose the Starobinsky-type inflationary model in the matter-coupled $N = 1$ four-dimensional supergravity with the massive vector multiplet that has inflaton (scalaron) and goldstino amongst its field components, whose action includes the Dirac-Born-Infeld-type kinetic term and the generalized (new) Fayet-Iliopoulos-type term, without gauging the R-symmetry. The $N = 1$ chiral matter (“hidden sector”) is described by the modified Polonyi model needed for spontaneous supersymmetry breaking after inflation. We compute the bosonic action and the scalar potential of the model, and show that it can accommodate the positive (observed) cosmological constant (as the dark energy) and the spontaneous supersymmetry breaking at high scale after the Starobinsky inflation.
1 Introduction

Supergravity is well motivated in theoretical high-energy physics of elementary particles beyond the Standard Model, in the gravitational theory, and in superstring theory. Supergravity theory is also considered as the promising framework in theoretical cosmology, because it has the natural candidate for dark matter particle, known as the Lightest Supersymmetric Particle (LSP). It is also worth mentioning that the phenomenological supergravity model building, both for particle physics and cosmology, is highly non-trivial, because the supergravity interactions are very restrictive (due to local supersymmetry), and the low-energy bounds (given by the Standard Model of elementary particles and the $\Lambda\text{CDM}$ Model of Cosmology) are tight.

The standard approach to the supergravity-based cosmology and the related inflationary model building employs chiral superfields for matter, inflaton and goldstino. The inflaton field is the real scalar field driving inflation, whereas a chiral multiplet has a complex physical scalar. Hence, another (non-inflaton) real scalar has to be stabilized during a single-field inflation. Since Supersymmetry (SUSY) is spontaneously broken during inflation, there must be also the goldstino (fermionic) field that is usually assigned to a chiral multiplet too. Then the input is provided by a (non-holomorphic) Kähler potential and a (holomorphic) superpotential of the chiral superfields involved. The inflationary scalar potential in supergravity generically suffers from the so-called $\eta$-problem, so that the input should be carefully designed to avoid this problem — see e.g., Refs. [1, 2, 3, 4, 5, 6, 7] for many examples.

The viable alternative to the standard approach in the inflationary cosmology based on supergravity is possible by employing a massive vector multiplet that has only one (real) physical scalar to be identified with inflaton, and its fermionic superpartner to be identified with goldstino (there is the massive vector field too). This alternative approach in its most minimal form (without chiral matter) can accommodate cosmological inflation with any values of the Cosmic Microwave Background (CMB) radiation tilts $n_s$ and $r$ [8, 9, 10], but fails to achieve a positive cosmological constant (dark energy) and spontaneously broken SUSY after inflation. Phenomenological applications of supergravity to particle physics (reheating) after inflation demand adding the ”hidden sector” to be responsible for spontaneous SUSY breaking and, next, a mediation of the SUSY breaking from the hidden sector to the observable (low-energy) phenomena described by the Standard Model. The simplest candidate for the hidden sector is given by the so-called Polonyi model of a single chiral superfield with a linear superpotential [11]. The Polonyi model was employed for SUSY breaking in the supergravity model with inflaton in the massive vector multiplet in Refs. [12, 13].
Starobinsky inflation [14] can be also realized in the supergravity with inflaton (scalaron) in the massive vector multiplet, without chiral matter [8, 9]. However, in the presence of the Polonyi superfield, it was found to lead to instability of the inflationary trajectory [15]. A cure to the last problem was also proposed in Ref. [15] by modifying the embedding of the Starobinsky model into supergravity. The alternative possibility was proposed in Ref. [16] by removing the Polonyi superfield, but adding the generalized Fayet-Iliopoulos (PI) term instead, which does not require gauging the R-symmetry. Such new FI terms in supergravity were recently introduced in Refs. [17, 18, 19, 20]. Then it is possible to get a positive cosmological constant via the D-type spontaneous SUSY breaking by the use of the FI term [19]. However, because of the tiny observed value of the cosmological constant, the scale of such SUSY breaking appears to be very small and, hence, inappropriate for particle physics phenomenology.

In this paper we add the Polonyi superfield to the supergravity model of Ref. [16], in order to achieve a higher scale of spontaneous SUSY breaking via the F-type SUSY breaking, while also keeping the D-type SUSY breaking due to the FI term describing the observed positive cosmological constant. Our construction appears to be very delicate and rather complicated, and is based on the following theoretical resources (tools):

- the manifest (linearly realized) local $N = 1$ supersymmetry,
- the inflaton (scalaron) and the goldstino in the massive $N = 1$ vector multiplet,
- the vector multiplet kinetic terms of the Dirac-Born-Infeld (DBI) structure inspired by superstrings and D-branes [21, 22, 23],
- the generalized Fayet-Iliopoulos (FI) term, without gauging the R-symmetry [17, 18, 19, 20],
- the Kähler potential with the stabilizing (quartic) term, and the modified kinetic term of the Polonyi superfield,
- the Starobinsky (real) master function modified by a linear term in supergravity.

Our motivation and reasons for these modifications of the Polonyi-Starobinsky (PS) supergravity of Refs. [15, 16] are explained in the main text of this paper.

1The standard FI term in supergravity is known to be very restrictive, that is problematic for cosmological applications [24].

Our paper is organized as follows. The technical setup based on the superconformal tensor calculus with a chiral compensator is briefly reviewed in Sec. 2. In Sec. 3 we define our supersymmetric action for the vector multiplet, and calculate its bosonic terms, including the auxiliary field and the scalar potential. In Sec. 4 we introduce the modified Polonyi model in the context of PS supergravity, and study its properties during and after inflation. In Sec. 5 we propose a more general action in the curved superspace of the "old-minimal" supergravity, connect it to the actions defined in Secs. 2 and 3, and use the superfield formulation to verify our results found in the superconformal tensor calculus approach. Our conclusion is given by Sec. 6. Technical details are collected in Appendices A, B, C and D.

2 Superconformal tensor calculus

We use the conformal \( N = 1 \) supergravity techniques \cite{25, 26, 27, 28, 29}, and follow the notation and conventions of Ref. \cite{30}. In addition to the symmetries of Poincaré supergravity, one also has the gauge invariance under dilatations, conformal boosts and \( S \)-supersymmetry, as well as under \( U(1)_A \) rotations. The gauge fields of dilatations and \( U(1)_A \) rotations are denoted by \( b_\mu \) and \( A_\mu \), respectively. A multiplet of conformal supergravity has charges with respect to dilatations and \( U(1)_A \) rotations, called Weyl and chiral weights, respectively, which are denoted by the pair (Weyl weight, chiral weight) in what follows.

A chiral multiplet has field components

\[
S = \{ S, P_{L\chi}, F \},
\]

where \( S \) and \( F \) are complex scalars, and \( P_{L\chi} \) is a left-handed Weyl fermion (\( P_L \) is the chiral projection operator). In this paper, we use two types of chiral multiplets: the conformal compensator \( S_0 \) and the matter multiplets \( S^i \), where the index \( i = 1, 2, 3, \ldots \), counts the matter multiplets. The \( S_0 \) has the weights \((1, 1)\) and is used to fix some of the superconformal symmetries. The matter multiplets \( S^i \) have the weights \((0, 0)\). The anti-chiral multiplets are denoted by \( \bar{S}_0 \) and \( \bar{S}^i \).

As regards a general (complex or real) multiplet, it has the field content

\[
\Phi = \{ C, Z, \mathcal{H}, \kappa, \mathcal{B}_a, \Lambda, D \},
\]

where \( Z \) and \( \Lambda \) are fermions, \( \mathcal{B}_a \) is a (complex or real) vector, and others are (complex or real) scalars, respectively.

\(^2\)Though the superspace approach and the superconformal tensor calculus approach to supergravity are equivalent, their correspondence is non-trivial.
The (gauge) field strength multiplet \( W \) has the weights \((3/2, 3/2)\) and the following field components:

\[
\bar{\eta} W = \left\{ \bar{\eta} P_L \lambda, \frac{1}{\sqrt{2}} \left( -\frac{1}{2} \gamma_{ab} \tilde{F}^{ab} + i D \right) P_L \bar{\eta}, \bar{\eta} P_L \not{D} \phi \right\},
\]

where \( \eta \) is the dummy spinor, \( \tilde{F}_{ab} = \partial_a B_b - \partial_b B_a + \bar{\psi}_{[a} \gamma_{b]} \lambda \equiv F_{ab} + \bar{\psi}_{[a} \gamma_{b]} \lambda \) is the superconformally covariant field strength, the \( \psi_a \) is gravitino, the \( \lambda \) and \( D \) are Majorana fermion and the real auxiliary scalar, respectively. The related expressions of the multiplets \( W^2 \) and \( W^2 \bar{W}^2 \), which are embedded into the chiral multiplet (1) and the general multiplet (2), respectively, are

\[
W^2 = \left\{ \cdots, \cdots, \cdots + \frac{1}{2} (F F - F \tilde{F}) - D^2 \right\},
\]

\[
W^2 \bar{W}^2 = \left\{ \cdots, \cdots, \cdots, \cdots, \cdots, \cdots + \frac{1}{2} |(F F - F \tilde{F}) - 2 D^2|^2 \right\},
\]

where we have omitted the fermionic terms (denoted by dots) for simplicity. In addition, we use the book-keeping notation \( F F = F_{ab} F^{ab} \) and \( \tilde{F}^{ab} \equiv -\frac{1}{2} \epsilon^{abcd} F_{cd} \) throughout the paper.

We also need another chiral multiplet

\[
\Sigma \left( \bar{W}^2 / |S_0|^4 \right) = \left\{ \cdots, \cdots, \cdots, \frac{F_0}{|S_0|^4} (F F + F \tilde{F} - 2 D^2) + \cdots \right\},
\]

where \( \Sigma \) is the chiral projection operator \([28, 29]\). The argument of \( \Sigma \) requires specific Weyl and chiral weights: in order for \( \Sigma \Phi \) to make sense, \( \Phi \) must satisfy \( w - n = 2 \), where \( (w, n) \) are the Weyl and chiral weights of \( \Phi \). We adjust the correct weights of the argument, by inserting the factor \( |S_0|^4 \). Equation (6) is the conformal supergravity counterpart of the superfield \( D^2 W^2 \).

The covariant derivative of \( W \) is given by \([29]\)

\[
D W = \{ -2D, \cdots, \cdots, \cdots, \cdots, \cdots \}
\]

of the weights \((2, 0)\). Here, the dots in the higher components also include some bosonic terms, but we do not write down them here for simplicity (see Ref. \([17]\) for their explicit expressions).

A massive vector multiplet \( V \) has the field components

\[
V = \{ C, Z, H, K, B_a, \lambda, D \} ,
\]
while all of them are either real (bosonic) or Majorana (fermionic). The weights of $V$ are $(0, 0)$.

The bosonic parts of the F-term invariant action formulas are

$$[S]_F = \int d^4x \sqrt{-g} \frac{1}{2} (F + \bar{F}),$$

while they can be applied only when $S$ has the weights $(3, 3)$. The bosonic part of the D-term formula for a real multiplet $\phi$ of the weights $(2, 0)$ is

$$[\phi]_D = \int d^4x \sqrt{-g} \left( D_\phi - \frac{1}{3} C_\phi R(\omega) \right),$$

where $R(\omega)$ is the superconformal Ricci scalar in terms of spacetime metric and $b_\mu$ [30]. The $C_\phi$ and $D_\phi$ are the first and the last components of $\phi$, respectively.

### 3 Vector multiplet coupled to chiral matter

Let us consider a (chiral) matter coupled extension of the DBI and FI system investigated in Ref. [16], with the action

$$S = S_{V,M} + S_{DBI} + S_{FI},$$

where we have introduced

$$S_{V,M} = \left[ |S_0|^2 \mathcal{H} \right]_D + 2 |S_0|^2 |\mathcal{W}|_F,$$

$$S_{DBI} = -\frac{1}{2} |\mathcal{W}|_F + \left[ |S_0|^4 \left( 1 - 2\alpha \mathcal{A} + \sqrt{1 - 4\alpha \mathcal{A}} + 4\alpha^2 \mathcal{B}^2 \right) \right]_D,$$

$$S_{FI} = \left[ |S_0|^2 \mathcal{H} \mathcal{I} \frac{W^2 \bar{W}^2}{(DW)^2(D\bar{W})^2} D\mathcal{W} \right]_D.$$

Here, $\mathcal{H}, \mathcal{I}$, and $\alpha$ are arbitrary weightless real functions of $V, S^i$, and $\bar{S}$, the $\mathcal{W}$ is a holomorphic superpotential that depends on $S^i$ only. The $\mathcal{A}$ and $\mathcal{B}$ in $S_{DBI}$ are given by

$$\mathcal{A} = \bar{\Sigma} \left( \frac{W^2}{|S_0|^4} \right) + \text{h.c.}, \quad \mathcal{B} = \bar{\Sigma} \left( \frac{W^2}{|S_0|^4} \right) - \text{h.c.}$$

More general actions are defined in Sec. 5 and Appendix A.

The action [11] reduces to the action of Ref. [12] in the limit of the vanishing $\alpha$ and $\mathcal{I}$. It is remarkable that one can introduce a superpotential
without any restriction because the new FI term does not require the gauged R-symmetry. Under Kähler transformations, the $S_0$, $\mathcal{H}$, and $\mathcal{W}$ transform as

$$S_0 \rightarrow S_0 e^Y, \quad \mathcal{H} \rightarrow \mathcal{H} e^{-Y}, \quad \mathcal{W} \rightarrow \mathcal{W} e^{-3Y}, \quad (16)$$

where $Y$ is a holomorphic function of $S^i$. That is why we inserted the factor $\mathcal{H}$ in Eq. (14) for the Kähler invariance, as was pointed out in Refs. [19, 20].

To obtain the Lagrangian in field components, we have to eliminate the unwanted symmetries of the conformal supergravity by gauge fixing, and solve for the auxiliary fields by their (algebraic) equations of motion. These steps can be done separately in Eqs. (12)-(14), except for the integration of the auxiliary field $D$. Let us focus on the $S_{V,M}$ first, and demonstrate the result for the bosonic part only. The details of derivation are given in Appendix B. The resulting Lagrangian $\mathcal{L}_{V,M}$ reads

$$\mathcal{L}_{V,M} = \frac{1}{2} \mathcal{R} - \frac{1}{2} \mathcal{J}_{CC} (\partial a C)^2 - \frac{1}{2} \mathcal{J}_{CC} B_a^2 + \mathcal{J}_C D$$

$$- 2 \mathcal{J}_{ij} \partial_a S^i \partial_a \bar{S}^j + (- \mathcal{J}_{C} \partial_a C \bar{S}^i - i \mathcal{J}_{C} \bar{B}_a \partial_a S^i + \text{h.c.}) - V_F,$$

where we have defined

$$\mathcal{J} = - \frac{3}{2} \log \left( - \frac{2}{3} \mathcal{H} \right), \quad (18)$$

and the subscripts on $\mathcal{J}$ and $\mathcal{W}$ denote the derivatives with respect to $C$ and $S^i(\bar{S}^i)$. The $V_F$ is the F-type scalar potential,

$$V_F = e^{2\mathcal{J}} \left[ 2 \Delta \bar{i} \left\{ \frac{1}{2} \mathcal{W}_i + \left( \mathcal{J}_i - \frac{\mathcal{J}_{C} \mathcal{J}_{C}}{\mathcal{J}_{CC}} \right) \mathcal{W} \right\} \left\{ \frac{1}{2} \bar{\mathcal{W}}_j + \left( \mathcal{J}_j - \frac{\mathcal{J}_{C} \mathcal{J}_{C}}{\mathcal{J}_{CC}} \right) \bar{\mathcal{W}} \right\} \right.$$

$$+ 2 \frac{\mathcal{J}_{C}^2}{\mathcal{J}_{CC}} |\mathcal{W}|^2 - 3 |\mathcal{W}|^2 \right], \quad (19)$$

where we have used the notation

$$\Delta \bar{i} = (\Delta \bar{i})^{-1}, \quad \Delta \bar{j} = \mathcal{J}_{ij} - \frac{\mathcal{J}_{C} \mathcal{J}_{C}}{\mathcal{J}_{CC}}.$$

Under the superconformal gauge fixing conditions (109)-(111), the $S_{DBI}$ becomes (see also Appendix A)

$$\mathcal{L}_{DBI} = \frac{e^{4\mathcal{J}/3}}{8\alpha} \left[ 1 - \sqrt{1 - 8\alpha e^{-4\mathcal{J}/3} \left( D^2 - \frac{1}{2} \mathcal{F} \bar{\mathcal{F}} \right) + 4\alpha^2 e^{-8\mathcal{J}/3} (\mathcal{F} \bar{\mathcal{F}})^2} \right]. \quad (21)$$

The Lagrangian density is defined by $S = \int d^4x \sqrt{-g} \mathcal{L}.$
The bosonic part of $\mathcal{L}_{FI}$ takes the same form as that in Ref. [16], with

$$\mathcal{L}_{FI} = \mathcal{I} \left( \frac{D^2 - \frac{1}{2} FF - \frac{1}{4} (F \tilde{F})^2}{6D^3} \right),$$

(22)

though with the matter-dependent $\mathcal{I}$-function in general.

Having obtained the total Lagrangian $\mathcal{L} = \mathcal{L}_{V,M} + \mathcal{L}_{DBI} + \mathcal{L}_{FI}$ before integrating out $D$, we can discuss the elimination of $D$. The $D$ equation of motion gives

$$\mathcal{J}_C + \frac{1}{\sqrt{F^2 - 8\alpha e^{-4\mathcal{J}/3} D^2}} D + \frac{\mathcal{I}}{6} \left( 1 + \frac{FF}{D^2} - \frac{3(FF)^2 - (F \tilde{F})^2}{4D^4} \right) = 0,$$

(23)

where we have used the notation

$$\mathcal{F}(F) \equiv \sqrt{1 + 4\alpha e^{-4\mathcal{J}/3} FF + 4\alpha^2 e^{-8\mathcal{J}/3} (F \tilde{F})^2}.$$ 

(24)

Similarly to Ref. [16], we find a perturbative solution with respect to the FI term,

$$D = D^{(0)} + \mathcal{I} D^{(1)} + \mathcal{O}(\mathcal{I}^2),$$

(25)

where we have introduced

$$D^{(0)} = K \mathcal{F},$$

(26)

$$D^{(1)} = -\frac{K^3}{6\mathcal{J}_C^2} \mathcal{F} \left( 1 + \frac{FF}{K^2 F^2} - \frac{3(FF)^2 - (F \tilde{F})^2}{4K^4 F^4} \right),$$

(27)

and

$$K^2 \equiv \frac{\mathcal{J}_C^2}{1 + 8\alpha \mathcal{J}_C^2 e^{-4\mathcal{J}/3}}.$$ 

(28)

Substituting the solution [25] back into the Lagrangian, we obtain the Lagrangian in the first order with respect to $\mathcal{I}$ as follows:

$$\mathcal{L} = \frac{1}{2} R - \frac{1}{2} \mathcal{J}_{CC} (\partial_a C)^2 - \frac{1}{2} \mathcal{J}_{BB} B_a^2 - 2 \mathcal{J}_{ij} \partial_a S^i \partial^a \bar{S}^j$$

$$+ \left( - \mathcal{J}_{C} \partial_a C \partial^a S^i - i \mathcal{J}_{C} B_a \partial^a S^i + \text{h.c.} \right) - V_F$$

$$+ \frac{e^{4\mathcal{J}/3}}{8\alpha} \left( 1 \pm \frac{\mathcal{J}_C}{K \mathcal{F}} \right) \pm \frac{\mathcal{I}}{6} K \mathcal{F} \left( 1 - \frac{FF}{K^2 F^2} + \frac{(FF)^2 - (F \tilde{F})^2}{4K^4 F^4} \right) + \mathcal{O}(\mathcal{I}^2).$$
When $B_a = 0$, Eq. (23) can be solved exactly, and its solution is given by

$$D = \sqrt{\frac{(J_C + \frac{I}{6})^2}{1 + 8\alpha e^{-4J/3} (J_C + \frac{I}{6})^2}}. \quad (30)$$

Thus, the Lagrangian becomes

$$\mathcal{L} = \frac{1}{2} R - \frac{1}{2} J_{CC}(\partial_a C)^2 - 2 J_{ij}\partial_a S^i \partial^a \tilde{S}^j$$

$$+ \left( - J_{Ci}\partial_a C \partial^a S^i + \text{h.c.} \right) - V_F - V_D, \quad (31)$$

where we have

$$V_D = -\frac{e^{4J/3}}{8\alpha} \left( 1 - \sqrt{1 + 8\alpha e^{-4J/3} (J_C + \frac{I}{6})^2} \right). \quad (32)$$

The field dependent kinetic term of the vector field $B_a$ can be read off from Eq. (30), and it should be negative to avoid a ghost mode, i.e.

$$-\frac{1}{4K} \left( J_C - \frac{2}{3} I + \frac{4}{3} \alpha K \ell e^{-4J/3} \right) < 0. \quad (33)$$

4 Inflation and SUSY breaking

Let us apply the supergravity model constructed in the previous section to a description of cosmological inflation, spontaneous SUSY breaking after inflation, and dark energy (positive cosmological constant). In this section we restore the (reduced) Planck mass $M_P$ and the gauge coupling constant $g$. We specify our model for the Starobinsky-type inflation by identifying the real scalar $C$ of the massive vector multiplet with the inflaton (Starobinsky scalaron). We also add the SUSY breaking (hidden) sector $\Phi$ as the modified Polonyi model, with

$$\mathcal{J} = J(C) + \frac{1}{2} |\Phi|^2 + \zeta |\Phi|^4 + \epsilon C(\Phi + \bar{\Phi}), \quad (34)$$

$$\mathcal{W} = \mu (\Phi + \beta), \quad (35)$$

$^4$The vector multiplet gets its mass via Higgs effect, see Sec. 5 for more.
where the function $J$ is taken in the Starobinsky form [9, 10] modified by the linear term with the real coefficient $\gamma$ as follows:

$$J = -\frac{3}{2} M_P^2 \log \left( -\frac{C}{M_P} e^{C/M_P} \right) + \gamma M_P C. \quad (36)$$

Some comments are in order. The original Polonyi model [11] is obtained in the case of the vanishing parameters $\zeta$, $\epsilon$ and $\gamma$, with the $\mu$ and $\beta$ as the real parameters. We have added the quartic coupling with the $\zeta$ parameter to the Kähler potential, and the mixing of the kinetic terms of $C$ and $\Phi$ with the parameter $\epsilon$ in (34). As regards Eq. (36), we have also added the linear term in $C$. As is demonstrated below, all these modifications are necessary to achieve our goals given in the Abstract and Sec. 1.

The $F$- and $D$- type scalar potentials are given by

$$V_F = e^{2J/M_P^2} \mu^2 \left[ \frac{1}{\Delta} \left| 1 + \frac{1}{M_P^2} (\Phi + \beta)(\Phi + 4\zeta \Phi |\Phi|^2 + 2\epsilon C - 2\epsilon Q) \right|^2 - \frac{3}{M_P^2} |\Phi + \beta|^2 \right], \quad (37)$$

$$V_D = \frac{M_P^4}{a^4} \left( \sqrt{1 + a^4 \left( \frac{gJ_C}{M_P} + \frac{b}{6} \right)^2} - 1 \right), \quad (38)$$

where we have defined

$$\Delta = 1 + 8\zeta |\Phi|^2 - 2\frac{\epsilon}{J_{CC}}, \quad (39)$$

$$P = 1 - \frac{2}{3 M_P^2} \frac{J_C^2}{J_{CC}}, \quad (40)$$

$$Q = \frac{J_C}{J_{CC}}, \quad (41)$$

with the dimensionless parameters $a$ and $b$ according to

$$\alpha = \frac{a^4 e^{4J/3M_P^2}}{8 M_P^4} \quad \text{and} \quad \mathcal{I} = b M_P^2. \quad (42)$$

The kinetic term of the inflaton $C$ is given by

$$-\frac{3 M_P^2}{4 C^2} (\partial_a C)^2. \quad (43)$$
It should be noticed that the $\gamma$-modification above does not affect this kinetic term. Hence, we define the "canonical" inflaton (scalarm) as

$$C/M_P = -e^{\sqrt{2} x}.$$  \hfill (44)

4.1 During inflation

In this subsection, we discuss the stabilization of $\Phi$ during inflation. Expanding the scalar potential around $\Phi = 0$ up to the second order, we obtain

$$V_F = V_0 + \delta_R \Phi_R + \frac{1}{2} m_R^2 \Phi_R^2 + \frac{1}{2} m_I^2 \Phi_I^2 + \ldots,$$  \hfill (45)

where $\Phi \equiv \frac{1}{\sqrt{2}} (\Phi_R + i \Phi_I)$ and

$$V_0 \sim \frac{\mu^2 \beta^2}{12 M_P^2} e^{3-2\gamma} e^x (3 - 2\gamma)^2 (4 - \epsilon) + \frac{M_P^4}{a^4} (R - 1),$$  \hfill (46)

$$\delta_R \sim -\frac{\mu^2 \beta^2}{3 \sqrt{2} M_P^3} e^{(3-2\gamma)e^x (3 - 2\gamma)^2 e(4 - \epsilon) - \frac{3}{\sqrt{2}} g M_P^3 e \left(1 - \frac{2\gamma}{3} - \frac{b}{9g}\right)},$$  \hfill (47)

$$m_R^2 \sim -\frac{\mu^2 \beta^2 \zeta}{2 M_P^2} e^{(3-2\gamma)e^x - 3x (3 - 2\gamma)^2} + \frac{2g^2 M_P^2 \epsilon}{R^3},$$  \hfill (48)

$$m_I^2 \sim -\frac{\mu^2 \beta^2 \zeta}{2 M_P^2} e^{(3-2\gamma)e^x - 3x (3 - 2\gamma)^2},$$  \hfill (49)

$$x \equiv \sqrt{\frac{2}{3} \frac{\varphi}{M_P}},$$  \hfill (50)

$$R \equiv \sqrt{1 + \frac{9}{4} a^4 g^2 \left(1 - \frac{2\gamma}{3} - \frac{b}{9g}\right)^2}. \hfill (51)$$

Here we have assumed $x \gg 1$ and $M_P^2 |\zeta| \gg 1$ and have extracted only the dominant terms. The terms proportional to $\mu$ ($g$) come from $V_F$ ($V_D$), so that we call them as the $F$-($D$-) contributions, respectively.

Let us comment on the role of $\gamma$ in Eq. (36). As can be seen from Eq. (46), the first term spoils the flatness of the scalar potential without $\gamma$. This fact

\footnote{Strictly speaking, due to the kinetic mixing between $C$ and $\Phi$, the scalar $\varphi$ is not canonical. However, we can ignore this by keeping the parameter $\epsilon$ small, $|\epsilon| \ll 1$.}

\footnote{To avoid a tachyonic mass, the $\zeta$ should be negative.}
was already pointed out in Ref. [15]. Hence, we have to assume $\gamma \geq 3/2$ from now on.

For $\gamma \geq 3/2$, the $D$-term contributions are dominant in Eqs. (46) and (47). For the mass term in Eq. (48), however, the dominant term depends on the value of $\zeta$. Hence, we have to assume a relatively large $|\zeta|$ that strongly stabilizes $\chi_R$ at its origin. This requires $|\zeta| \gg \frac{g^2 M_P^2}{\mu^2} e^{-(3-2\gamma)x^2+3x}$. Thus, we obtain $\langle \Phi_R \rangle \sim \langle \Phi_I \rangle \sim 0$ during inflation. In the case of Starobinsky inflation, the $x$ varies between 5.5 and 0.5 according to [31], which implies a huge $|\zeta|$ in general. However, when $\gamma = 3/2 + \delta$ is very close to 3/2, e.g., with $\delta \leq 0.001$, it is enough to require $|\zeta| M_P^2 \gg \frac{g^2 M_P^4}{\mu^2} \times 10^7$.

If $\zeta = 0$, the $D$-term contribution is dominant in Eq. (48). Then, $\chi_R$ tends to be light ($\sim g^2 M_P^2$), and the deviation of $\langle \Phi_R \rangle$ from its origin is typically given by $O(M_P)$. In this case, the isocurvature perturbations should also be included into consideration.

After integrating out $\Phi$, we obtain the effective potential during inflation as

$$V_{eff} = V_D|_{\Phi=0}.$$ (52)

This potential is known to be viable for inflation — see Ref. [16] for details.

One may wonder whether the expansion of the scalar potential of Polonyi superfield up to the second order in Eq. (45) can be trusted. The full scalar potential $V$ given by a sum of Eqs. (37) and (38) is dictated by Eqs. (34) and (35). We checked the stabilization numerically (Figure 1) with the ratio $\beta/M_P = -1$ for the relevant values of $\varphi$.

7Should $\mu^2$ be much smaller than the inflationary scale $\sim V_D$, we could neglect the first term in $V_0$ even when $\gamma < 3/2$. However, the value of $\mu$ comparable with the observed dark energy scale (the cosmological constant) implies the very small SUSY breaking scale in the vacuum, so that we do not consider this situation here.
Figure 1: 3D plot of the potential \( V \) in the \((\varphi, \Phi_R)\) plane with the parameters \( \beta/M_P = -1, \delta = 0.1, \epsilon = -5 \times 10^{-3}, g = 10^{-6}, a = 1, \mu = 10^{-3}, b = -6 \gamma g, \) and \( \zeta = 10^2. \)

4.2 After inflation

In this subsection, we investigate the vacuum structure after inflation in our model. We expand the scalar potential around \( \varphi = 0 \) and \( \Phi = 0, \)

\[
V = V_0 + \delta \varphi \varphi + \delta R \Phi_R + \frac{1}{2} m_\varphi^2 \varphi^2 + m_R^2 \Phi_R \Phi_R + \frac{1}{2} m_{\varphi R}^2 \Phi_R \Phi_{\varphi R} + \frac{1}{2} m_{RR}^2 \Phi_R^2 + \frac{1}{2} m_{II}^2 \Phi_I^2 + \ldots ,
\]

(53)

where we have

- \( V_0 \sim O(\mu^2), \)
- \( \delta \varphi \sim O(\mu^2/M_P), \)
- \( \delta R \sim O(\mu^2/M_P), \)
- \( m_\varphi^2 \sim O(\mu^2/M_P^2), \)
- \( m_R^2 \sim m_I^2 \sim O(\mu^2 \zeta). \)

The explicit values are shown in Appendix C. Also, we neglected the \( D \)-term contributions. This is valid when the SUSY breaking scale \( \mu \) is much larger than the inflation scale \( M_P^2 g^9 \)

\(^8\)Due to \( \zeta \neq 0, \) the value \( \langle \Phi \rangle = (\sqrt{3} - 1)M_P \) of the original Polonyi model does not correspond to a minimum \(^9\)

\(^9\)The \( g \) is set to be \( \sim 10^{-6} \) by the amplitude of the CMB power spectrum. \(^10\)
For large $|\zeta|$, we obtain the following vacuum expectation values:

$$\langle \phi \rangle \sim -\frac{\delta \phi}{m^2_\phi} \sim O(M_P), \quad \langle \Phi_R \rangle \sim O(1/M_P\zeta), \quad \langle \Phi_I \rangle \sim 0.$$  \hfill (59)

Inserting them into Eq. (53) and neglecting the terms suppressed by $\zeta$, we find the minimum of the potential at

$$\langle V \rangle \sim \langle V_F \rangle \sim V_0 - \frac{\delta^2 \phi}{2m^2_\phi} + \frac{\mu^2\epsilon^{3-2\gamma}}{3M^2(4\epsilon - 3)} V(\beta, \gamma, \epsilon).$$  \hfill (60)

Here $V$ is a complicated function of $(\beta, \gamma, \epsilon)$ that can be read off from Eqs. (122) and (124). Like the original Polonyi model, we demand $\langle V_F \rangle = 0$ by tuning the parameter $\beta$.

Note that in the special case $\epsilon = \gamma = 0$, we have

$$V(\beta, 0, 0) = 3 \left( 1 - \frac{3 \beta^2}{M_P^2} \right),$$  \hfill (61)

so that $\beta = M_P/\sqrt{3}$ leads to $\langle V_F \rangle = 0$.

Now $\gamma$ satisfies $\gamma > 3/2$ and, therefore, the value of $\beta$ to obtain $V = 0$ is changed. Unfortunately, we found that there is no real solution $\beta$ of $V = 0$, when $\epsilon = 0$. One can check it by explicitly solving the equation $V(\beta, \gamma, 0) = 0$ that has the following four solutions:

$$\beta/M_P = \pm \sqrt{\frac{\Lambda \pm \sqrt{\Sigma}}{\Gamma}},$$  \hfill (62)

$$\Gamma = 32\gamma^6 - 96\gamma^5 + 32\gamma^4 + 240\gamma^3 - 414\gamma^2 + 162\gamma + 81,$$

$$\Lambda = 48\gamma^3 - 42\gamma^2 + 54,$$

$$\Sigma = 128\gamma^6 - 576\gamma^5 + 1120\gamma^4 - 1296\gamma^3 + 900\gamma^2 - 324\gamma + 81.$$  \hfill (63, 64, 65)

As one can see in Figure 2, there is no real solution of $\beta$ since the argument of the square root is always negative for $\gamma > 3/2$. This is also the case when $\beta$ is extended to a complex parameter, because then the solution (62) is replaced by

$$|\beta|/M_P = \sqrt{\frac{\Lambda \pm \sqrt{\Sigma}}{\Gamma}}.$$  \hfill (66)

The special case with $\gamma = 3/2$ and $\epsilon = \zeta = 0$ yielding a Minkowski vacuum is considered in Appendix D.
that cannot be imaginary.

Figure 2: The red and green (dashed) lines show the “+” and “−” branches, respectively.

However, for a non-vanishing $\epsilon$, the condition $\mathcal{V}(\beta, \gamma, \epsilon) = 0$ can be satisfied. First, let us comment on the allowed regions of $\epsilon$. They are determined by requiring no ghost mode in the system. Since $\epsilon$ induces the kinetic mixing of $\varphi$ and $\Phi$ (see Eq. (34)), we have the following eigenvalues:

$$\lambda_{\pm} = 1 + 4\zeta|\Phi|^2 \pm \sqrt{16\zeta^2|\Phi|^4 + \frac{\epsilon^2}{3}e^{2x}}$$

in the diagonal basis of the kinetic terms. Therefore, to avoid a ghost mode during inflation, the $\epsilon$ must satisfy

$$|\epsilon| < \sqrt{3}e^{-x},$$

where we have neglected the terms including $\Phi$, because they are suppressed by $\zeta$. At the point $x = 5.5$, we have a constraint $|\epsilon| < \sqrt{3}e^{-x} \sim 7 \times 10^{-3}$.

For example, given $\delta = 10^{-5}$ and $\epsilon = -5.0 \times 10^{-3}$ we obtain the two solutions numerically:

$$(i) \quad \beta \sim -76M_P,$$

$$(ii) \quad \beta \sim -266M_P.$$

$^{11}$In the case of $\epsilon > 0$, there is a singularity in the scalar potential.
In both cases, the masses of $\varphi, \Phi_R$ and $\Phi_I$ are estimated as

\begin{align}
(i) & \quad m_{\varphi}^2 \sim 6.0 \times 10^{-2} \mu^2/M_P^2, \quad m_R^2 \sim m_I^2 \sim 2.1 \mu^2 |\zeta|, \\
(ii) & \quad m_{\varphi}^2 \sim 3.9 \mu^2/M_P^2, \quad m_R^2 \sim m_I^2 \sim 148 \mu^2 |\zeta|,
\end{align}

while they are all positive.

To summarize this subsection, due to the existence of $\epsilon \neq 0$, we can find solutions of $\beta$ to get $V = 0$.\footnote{We neglected the $\zeta$-suppressed terms in our analysis. When taking into account their contribution, our solutions for $\beta$ are going to be corrected by the $O(1/\zeta)$ terms.} Then, the $D$-term contributions to the vacuum energy, neglected in Eqs. (53) and (60), become important. The vacuum energy produced by $V_D$ is given by

\begin{equation}
\langle V \rangle \sim \langle V_D \rangle \\
\sim \frac{M_P^4}{a^4} \left( \sqrt{1 + \frac{9}{4} a^4 g^2} \left( 1 - \langle e^{-x} \rangle - \frac{2}{3} \gamma - \frac{b}{9g} \right)^2 - 1 \right),
\end{equation}

where $\langle x \rangle$ can be read off from Eq. (59). To obtain a tiny positive cosmological constant, we have to tune the value of $b$. For example, by setting $b = 9g \left( 1 - \langle e^{-x} \rangle - \frac{2}{3} \gamma \right) + \tilde{\epsilon}$, we find

\begin{equation}
\langle V \rangle \sim \frac{M_P^4}{72} \tilde{\epsilon}^2.
\end{equation}

To this end, we comment on the SUSY breaking scale in the vacuum. The vacuum expectation values of $F\Phi$ and $D$ are given by

\begin{align}
|\langle F\Phi \rangle| & \sim \mu e^{2\langle J \rangle/M_P^2} \frac{1}{\langle \Delta \rangle} \left( 1 - 2\epsilon \frac{\beta}{\langle M_P e^x + Q \rangle} \right), \\
|\langle D \rangle| & \sim \frac{M_P^2}{6} \tilde{\epsilon}.
\end{align}

As is clear from these equations, the F- and D- type breaking scales are decoupled, while the $D$-type scale can be arbitrarily small. In the examples (i) and (ii) above, the $|\langle F\Phi \rangle|$ is given by

\begin{align}
(i) & \quad |\langle F\Phi \rangle| \sim 1.8 \times 10^2 \mu, \\
(ii) & \quad |\langle F\Phi \rangle| \sim 9.8 \times 10^{-3} \mu.
\end{align}
5 Superspace actions and super-Higgs effect

It is instructive to reformulate our model in the curved superspace of the old-minimal supergravity, because it helps in construction of more general supergravity models, as well as for verifying our calculations based on the superconformal tensor calculus.

5.1 Superspace Lagrangian

Let us consider the following superspace Lagrangian by using the standard notation and conventions of Ref. [33],

\[ \mathcal{L} = \int d^2\Theta^2 \mathcal{E} \left[ \frac{3}{8}(\bar{D}^2 - 8\mathcal{R})e^{-\frac{1}{2}j} \left( 1 + \frac{4gI}{9} \frac{W^2\bar{W}^2}{(D\bar{W})^3} \right) + W(S_i) + L_{DBI} \right] + \text{h.c.} \]  

(79)

Here \( J \) and \( I \) are arbitrary real functions of chiral superfields \( S_i, \bar{S}_i \) and a real massive superfield \( V \), whose gauge coupling is \( g \); \( W(S_i) \) is a holomorphic superpotential; \( W \equiv W^aW_a \) and \( DW \equiv D^aW_a \) where \( W_a \) is the chiral superfield strength of \( V \). The \( L_{DBI} \) is the DBI contribution,

\[ L_{DBI} = \frac{1}{4} f W^2 - \frac{\alpha}{16}(\bar{D}^2 - 8\mathcal{R}) \frac{|f + \bar{f}|^2 W^2 \bar{W}^2}{1 + 4\alpha(\omega + \bar{\omega}) + \sqrt{1 + 8\alpha(\omega + \bar{\omega}) + \frac{\alpha^2}{16}(\omega - \bar{\omega})^2}} \]  

(80)

where \( f = f(S_i) \) is a holomorphic gauge kinetic function, \( \alpha \) is the DBI parameter extended to a function of the chiral (matter) superfields, and \( \omega \) is

\[ \omega = \frac{1}{16}(f + \bar{f})(\bar{D}^2 - 8\mathcal{R})W^2. \]  

(81)

The Lagrangian (79) is written after the superconformal gauge fixing in the so-called Jordan frame. After eliminating the auxiliary fields and the Weyl rescaling to the Einstein frame, the bosonic component Lagrangian reads

\[ e^{-1}\mathcal{L} = \frac{1}{2} R + \sqrt{1 + 8\alpha e^{-\frac{1}{4}J} \frac{Z^2}{f_R} \left( \frac{gI}{6Z} - \frac{1}{4} \right) f_R F^2 + \frac{i}{4} f_I F \bar{F} - \frac{1}{2} J_{CC} \partial^a C \partial_a C - 2J_{ij} \partial^a S^i \partial_a \bar{S}^j - \frac{1}{2} g^2 J_{CC} B^a B_a - J_{Gi} \partial^a C \partial_a S^i - J_{G\bar{i}} \partial^a C \partial_a \bar{S}^i + iB^a (J_{Gi} \partial_a S^i - J_{G\bar{i}} \partial_a \bar{S}^i) - V_D - V_F + \mathcal{O}(F^4) \]  

(82)
with the F- and D-type scalar potentials

\[ V_F = e^{2\sigma} \left[ 2\Delta^{ij} \left( \chi_i W + \frac{1}{2} W_i \right) \left( \bar{\chi}_j \bar{W} + \frac{1}{2} \bar{W}_j \right) - \left( 3 - 2 \frac{J^2}{J_{CC}} \right) |W|^2 \right], \quad (83) \]

\[ V_D = \frac{e^{4\sigma}}{8\alpha} \left[ \sqrt{1 + 8\alpha e^{-\frac{4\sigma}{3}Z^2/f_R}} - 1 \right] = \frac{Z^2}{2f_R} + \mathcal{O}(\alpha), \quad (84) \]

where we have introduced the notation

\[ Z \equiv -g \left( J_C + \frac{T}{6} \right), \quad f_R \equiv \text{Re} f, \quad f_I \equiv \text{Im} f, \]

\[ \Delta^{ij} \equiv \left( J_{Ci} J_{Cj} \right)^{-1}, \quad \chi_i \equiv J_i - \frac{J_{Ci} J_{Cj}}{J_{CC}}, \quad \bar{\chi}_i \equiv \bar{J}_i - \frac{J_{C\bar{i}} J_{C\bar{j}}}{J_{CC}}. \quad (85) \]

The Lagrangian [79] is more general than the one of Sec. 2 because it includes the gauge kinetic function \( f(S_i) \). See Appendix A for more.

### 5.2 The gauge-invariant reformulation

The Lagrangian [79] can be rewritten to the manifestly \( U(1) \) gauge-invariant and manifestly supersymmetric form. Let us begin with another action

\[ \mathcal{L} = \int d^2\Theta \mathcal{E} \left[ \frac{3}{8} (\mathcal{D}^2 - 8\mathcal{R}) e^{-\frac{1}{4}K + \Gamma} \left( 1 + \frac{4g I}{9} \frac{W^2 \bar{W}^2}{(D W)^3} \right) + W(S_i) + L_{DBI} \right] + \text{h.c.} \quad (86) \]

where we have introduced the Kähler potential \( K = K(T, \bar{T}, S_i, \bar{S}_i) \) depending upon the Higgs/Stückelberg superfield \( T \) and the neutral chiral superfields \( S_i \). The \( \Gamma = \Gamma(T, \bar{T}, V) \) is the counterterm for the gauge transformations of \( T \) (so that a sum \( K + \Gamma \) is gauge invariant), while the vector gauge superfield \( V \) is massless. Here the \( \mathcal{I} \) is an arbitrary gauge-invariant real function \( \mathcal{I}(T, \bar{T}, S_i, \bar{S}_i, V) \). The \( L_{DBI} \) is unchanged.

The linearly realized \( U(1) \) gauge symmetry can be made explicit by requiring that the function \( K + \Gamma \) only depends upon the product \( Te^{2gV} \bar{T} \) that is invariant with respect to the gauge transformations

\[ V \to V - \frac{1}{2}(\Lambda + \bar{\Lambda}) , \quad T \to Te^{g\Lambda} , \quad (87) \]
where $\Lambda$ is an arbitrary chiral superfield. However, for our purposes, it is more convenient to let $T$ transform under the $U(1)$ as follows:

$$V \to V - \frac{1}{2}(\Lambda + \bar{\Lambda})\; , \; T \to T - g\Lambda\; , \tag{88}$$

so that the gauge-invariant combination is given by $T + \bar{T} - 2gV$.

It is straightforward to calculate the bosonic terms of the Lagrangian (86). We find

$$e^{-1}\mathcal{L} = \frac{1}{2}R + \sqrt{1 + 8\alpha e^{-\frac{2}{3}K}\bar{Z}^2 \left( \frac{g\bar{T}}{6\bar{Z}} - \frac{1}{4} \right) f_R F^2 + \frac{i}{4} f_I F \tilde{F} - K_{ij} \partial^a S^j \partial_a S^j - K_{T\bar{T}} D^a T \bar{D}^a \bar{T} - K_{t\bar{t}} \partial^a S^i \bar{D}^a T - K_{T\bar{t}} D^a T \partial_a \bar{S}^j - V_D - V_F + \mathcal{O}(F^4) } \; , \tag{89}$$

where the $F$- and $D$-type scalar potentials are

$$V_F = e^K \left[ K^{IJ} (K_I W + W_I) (K_J \bar{W} + \bar{W}_J) - 3|W|^2 \right] \; , \tag{90}$$

$$V_D = \frac{e^{\frac{2}{3}K}}{8\alpha} \left[ \sqrt{1 + 8\alpha e^{-\frac{2}{3}K}\bar{Z}^2 \left( \frac{g\bar{T}}{6\bar{Z}} - \frac{1}{4} \right) f_R F^2 + \frac{i}{4} f_I F \tilde{F} - K_{ij} \partial^a S^j \partial_a S^j - K_{T\bar{T}} D^a T \bar{D}^a \bar{T} - K_{t\bar{t}} \partial^a S^i \bar{D}^a T - K_{T\bar{t}} D^a T \partial_a \bar{S}^j - V_D - V_F + \mathcal{O}(F^4) } \right] = \frac{\bar{Z}^2}{2f_R} + \mathcal{O}(\alpha) \; . \tag{91}$$

Here $D_a$ is the $U(1)$ covariant derivative

$$D_a \equiv \partial_a - gB_a X^T \; , \tag{92}$$

where $X^T$ is the Killing vector. The indices $I, J$ in Eq. (90) denote the scalar field $T$ and $S$, together, i.e. $I = \{i, T\}$. Since we do not consider the $U(1)$ gauge symmetry as the $R$-symmetry, the superpotential $W$ should be independent of $T$, i.e. $W_T = 0$. The $\bar{Z}$ stands for

$$\bar{Z} \equiv -g \left( \mathcal{D} + \frac{T}{6} \right) \; , \tag{93}$$

where $\mathcal{D}$ is the Killing potential, which can be expressed as

$$\mathcal{D} = iK_T X^T = -iK_{\bar{T}} X^{\bar{T}} \; . \tag{94}$$

A correspondence between the gauge-invariant Lagrangian (86) and its "massive" counterpart (79) can be established as follows. The "massive"

\footnote{We use the same letters to denote the chiral superfields and their leading components.}

\footnote{The Killing vector of the linear gauge transformation (87) is $X^T = -iT$, whereas for the gauge transformation (88) we get $X^T = i$.}
description corresponds to the unitary gauge, \( T = 0 \), in the convention (88), of the Lagrangian (86). In this gauge, we choose \( \Lambda = T/g \) in Eq. (88), which leads to
\[
V \rightarrow \tilde{V} = V - \frac{1}{2g} (T + \bar{T}) \quad \text{and} \quad T \to 0.
\]

On the other hand, in the gauge-invariant formulation, we can use the Wess-Zumino gauge where the auxiliary components of \( V \) are gauged away. In the leading order, this leads to
\[
T + \bar{T} - 2gC \to \tilde{T} + \bar{\tilde{T}},
\]
where \( \tilde{T} = T - gC \) is the Stückelberg scalar that absorbs the auxiliary real scalar \( C \). Since the left-hand-side of Eq. (95) in the leading order reads \( T + \bar{T} - 2gC \to -2g\tilde{C} \), comparing it with Eq. (96) yields the relation
\[
2g\tilde{C} = -\tilde{T} - \bar{\tilde{T}}.
\]

Thus we can easily switch between the Lagrangian (79) in terms of \( J(C, \ldots) \) and the Lagrangian (86) in terms of \( K(T, \bar{T}, \ldots) \).

As an example, let us consider the Starobinsky case with \( J \) in Eq. (79) given by
\[
2J = -3 \log(-C) - 3C.
\]
This can be translated into the gauge-invariant notation by sending \( C \to 2gC \) in Eq. (98) and using Eq. (97). The resulting theory is now described by the Lagrangian (86) with
\[
K = -3 \log(T + \tilde{T}) + 3(T + \tilde{T}),
\]
whose first term corresponds to the no-scale supergravity.

6 Conclusion

Our purpose and motivation are to unify (i) Starobinsky inflation, (ii) high-scale spontaneous SUSY breaking after inflation, and (iii) a positive cosmological constant, in supergravity. To achieve the goal (i), we employ the minimal number of the physical degrees of freedom contained in the single massive vector multiplet: the inflaton, the goldstino, and the massive vector field. The possible observational signatures of the massive vector field may

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15The gauge couplings \( g \) in the two formulations are related to each other by the factor of two, i.e. one should take \( g \to g/2 \) in Eq. (86), or \( g \to 2g \) in Eq. (79).
arise as non-gaussianities of the CMB spectrum [34]. To achieve the goal (ii), needed for reheating and viable particle phenomenology, we employ the minimal hidden sector described by the single chiral (Polonyi) multiplet. The goal (iii) is needed to describe the dark energy (the accelerating Universe) via a de Sitter vacuum.

Despite the considerable progress, mentioned in our references and in the main text, this line of research faced several problems, such as (a) instability of the inflationary trajectory due to the mixing of the inflaton and Polonyi scalar, (b) the limited use of the standard FI term in the supergravity-based cosmology, and (c) the problem of decoupling the F-type and D-type contributions to the scalar potential, needed for the hierarchy between the observed cosmological constant and a much higher SUSY breaking scale.

Overcoming these problems is apparently possible only with additional theoretical resources provided by the DBI structure and the new (generalized) FI terms that do not require the R-symmetry gauging and, hence, allow for more general couplings.

We demonstrated in this paper that the goals (i), (ii) and (iii) can be simultaneously achieved, though in the rather complicated way, via the extension of the original Polonyi model and the tuning of the parameters.

Our model after stabilization of the Polonyi scalar by construction leads to the same predictions for the CMB observables as the Starobinsky model due to the choice (36) of the $J$-function and the chosen value of the coupling constant $g$ in Subsect. 4.1. As regards the tiny value of the cosmological constant $\Lambda^{1/4} \approx 0.0024$ eV, it arises entirely due to the D-term by demanding (fine-tuning) the vanishing F-type term, i.e. the vanishing difference between the two relatively large contributions with the opposite signs in Eq. (60).

As regards the so-called Polonyi problem (i.e. possible overproduction of Polonyi particles in the early Universe after inflation), it may be avoided in the supergravity model under consideration [38].

We used the more general Polonyi couplings demanded by consistency of our approach. It is still desirable to fix interactions of the hidden (Polonyi) sector by some organizing (symmetry) principles beyond local supersymmetry. It could be S-duality or extra (hidden and non-linearly realized) supersymmetry of the relevant actions in the context of partial SUSY breaking, e.g., by exploiting possible connections to $D3$-branes and their anti-branes in string theory [35], though this is beyond the scope of this investigation.
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A Extended DBI sector

Let us study some matter coupled generalizations of the DBI sector, $S_{DBI}$, in the superconformal setting. These extensions can be found, e.g., in Refs. [36, 37], where the constraint related to the partial breaking of $N = 2$ supersymmetry was used. As a specific extension of the DBI sector, let us take

$$S_{DBI} = -\frac{1}{2} [f W^2]_F + \left[ \frac{\alpha \Psi W^2 \bar{W}^2}{|S_0|^4 (1 - 2\alpha \mathcal{A} + \sqrt{1 - 4\alpha \mathcal{A} + 4\alpha^2 \mathcal{B}^2})} \right] D,$$

where

$$\mathcal{A} = \bar{\Sigma} \left( \frac{W^2}{|S_0|^4} \Omega \right) + \text{h.c.}, \quad \mathcal{B} = \bar{\Sigma} \left( \frac{W^2}{|S_0|^4} \Omega \right) - \text{h.c.},$$

and $f$ is a holomorphic gauge kinetic function of $S^i$. In general, $\alpha$ (real), $\Psi$ (real) and $\Omega$ (complex) depend on $V, S^i$ and $\bar{S}^i$. All of them are the weightless multiplets. The bosonic expansion of Eq. (100) reads

$$L_{DBI} = -\frac{1}{4} \left( f_R - \frac{\Psi \Omega_R}{|\Omega|^2} \right) FF + \frac{i}{4} \left( f_I - \frac{\Psi \Omega_I}{|\omega|^2} \right) F\tilde{F} + \frac{1}{2} \left( f_R - \frac{\Psi \Omega_R}{|\Omega|^2} \right) D^2 + \frac{|S_0|^4 \Psi}{8\alpha |\Omega|^2} \left[ 1 - \left\{ 1 + \frac{4\alpha}{|S_0|^4} \left( -2\Omega_R D^2 + \Omega_R FF - i\Omega_I F\tilde{F} \right) \right. \right.$$

$$\left. \left. + \frac{4\alpha^2}{|S_0|^8} \left( -2i\Omega_I D^2 + i\Omega_I FF - \Omega_R F\tilde{F} \right)^2 \right\} \right]^{1/2},$$

(102)
where we have used the same letters to denote the first components of \( f, \alpha, \Psi \) and \( \Omega \) as for the corresponding multiplets or superfields. In addition, the “\( R \)” and “\( I \)” subscripts denote the real and imaginary parts.

First, we demand the above Lagrangian to have the DBI form (for \( F_{ab} \)) when \( D = 0 \), by using the identity
\[
-\det (\eta_{ab} + aF_{ab}) = 1 + \frac{a^2}{2} FF + \frac{a^4}{16} (F \tilde{F})^2.
\] (103)

This gives rise to the following consequences:

1. The first term in Eq. (102) should vanish, because Maxwell term comes from the expansion of the DBI term in the weak field limit. On the other hand, the \( F \tilde{F} \) does not come from the DBI expansion, so that the second term in Eq. (102) should not vanish.

2. Inside the square root (focusing on the third and fourth lines of Eq. (102)), we need \( FF \), but not \( (FF)^2 \). In the same way, we need \( (F \tilde{F})^2 \), but not \( F \tilde{F} \). Otherwise, the expansion formula (103) can never be applied.

After taking the both consequences into account, we obtain
\[
\Omega_R = \frac{\Psi}{f_R} \quad \text{and} \quad \Omega_I = 0.
\] (104)

For example, the case
\[
\Psi = f_R^2
\] (105)
leads to
\[
\mathcal{L}_{DBI} = \frac{i}{4} f_I F \tilde{F} + \frac{|S_0|^4}{8\alpha} \left[ 1 - \sqrt{-\det \left( g_{\mu\nu} + \frac{8\alpha f_R}{|S_0|^4} F_{\mu\nu} \right)} \right].
\] (106)

Hence, when \( D \neq 0 \), as in our case, the DBI structure becomes generically more complicated.
B Derivation of $S_{V,M}$

The bosonic expansion of Eq. (12) reads

$$L_{V,M} = -{1 \over 3} |S_0|^2 \mathcal{H} R(\omega) + 2 \mathcal{H} (|F_0|^2 - |D_a S_0|^2) + |S_0|^2 \mathcal{H}_C D_a S_0 + 2 |S_0|^2 \mathcal{H}_{ij} \left( F^i \bar{F}^j - \partial_a S^i \partial^a \bar{S}^j \right)$$

$$+ \left\{ -\mathcal{H}_C N S_0 \bar{F}_0 + \mathcal{H}_C i (B_a + iD_a C) S_0 D^a \bar{S}_0 - |S_0|^2 \mathcal{H}_C i (F^i \bar{N} + i\partial_a S^i (B^a - iD^a C)) \right.$$ 

$$+ 2 S_0 \mathcal{H}_i \bar{F}_0 F^i - 2 S_0 \mathcal{H}_i \bar{S}_0 \partial^a S^i + 3 S_0^2 F_0 W + S_0^3 F^i W_i + h.c. \right\} ,$$

$$\text{(107)}$$

where $N \equiv H + iK$. The subscripts on $\mathcal{H}$ and $W$ denote the derivatives with respect to $C, S^i$ and $\bar{S}^i$, e.g., $\mathcal{H}_C = \partial \mathcal{H}/\partial C$ and $\mathcal{H}_i = \partial \mathcal{H}/\partial S^i$. The $D_a$ is the superconformal covariant derivative $[30]$, whose bosonic part is given by

$$D_a S_0 = \partial_a S_0 - i A_a S_0 - b_a S_0, \quad D_a C = \partial_a C - 2 b_a C. \quad (108)$$

We take the following superconformal gauge fixing conditions:

$$\text{D – gauge : } -{1 \over 3} |S_0|^2 \mathcal{H} = {1 \over 2} ,$$

$$\text{A – gauge : } S_0 = \bar{S}_0 , \quad (110)$$

$$\text{K – gauge : } b_\mu = 0 , \quad (111)$$

which ensure that the Ricci scalar is canonically normalized in the action. $^{16}$ Then the $R(\omega)$ becomes the usual Ricci scalar that we denote by $R$. Under these gauge conditions, we can rewrite Eq. (107) as

$$L_{V,M} = {1 \over 2} R + 2 \mathcal{H} |F_0|^2 - 3 \mathcal{H} |F_0|^2 - \mathcal{H} A_a J^a + 3 A_a^2$$

$$- {3 \over 2 \mathcal{H}} \left\{ \mathcal{H}_C D_a S_0 + {1 \over 2} \mathcal{H}_{ij} \left( |N|^2 - B_a^2 - (\partial_a C)^2 \right) + 2 \mathcal{H}_{ij} \left( F^i \bar{F}^j - \partial_a S^i \partial^a \bar{S}^j \right) \right\}$$

$$+ \left\{ {3 \mathcal{H}_C i \over 2 \mathcal{H}} (F^i \bar{N} + i\partial_a S^i (B^a - iD^a C)) + \sqrt{-3 \over 2 \mathcal{H}} (-\mathcal{H}_C N \bar{F}_0 + 2 \mathcal{H}_i \bar{F}_0 F^i) \right.$$ 

$$- {9 \over 2 \mathcal{H}} F_0 W + \left( -{3 \over 2 \mathcal{H}} \right)^{3/2} F^i W_i + h.c. \right\} ,$$

$$\text{(112)}$$

$^{16}$The gauge fixing condition for special SUSY is irrelevant for the bosonic part.
where we have used the notation
\[ I_a = \mathcal{H}_C \partial_a C + \mathcal{H}_i \partial_a S^i + \mathcal{H}_j \partial_a \bar{S}^j, \quad (113) \]
\[ J_a = \mathcal{H}_C B_a + i\mathcal{H}_i \partial_a S^i - i\mathcal{H}_j \partial_a \bar{S}^j. \quad (114) \]

Next, the auxiliary fields \( A_a, F^i, N \) and \( D \) should be eliminated. As was already mentioned in the main text, only \( D \) non-trivially appears in \( \mathcal{L}_{DBI} \) and \( \mathcal{L}_{FI} \), and therefore, we can easily integrate over the other auxiliary fields. Their algebraic solutions are
\[ A_a = -\frac{1}{2\mathcal{H}} J_a, \quad (115) \]
\[ F^i = -e^\mathcal{J} \Delta^{ij} \left\{ \frac{1}{2} \bar{W}_j + \left( J_j - \frac{\mathcal{J}_C \mathcal{J}_C}{\mathcal{J}_{CC}} \right) \bar{W} \right\}, \quad (116) \]
\[ N = \frac{2}{\mathcal{J}_{CC}} \left( \mathcal{J}_C F^i + e^{\mathcal{J}} \mathcal{J}_C \bar{W} \right), \quad (117) \]
\[ F_0 = \frac{2}{3} e^{\mathcal{J}/3} \left( \mathcal{J}_i - \frac{\mathcal{J}_C \mathcal{J}_C}{\mathcal{J}_{CC}} \right) F^i - e^{4\mathcal{J}/3} \left( \frac{2\mathcal{J}_C^2}{3\mathcal{J}_{CC}} - 1 \right) \bar{W}, \quad (118) \]
where we have defined
\[ \mathcal{J} = -\frac{3}{2} \log \left( -\frac{2}{3} \mathcal{H} \right), \quad (119) \]
\[ \Delta^{ij} = (\Delta_{ij})^{-1}, \quad \Delta_{ij} = \mathcal{J}_{ij} - \frac{\mathcal{J}_C \mathcal{J}_C}{\mathcal{J}_{CC}}. \quad (120) \]

After substituting these solutions into Eq. (112), we finally obtain
\[ \mathcal{L}_{V,M} = \frac{1}{2} R - \frac{1}{2} \mathcal{J}_{CC} (\partial_a C)^2 - \frac{1}{2} \mathcal{J}_{CC} B_a^2 + \mathcal{J}_C D \]
\[ - 2\mathcal{J}_{ij} \partial_a S^i \partial^a \bar{S}^j + (-\mathcal{J}_C \partial_a C \partial^a S^i - i\mathcal{J}_C B_a \partial^a S^i + \text{h.c.}) \]
\[ - e^{2\mathcal{J}} \left\{ 2\Delta^{ij} \left\{ \frac{1}{2} \bar{W}_i + \left( \mathcal{J}_i - \frac{\mathcal{J}_C \mathcal{J}_C}{\mathcal{J}_{CC}} \right) \bar{W} \right\} \right\} \left\{ \frac{1}{2} \bar{W}_j + \left( \mathcal{J}_j - \frac{\mathcal{J}_C \mathcal{J}_C}{\mathcal{J}_{CC}} \right) \bar{W} \right\} \]
\[ + 2\frac{\mathcal{J}_C^2}{\mathcal{J}_{CC}} |\bar{W}|^2 - 3|\bar{W}|^2 \right\}. \quad (121) \]

C Derivation of the coefficients in Eq. (53)

Here we compute the explicit coefficients of Eq. (53). We skip writing down \( \delta_R \) and \( m^2_{Rc} \), whose explicit expressions are not used in the main text.
First, the $V_0$ is given by
\[
V_0 = \frac{e^{3-2\gamma}\mu^2}{3M_P^2(3-4\epsilon)} \left[ 9M_P^2 - 12M_P\beta(3+2\gamma)\epsilon + \beta^2(3+2\gamma)(-9+12\epsilon+12\epsilon^2) + 6\gamma + 8\gamma\epsilon(\epsilon-1) \right].
\] (122)

The linear term $\delta \varphi$ reads
\[
\delta \varphi = \frac{2\mu^2}{3M_P^2(3-4\epsilon)^2} \sqrt{\frac{2}{3}} e^{3-2\gamma} \left[ 9M_P^2(4\epsilon\epsilon + 4\epsilon - 3\gamma) + \beta^2 \left\{ 4\gamma^3(4\epsilon - 3)(4\epsilon - 4\epsilon^2 + 3) 
\right.
\right.
\]
\[
+ 4\gamma^2 \left( 4\epsilon \left( 8\epsilon^2 + \epsilon - 6 \right) + 9 \right) + 3\gamma(4\epsilon(12\epsilon^2 + 7\epsilon - 6) + 9) + 144\epsilon^3 \bigg] 
\]
\[
- 12\beta M_P \epsilon \left( 2\gamma^2(-3 + 4\epsilon) + 3\gamma(-1 + 4\epsilon) + 12\epsilon \right) \right].
\] (123)

The $m^2_\varphi$ is given by the following complicated expression:
\[
m^2_\varphi = -\frac{2\mu^2 e^{3-2\gamma}}{9M_P^2(4\epsilon-3)^3} \left[ 12\beta M_P \epsilon \left\{ 8\gamma^3(-16\epsilon^2 + 24\epsilon - 9) + 8\gamma^2(-16\epsilon^2 + 9) 
\right.
\right.
\]
\[
+ 24\gamma(-16\epsilon^2 + 8\epsilon - 3) - 3(80\epsilon^2 + 24\epsilon + 9) \bigg] 
\]
\[
+ 9M_P^2 \left( 4\gamma^2(3-4\epsilon)^2 + 6\gamma(4\epsilon - 3)(4\epsilon + 1) + 8\epsilon(14\epsilon - 3) + 27 \right) 
\]
\[
+ \beta^2 \left\{ 16\gamma^4(3+4\epsilon - 4\epsilon^2)(3-4\epsilon)^2 + 8\gamma^3(4\epsilon - 3)(4\epsilon^2(29 + 4\epsilon) - 120\epsilon + 45) 
\right.
\]
\[
+ 8\gamma^2(4\epsilon^2(171 - 80\epsilon + 80\epsilon^2) - 540\epsilon + 135) 
\]
\[
+ 6\gamma(4\epsilon^2(-45 + 24\epsilon + 176\epsilon^2) + 324\epsilon - 81) 
\]
\[
+ 9(4\epsilon^2(9 + 4\epsilon)(-5 + 12\epsilon) + 108\epsilon - 27) \bigg] \right].
\] (124)

Finally, as regards the $m^2_R$ and the $m^2_I$, we only show their dominant terms with respect to $\zeta$ as follows:
\[
m^2_R = m^2_I = -\frac{8\zeta\mu^2}{M_P^2(3-4\epsilon)^2} e^{3-2\gamma}(3M_P^2 - 2\beta(2\gamma + 3)\epsilon)^2.
\] (125)

**D Minkowski vacuum after inflation**

Let us discuss in more detail the possibility of Minkowski vacuum for the model described in Sec. 4 with $\epsilon = \zeta = 0$. Consider the superspace action
with \( f = 1 \) and a constant FI parameter \( I \) (we take \( M_P = 1 \) below):

\[
2\mathcal{J} = A \log(-C) - BC + \Phi \bar{\Phi} + d, \tag{126}
\]

\[
W = \mu(\Phi + \beta), \tag{127}
\]

where \( C \) is the real scalar component \( V, \Phi \) is the Polonyi complex scalar, and \( A, B, d, \mu, \beta \) are real constant parameters. This choice is related to the \( \mathcal{J} \)-function of Sec. 4 by setting \( A = -3 \) and \( B = 3 - 2\gamma \).

The choice (126) and (127) leads to the scalar potentials

\[
V_F = \mu^2(-C)^A e^{\Phi \bar{\Phi} - BC + d} \left(|\Phi \bar{\Phi} + \beta \Phi + 1|^2 - X|\Phi + \beta|^2\right), \tag{128}
\]

\[
V_D = \frac{g^2}{8} \left(B - \frac{A}{C} - \frac{I}{3}\right)^2 + O(\alpha), \tag{129}
\]

\[
X \equiv 3 + A - 2BC + \frac{B^2C^2}{A}. \tag{130}
\]

Requiring the canonical kinetic term for the inflaton yields \( C = \lambda e^{\sigma \varphi} \), with \( \sigma = \sqrt{-2/A}. \) Using the notation of Ref. [16], and can set \( \lambda = -1 \) without loss of generality because \( \lambda \) can be absorbed by \( B \) and \( d \) in Eq. (126). During inflation, unless \( B \leq 0 \), the double exponential factor creates inflationary instability [15]. In addition, the absence of ghosts requires \( A < 0 \). Taking all this into account, it follows from Eq. (84) that, in order to obtain the Starobinsky inflationary potential, we must require \( B - I/3 > 0 \), or \( I < -3|B| \).

The Minkowski vacuum conditions for the potential of Eqs. (128) and (129) read

\[
\mathcal{V} = \mu^2(-C_0)^A e^{\Phi_0^2 - BC_0 + d} \left[(\Phi_0^2 + \beta \Phi_0 + 1)^2 - X_0(\Phi_0 + \beta)^2\right] + \frac{g^2}{8} \left(B - \frac{A}{C_0} - \frac{I_0}{3}\right)^2 = 0, \tag{131}
\]

\[
\mathcal{V}_\Phi = \Phi_0 \mathcal{V}_F + \mu^2(-C_0)^A e^{\Phi_0^2 - BC_0 + d} \times \left[(2\Phi_0 + \beta)(\Phi_0^2 + \beta \Phi_0 + 1) - X_0(\Phi_0 + \beta)\right] = 0, \tag{132}
\]

\[
\mathcal{V}_\varphi = \frac{\sigma A}{C_0} \left(B - \frac{A}{C_0} - \frac{I_0}{3}\right) + \sigma(A - BC_0) \mathcal{V}_F + 2\sigma \mu^2(-C_0)^A e^{\Phi_0^2 - BC_0 + d}(\Phi_0 + \beta)^2 BC_0 \left(1 - \frac{BC_0}{A}\right) = 0, \tag{133}
\]

where the subscript 0 denotes the vacuum expectation values.
Let us consider the case when $V_D$ and $V_F$ separately vanish at the minimum. Then, as regards $C_0$, we have

$$C_0 = -e^{\sigma \phi_0} = \frac{A}{B - \frac{2}{3}},$$

(134)

while Eqs. (131)-(133) are reduced to

$$(\Phi_0^2 + \beta \Phi_0 + 1)^2 = X_0 (\Phi_0 + \beta)^2,$$

(135)

$$(2\Phi_0 + \beta)(\Phi_0^2 + \beta \Phi_0 + 1) = X_0 (\Phi_0 + \beta),$$

(136)

$$(\Phi_0 + \beta)^2 BC_0 \left(1 - \frac{BC_0}{A}\right) = 0,$$

(137)

respectively.

As is clear from Eq. (137), if we assume a non-vanishing $B$, then we get $B, C_0, A < 0$, and the only way to satisfy Eq. (137) is to require $\Phi_0 + \beta = 0$ that is inconsistent with Eq. (135). Hence, there is no Minkowski vacuum (at least for the case $V_D = V_F = 0$) unless $B = 0$ (i.e. $\gamma = 3/2$).

Given $B = 0$, we have

$$X_0 = X = 3 + A \quad \text{and} \quad C_0 = -\frac{3A}{I},$$

(138)

and the solutions to Eqs. (135)(136) are given by

(a) $\Phi_0 = \sqrt{3 + A} \pm 1, \quad \beta = -\sqrt{3 + A} \mp 2,$

(139)

(b) $\Phi_0 = -\sqrt{3 + A} \pm 1, \quad \beta = \sqrt{3 + A} \mp 2$.

(140)

Since $A$ must be negative, there is the additional restriction $-3 \leq A < 0$. Our choice of Sec. 4, $A = -3$, is the lowest allowed value in this case.

To summarize, we began with 7 parameters $A, B, d, g, \mu, \beta, I$. Requiring the Starobinsky inflationary potential and Minkowski vacuum with spontaneously broken SUSY after inflation leaves only two parameters $g$ and $\mu$. The $B$ should vanish, and the $d$ can also vanish because it does not have a significant impact. The orders of magnitude of $A$ and $\beta$ are also fixed, whereas the $I$ can be absorbed into a redefinition of $\varphi$ and $g$. The gauge coupling $g$ controls the mass of the vector multiplet, including the mass of the inflaton $\varphi$, while the parameter $\mu$ is proportional to the gravitino mass and remains arbitrary in our models.

\footnote{In this case, the DBI corrections can be dropped because they are all proportional to powers of $(B - \frac{A}{2} - \frac{2}{3})$.}
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