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Starobinsky-like inflation and soft-SUSY breaking

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ABSTRACT: We study a version of Starobinsky-like inflation in no-scale supergravity (SUGRA) where a Polonyi term in the hidden sector breaks supersymmetry (SUSY) after inflation, providing a link between the gravitino mass and inflation. We extend the theory to the visible sector and calculate the soft-SUSY breaking parameters depending on the modular weights in the superpotential and choice of Kähler potential. We are led to either no-scale SUGRA or pure gravity mediated SUSY breaking patterns, but with inflationary constraints on the Polonyi term setting a strict upper bound on the gravitino mass $m_{3/2} < 10^3$ TeV. Since gaugino masses are significantly lighter than $m_{3/2}$, this suggests that SUSY may be discovered at the LHC or FCC.

KEYWORDS: Beyond Standard Model, Cosmology of Theories beyond the SM, Supersymmetric Standard Model

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1 Introduction

Inflation [1–6] is well known to solve the flatness and horizon problems, diluting cosmological relics and providing an origin of cosmological fluctuations. In slow-roll inflation [7, 8], the inflaton rolls along a quite flat potential and inflation end as it falls into some basin. Inflation is supported and constrained by current observational data [9], which measures a spectral index $n_s \approx 0.96 \pm 0.007$ and low tensor-to-scalar ratio $r < 0.08$, excluding the simplest chaotic models based on polynomial potentials such as ϕ^2 or ϕ^4 [10]. Surviving models include Starobinsky inflation [3, 11, 12], Higgs inflation [13] and related models [14, 15], and low scale hybrid inflation [16, 17].

Supersymmetry (SUSY) may be naturally combined with inflation since it allows better control over the high energy dynamics of scalars [18–20]. SUSY inflation is also motivated by the Lyth bound [21] on the low tensor-to-scalar ratio, which prefers a scale of inflation below the Planck scale. Since inflation is sensitive to UV scales, it is necessary to consider supergravity (SUGRA) inflation, as in e.g. [22–42]. In no-scale SUGRA [44], the Kähler potential takes a logarithmic form which circumvents the η problem. Alternatives to no-scale SUGRA have also been proposed which also address the η problem based on a Heisenberg symmetry [45–48] or a shift symmetry [49–51] (see also [52–57]).

It has been shown by Ellis, Nanopoulos, Olive (ENO) that no-scale SUGRA can behave like a Starobinsky inflationary model [58–60]. However, in this approach, a term with constant modular weight is used to break SUSY, and there is no connection between

inflation and SUSY breaking. Recently we considered the above ENO model, but with a linear Polonyi term added to the superpotential [61]. The purpose of adding this term was to provide an explicit mechanism for breaking SUSY in order to provide a link between inflation and SUSY breaking. Indeed we showed that inflation requires a strict upper bound for the gravitino mass $m_{3/2} < 10^3$ TeV [61].

In the present paper we show how the Polonyi-extended ENO model may be generalised to include the fields in the visible sector of the minimal supersymmetric standard model (MSSM). Such a generalisation has been done for the ENO model [58–60] and we perform a similar analysis for the Polonyi-extended ENO model. We calculate the soft-SUSY breaking parameters depending on the modular weights in the superpotential and choice of Kähler potential and we are led to new phenomenological possibilities for supersymmetry (SUSY) breaking, based on generalisations of no-scale SUSY breaking and pure gravity mediated SUSY breaking. The Polonyi-extended ENO model is especially interesting to consider because of the upper bound on the gravitino mass discussed in the previous paragraph which allows the much lighter gauginos to be discovered in future collider experiments. This motivates the present investigation of the soft SUSY breaking parameters, which could form the basis for future phenomenological studies.

The layout of the remainder of the paper is as follows. In section 2 we discuss the hidden sector of the supergravity theory where inflation takes place. In section 3 we discuss the visible sector of the supergravity theory and show how the MSSM matter and Higgs fields may be included. In section 4 we discuss the supergravity scalar potential, showing how inflation emerges from the hidden sector and soft-SUSY breaking parameters emerge from the full theory including the visible sector, leading to new examples of no-scale SUSY breaking and pure gravity mediated SUSY breaking. Section 5 concludes the paper.

2 The hidden sector

In general supergravity theory, the tree-level supergravity scalar potential can be found using the Kähler function G , which is given in terms of the Kähler potential K and the superpotential W as,

$$G = \frac{K}{M_P^2} + \ln \left| \frac{W}{M_P^3} \right| + \ln \left| \frac{W}{M_P^3} \right|^* . \tag{2.1}$$

The effective scalar potential is then given by,

$$V = e^G \left[\frac{\partial G}{\partial \phi_i} K_{ij^*} \frac{\partial G}{\partial \phi_{j^*}} - 3 \right] M_P^4, \tag{2.2}$$

where K_{ij^*} is the inverse of the Kähler metric $K^{ij^*} \equiv \partial^2 K / \partial \phi_i \partial \phi_{j^*}$. When, at the minimum of the scalar potential, some of the hidden sector fields acquire VEVs in such a way that at least one of their auxiliary fields, F^i , is non-vanishing, then SUSY is spontaneously broken and soft SUSY-breaking terms are generated in the observable sector. The gravitino becomes massive and its mass

$$m_{3/2}^2 = e^G = e^{K/M_P^2} \frac{|W|^2}{M_P^4} \tag{2.3}$$

sets the overall scale of the soft parameters. In general, the Kähler potential K and the superpotential W involve superfields in the hidden and visible sectors. The hidden sector superfields are gauge singlets and do not have Yukawa couplings to the charged fields in the visible sector, coupling only indirectly to them via Planck suppressed operators.

The simplest no-scale Kähler potential in the hidden sector is given by two complex fields (T, ϕ) , where T is a modulus field while ϕ is the field responsible for inflation and SUSY breaking. The Kähler potential in the hidden sector takes the form

$$K_{\text{hid}}(\phi, T) = -3M_P^2 \ln \left(\frac{T + T^*}{M_P} - \frac{|\phi|^2}{3M_P^2} \right), \quad (2.4)$$

where M_P is the reduced Planck scale.

It was found in [58] this Kähler potential together with the Wess-Zumino superpotential [62, 63] can lead to the Starobinsky-like inflationary potential. When the modulus field T is fixed with a vacuum expectation value of $\langle \text{Re } T \rangle = 1/2$ and $\langle \text{Im } T \rangle = 0$, the no-scale Kähler potential together with the Wess-Zumino superpotential is equivalent of an $R + R^2$ model of gravity, in which Starobinsky inflation emerges at a particular point in parameter space [59]. A simple modification to this superpotential has been done in [61], adding the Polonyi term to provide an explicit and simple mechanism for supersymmetry breaking at the end of inflation. The Wess-Zumino superpotential [62] in the hidden sector, with quadratic and trilinear terms, together with a linear Polonyi term looks like

$$W_{\text{hid}}(\phi) = m^2 \phi + \frac{\mu}{2} \phi^2 - \frac{\lambda}{3} \phi^3. \quad (2.5)$$

In the following, it is convenient to introduce the change of variables [64]

$$T = \frac{M_P}{2} \left(\frac{1 - \frac{y_2}{\sqrt{3}M_P}}{1 + \frac{y_2}{\sqrt{3}M_P}} \right), \quad \phi = \left(\frac{y_1}{1 + \frac{y_2}{\sqrt{3}M_P}} \right), \quad (2.6)$$

with the inverse relations

$$y_1 = \left(\frac{2\phi}{1 + 2T/M_P} \right), \quad y_2 = \sqrt{3}M_P \left(\frac{1 - 2T/M_P}{1 + 2T/M_P} \right). \quad (2.7)$$

After this change of variables, the hidden sector superpotential becomes

$$W_{\text{hid}}(y_1, y_2) = \left(1 + \frac{y_2}{\sqrt{3}M_P} \right)^{-3} \widetilde{W}_{\text{hid}}(y_1, y_2), \quad (2.8)$$

where the rescaled superpotential is

$$\widetilde{W}_{\text{hid}}(y_1, y_2) = m^2 y_1 \left(1 + \frac{y_2}{\sqrt{3}M_P} \right)^2 + \frac{\mu}{2} y_1^2 \left(1 + \frac{y_2}{\sqrt{3}M_P} \right) - \frac{\lambda}{3} y_1^3. \quad (2.9)$$

The Kähler potential in the hidden sector becomes

$$K_{\text{hid}}(y_1, y_2) = \widetilde{K}_{\text{hid}}(y_1, y_2) + 3 \ln \left(\left| 1 + \frac{y_2}{\sqrt{3}M_P} \right|^2 \right) \quad (2.10)$$

where

$$\tilde{K}_{\text{hid}}(y_1, y_2) = -3M_P^2 \ln \left(1 - \frac{|y_1|^2 + |y_2|^2}{3M_P^2} \right), \quad (2.11)$$

Note that the combination of \tilde{W}_{hid} and \tilde{K}_{hid} is equivalent to using W_{hid} and K_{hid} , since the physical quantities are given by the Kähler function G (2.1) which is the same in both cases (the extra term in the Kähler potential cancels with an opposite term coming from the superpotential). Therefore, we use the symmetric representation of the Kähler \tilde{K}_{hid} in eq. (2.11) and the rescaled superpotential \tilde{W}_{hid} in eq. (2.9) in the following.

This superpotential can reproduce the Starobinsky model for the real part of y_1 and fixed $\langle y_2 \rangle = 0$, when a suitable stabilizing term $\propto y_2^4$ is added to the no-scale Kähler potential [59, 64, 65]:

$$\tilde{K}_{\text{hid}}(y_1, y_2) = -3M_P^2 \ln \left(1 - \frac{|y_1|^2 + |y_2|^2}{3M_P^2} + \frac{|y_2|^4}{\Lambda^2 M_P^2} \right), \quad (2.12)$$

with $\Lambda \lesssim 0.1M_P$, as discussed in [59]. In section 3, we ignore the term $\propto y_2^4$ for simplicity, but it should be kept in mind that it is necessary to stabilize the field $\langle y_2 \rangle = 0$.

The exact Starobinsky potential is obtained when dropping the Polonyi term, $m = 0$, fixing $\langle y_2 \rangle = 0$ and using the relationship $\lambda = \mu/\sqrt{3}M_P$ [59]. To quantify how much the Starobinsky limit deviates when including the Polonyi term, we use the parameter $b = m^2/3\lambda M_P^2$ while keeping the relation $\lambda = \mu/\sqrt{3}M_P$, which gives

$$\mu^{-1} \tilde{W}_{\text{hid}}(y_1, y_2) = \sqrt{3}bM_P y_1 \left(1 + \frac{y_2}{\sqrt{3}M_P} \right)^2 + \frac{y_1^2}{2} \left(1 + \frac{y_2}{\sqrt{3}M_P} \right) - \frac{y_1^3}{3\sqrt{3}M_P}. \quad (2.13)$$

For non-zero b , SUSY is broken and the gravitino mass in eq. (2.3) becomes non-zero at the end of inflation, as was shown in [61]. As discussed in [61], there is an upper limit on the parameter $|b|$ in order to have a viable inflationary scenario, suggesting a gravitino mass $m_{3/2} < 10^3$ TeV with favoured values of $m_{3/2} \sim \mathcal{O}(1)$ TeV. A more quantitative discussion (which we do not repeat here) can be found in [61], where it is shown that this limit comes from the two crucial dimensionless observables: the tensor-to-scalar ratio, r , and the scalar tilt, n_s , given by the Planck satellite [9]. Furthermore, the scalar amplitude observable, A_s , is sensitive to the overall scale of the potential, *i.e.* to the parameter μ in eq. (2.13). It is shown there [61] that, in the Starobinsky limit when $\lambda = \mu/\sqrt{3}M_P$, the bilinear mass term parameter becomes $\mu \simeq 10^{-5}M_P$ and we use these values in the following computations.

3 The visible sector

We now extend the hidden sector inflation model in the previous section to include the matter fields in the visible sector, such as those of the minimal supersymmetric standard model (MSSM), which includes the Standard Model (SM) quarks and leptons together with their supersymmetric partners, written generically in terms of matter and Higgs superfields $y_{\text{vis}} = \hat{Q}_i, \hat{U}_i^c, \hat{D}_i^c, \hat{L}_i, \hat{E}_i^c, \hat{H}_u, \hat{H}_d$. The visible sector superfields carry gauge charges under the SM gauge group, unlike the hidden sector superfields which are gauge singlets. We

consider two possibilities for extending the hidden sector supergravity theory of the previous section to include such visible sector superfields.

Motivated by the no-scale approach [66, 67], the first possibility (Case I) is to embed the visible sector matter superfields within the logarithm in the Kähler potential (Case I), such that

$$K_I = -3M_P^2 \ln \left(1 - \frac{|y_1|^2 + |y_2|^2 + |y_{\text{vis}}|^2}{3M_P^2} \right), \quad \text{Case I.} \quad (3.1)$$

Another possibility (Case II) we also explore is to have the visible sector superfields outside the logarithm in the Kähler potential via minimal kinetic terms

$$K_{II} = -3M_P^2 \ln \left(1 - \frac{|y_1|^2 + |y_2|^2}{3M_P^2} \right) + |y_{\text{vis}}|^2, \quad \text{Case II.} \quad (3.2)$$

Furthermore, in the same spirit as in [60], we assume that the superpotential for the visible sector superfields has the form

$$W_{\text{vis}} = W_2(y_{\text{vis}}) \left(1 + \frac{y_2}{\sqrt{3}M_P} \right)^\beta + W_3(y_{\text{vis}}) \left(1 + \frac{y_2}{\sqrt{3}M_P} \right)^\alpha, \quad (3.3)$$

where α and β are modular weights and $W_{2,3}(y_{\text{vis}})$ are bi/trilinear parts of the superpotential of the MSSM in terms of the visible sector superfields y_{vis} . The total superpotential is then given by

$$W = \widetilde{W}_{\text{hid}} + W_{\text{vis}}, \quad (3.4)$$

where $\widetilde{W}_{\text{hid}}$ was given in eq. (2.13) and W_{vis} is given in eq. (3.3).

We want to make the connection between the Wess-Zumino-Polonyi superpotential (2.13), which has a parameter space in which the Starobinsky inflation is recovered, and the soft supersymmetry-breaking parameters after the MSSM superfields are included. The main difference between the model of [61] which we are developing here, and that proposed in [60] is that, in the present case, supersymmetry is broken through the Polonyi term while in [60] a term with constant modular weight 3 term was used to break SUSY. In the present model there is a constraint from inflation in how big the parameter b (accounting for the Polonyi term) can be, leading to an upper bound on the gravitino mass, which sets the SUSY breaking scale. The present model therefore suggests a more constrained region of parameter space which may be confronted with LHC, Higgs and dark matter constraints.

The main difference between the two cases I and II is that in the pure no-scale case (3.1), the soft supersymmetry-breaking mass squared m_0 is zero while in the latter case (3.2), m_0 is different from zero and equal to the gravitino mass, as we will show in section 4.

4 Potential and soft-SUSY breaking parameters

We can compute the supergravity scalar potential from eqs. (2.1) and (2.2), using either eq. (3.1) (case I) or eq. (3.2) (case II), with the total superpotential in eq. 3.4. At the minimum of the potential some of the hidden sector fields, y_1 and y_2 , acquire VEVs such

that the F term is non-vanishing, then supersymmetry is spontaneously broken and soft supersymmetry-breaking terms are generated in the observable sector, y_{vis} . This is just the usual gravity mediated SUSY breaking mechanism. However the special forms of superpotential and Kähler potential here lead to a special form of SUSY breaking, referred to as either no-scale SUSY breaking (case I) or pure gravity mediated SUSY breaking (case II).

4.1 Hidden sector potential and inflation

We begin by computing the supergravity scalar potential from eqs. (2.1) and (2.2), in the hidden sector using $\widetilde{W}_{\text{hid}}$ (2.13) and $\widetilde{K}_{\text{hid}}$ (2.11) only. We follow the same treatment as in [58] and assume that the T field is fixed with a vacuum expectation value of $\langle \text{Re}T \rangle = c/2$ and $\langle \text{Im}T \rangle = 0$, corresponding to $\langle y_2 \rangle = 0$.¹ The minimum of the potential is always given by $V = 0$, for both cases I and II, $V = 0$ is found for

$$y_1 = \frac{\sqrt{3}}{2} \left(1 \pm \sqrt{1 + 4b} \right) M_P, \tag{4.1}$$

such that supersymmetry is spontaneously broken in the hidden sector. For $b = 0$, $y_1 = 0$ only if we have the minus sign and we restrict ourselves to this case in the following. This model reproduces the effective potential of the Starobinsky model for $\mathcal{R}e y_1$. The dynamical field y_1 can be converted into a canonically-normalized inflaton field x by the transformation [58, 59]

$$y_1 = \pm \sqrt{3} M_P \tanh \left(\frac{\chi}{\sqrt{3} M_P} \right) = \pm \sqrt{3} M_P \tanh \left(\frac{x}{\sqrt{6} M_P} \right), \tag{4.2}$$

where $\chi = (x + iy)/\sqrt{2}$ and the latter equality holds for $y = 0$. The imaginary part of the inflaton is fixed to $y = 0$ by the potential [58, 61] since the potential is always minimized by $y = 0$ in the range of interest of the inflaton field x . We use the positive sign in eq. (4.2) and write the potential in terms of the inflaton field x

$$V = \frac{3}{4} \mu^2 M_P^2 \left[1 + b + (b - 1) \cosh \left(\sqrt{\frac{2}{3}} \frac{x}{M_P} \right) + \sinh \left(\sqrt{\frac{2}{3}} \frac{x}{M_P} \right) \right]^2, \tag{4.3}$$

where we have fixed $\langle y_2 \rangle = 0$. In terms of the inflaton field x , the minimum of the potential is found for

$$x_0 = \sqrt{6} M_P \tanh^{-1} \left(\frac{1}{2} (1 - \sqrt{1 + 4b}) \right). \tag{4.4}$$

The exact Starobinsky limit is realized for $b = 0$,² while small values for b represents a small deviation from the Starobinsky limit as shown in figure 1.

As b deviates from zero, the value of the field at the global minimum x_0 shifts away from zero, while maintaining $V = 0$. Although the scalar potential vanishes, the superpotential

¹The modulus field y_2 is stabilized when including a term $\propto y_2^4$ as in eq. (2.12), for a better understanding see [59] and section 4.1.1.

²For $b = 0$, the scalar potential can be written as $V = \frac{3}{4} \mu^2 M_P^2 \left(1 - e^{-\sqrt{2/3} x/M_P} \right)^2$, which is exactly the Starobinsky potential [11, 12].

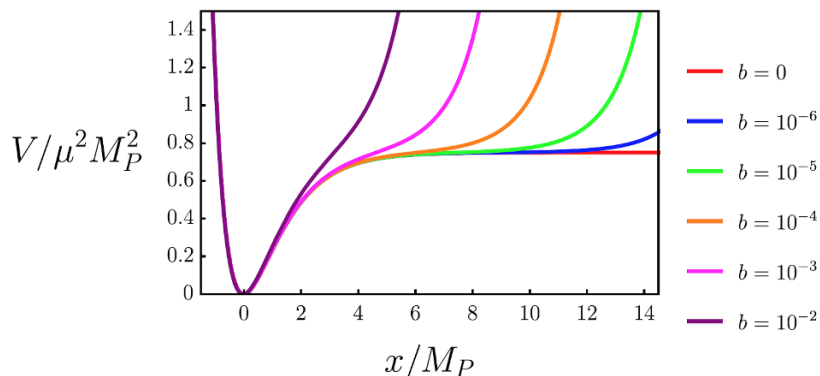


Figure 1. Potential for different values of b . For $b = 0$ the exact Starobinsky limit is obtained and the potential has a flat plateau. For $b \sim 10^{-5}$ the potential retains a quite flat region, while for $b \gtrsim 10^{-4}$ the potential loses its flatness.

W and the Kähler function G is different from zero, generating explicit soft supersymmetry-breaking terms of the required form in the effective low-energy Lagrangian. The gravitino becomes massive, see eq. (2.3), and sets the overall scale of the soft parameters

$$m_{3/2} = -\frac{3}{2}\mu b^2 + \mathcal{O}(b^3), \tag{4.5}$$

where we have expanded at lowest order in b . The exact analytical expression can be found in appendix A.

As explained in [61], the limits from inflation suggest that the parameters $\mu \simeq 10^{-5}M_P \simeq 2 \times 10^{13}$ GeV and $b \lesssim 10^{-4}$. For fixed $x_* = 5.35M_P$ (where x_* is the value of the field when inflation starts) we see that there is an approximate quadratic dependence of the gravitino mass on the parameter b , as shown in figure 2, where we have rescaled the results around the origin as in [61]. For $b = 0$, we see that $m_{3/2} = 0$, so supersymmetry is unbroken and we are left with the Wess-Zumino superpotential limit leading to Starobinsky inflation. However, from figure 1, we see that for small values of $b \lesssim 10^{-4}$, the potential retains a plateau where inflation will happen and the corresponding limit on the gravitino mass may be read off from figure 2 as $m_{3/2} < 10^3$ TeV.

4.1.1 Stabilizing the modulus field

We introduced a term in the Kähler potential $\propto y_2^4/\Lambda^2$ to assure the stabilization of the field y_2 during inflation, see eq. (2.12). Additionally, we can compute the mass of the modulus field y_2 , m_{y_2} , and the mass of the field y_1 , m_{y_1} , during inflation. As a benchmark point, we choose $x = 5M_P$ and $\Lambda = 0.01M_P$, and we find that the masses are $m_{y_1} \sim 10^{13}$ GeV and $m_{y_2} \sim 10^{17}$ GeV. The fact that $m_{y_2} \gg m_{y_1}$ is valid, not only for this benchmark point, but during the whole inflationary trajectory for $\Lambda \lesssim 0.1M_P$, means that the single field approximation is justified during inflation.

At the end of inflation, when y_1 is at its minimum (see eq. (4.1)), the masses of the fields y_1 and y_2 are given by

$$m_{y_1} = \sqrt{2}\mu(1 - b) + \mathcal{O}(b^2), \quad m_{y_2} = \sqrt{2}\mu b + \mathcal{O}(b^2). \tag{4.6}$$

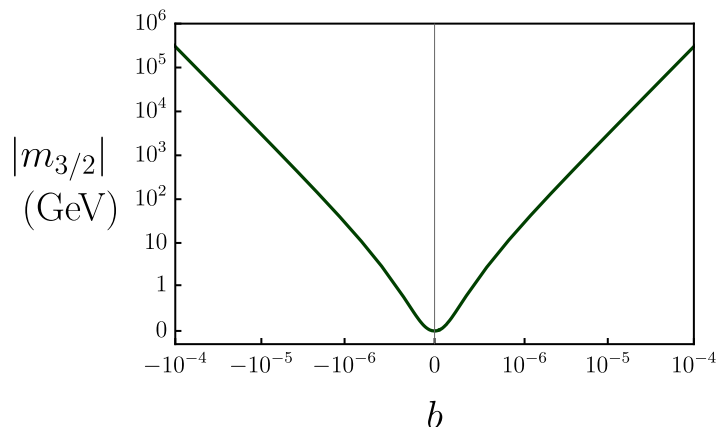


Figure 2. Gravitino mass as a function of the parameter b accounting for the Polonyi term. As shown in the previous figure 1 (for details see [61]), successful inflation requires $b \lesssim 10^{-4}$, corresponding here to a strict upper bound for the gravitino mass $m_{3/2} < 10^3$ TeV.

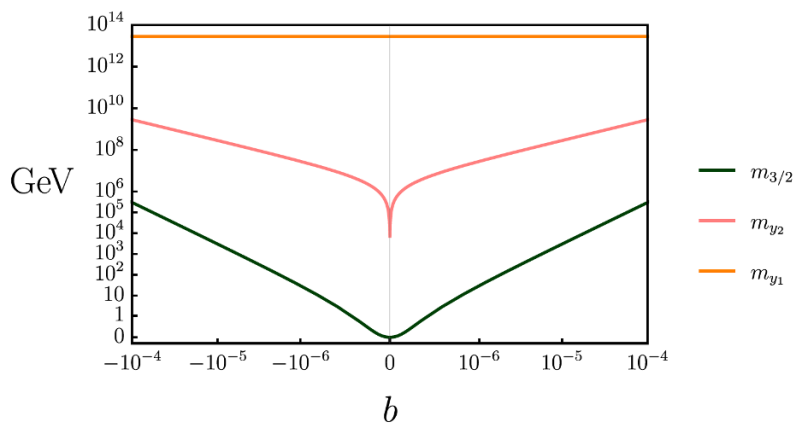


Figure 3. Masses of the gravitino, the modulus field y_2 and the field y_1 , for different values of b , at the end of inflation. The masses are evaluated when y_1 is at its minimum, given by eq. (4.1). The modulus field is strongly stabilized with a mass $m_{y_2} \gg m_{3/2}$.

Figure 3 shows the gravitino mass as well as the mass of the modulus field y_2 and the mass of the field y_1 . The modulus y_2 is strongly stabilized at the end of inflation, where $m_{y_2} \gg m_{3/2}$.

4.2 Visible sector potential and SUSY breaking

We now compute the supergravity scalar potential from eqs. (2.1) and (2.2), including the visible sector superfields, using either eq. (3.1) (case I) or eq. (3.2) (case II), with the total superpotential in eq. 3.4. With the addition of the visible sector superpotential in eq. (3.3), we are now able to compute the soft supersymmetry breaking mass-squared, bilinear, and trilinear parameters, m_0, B_0 and A_0 respectively.

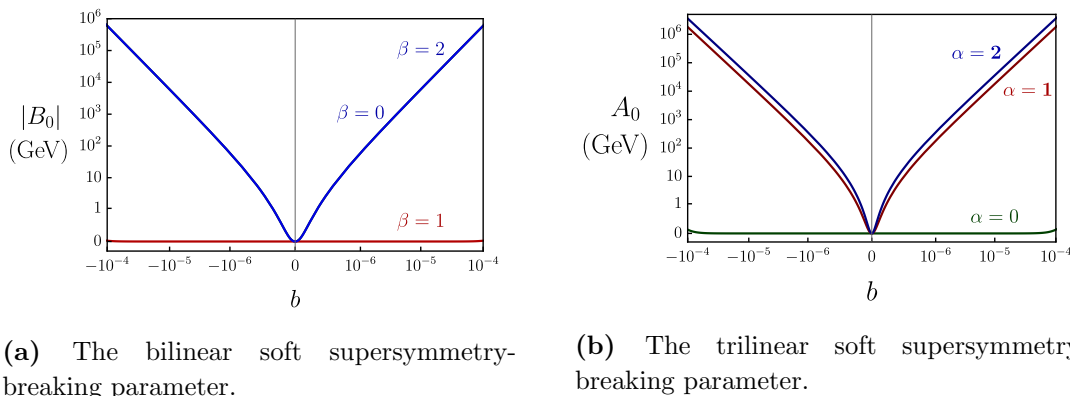


Figure 4. Soft supersymmetry-breaking parameters as a function of the parameter b , accounting for the Polonyi term, for different modular weights α and β .

4.2.1 Case I: no-scale SUSY breaking

For the Case I, where the SM superfields are inside the Log of the Kähler potential, the soft supersymmetry breaking parameters become

$$\begin{aligned}
 m_0 &= 0, \\
 \frac{A_0}{m_{3/2}} &= -6\alpha - 3(4 + \alpha)b^2 + \mathcal{O}(b^3), \quad \text{Case I} \\
 \frac{B_0}{m_{3/2}} &= 2(1 - \beta) - (1 + \beta)b^2 + \mathcal{O}(b^3).
 \end{aligned} \tag{4.7}$$

The prediction for $m_0 = 0$ is the familiar result of no-scale SUSY breaking. The exact analytical functions are given in appendix A, while the expressions in eq. (4.7) are found when expanding in powers of b and hold for our range of interest $b \lesssim 10^{-4}$. The mass-squared term m_0 is zero while the bilinear and trilinear parameters will depend on the choice of the modular weights α and β , as shown in figure 4. The main effect of switching on b is to increase the gravitino mass, since the terms proportional to b^2 are negligible. The special choice of $\alpha = 0$ and $\beta = 1$, corresponds to the pure no-scale option where $m_0 = B_0 = A_0 = 0$ for $b \lesssim 10^{-4}$, in which case the supersymmetry breaking in the low scale energy can be produced via a non-minimal gauge kinetic term generating non-zero gaugino masses $M_{1/2} \neq 0$.

For the rest of cases, where $\alpha \neq 0$ and $\beta \neq 1$, we can set some limits in the bilinear and trilinear couplings due to the constraint on the parameter $b \lesssim 10^{-4}$ coming from inflation such that $B_0 < 10^3$ TeV and $A_0 \lesssim 10^3$ TeV. We assume the condition $m_{3/2} \gg M_{1/2}$ to avoid constraints from Big Bang nucleosynthesis on gravitino decays. In this case, we cannot find a viable phenomenological region due to the relation between the trilinear coupling and the gravitino mass $A_0 \simeq -6\alpha m_{3/2}$ in eq. (4.7), implying $A_0 \gg M_{1/2}$ which leads to tachyonic sfermion masses. Furthermore, for $\beta \neq 1$, there is no value of $\tan\beta$ which minimizes the Higgs potential and which could give the correct Higgs Mass $m_H = 125 \pm 0.24$ GeV. Therefore, the only case we can further explore is the case $\alpha = 0$ and $\beta = 1$, corresponding to $m_0 = A_0 = B_0 = 0$.

It has been proven in [68] the no-scale boundary conditions $m_0 = A_0 = B_0 = 0$ when becoming universal at some unification scale M_{in} above the GUT scale, are compatible with low-energy constraints. They show a scenario based on SU(5) model where the superpotential contains terms $W \ni \lambda H \Sigma \bar{H} + (\lambda'/6) \text{Tr} \Sigma^3$, where H, \bar{H} and Σ are **5**, $\bar{\mathbf{5}}$ and **24** Higgs representations. For different values of λ and λ' , the region of M_{in} , $M_{1/2}$ and $\tan \beta$ has been studied, taking into account phenomenological constraints on supersymmetric particles and the cosmological LSP density. Although, the LHC has imposed additional constraints, via the measurement of the Higgs Mass and the decay $B_s \rightarrow \mu^+ \mu^-$, the model is still consistent with the LHC data [60, 69, 70]. For example, for the values $\lambda = -0.1$ and $\lambda' = 2$, the region with $M_{1/2} \in (1000, 1500)$ GeV and $M_{\text{in}} \in (10^{17}, 10^{18})$ GeV is consistent with the relic cold dark matter density, a Higgs mass of $m_H \sim 125$ GeV and the experimental measurements of $B_s \rightarrow \mu^+ \mu^-$.

4.2.2 Case II: pure gravity SUSY breaking

As an alternative to the no-scale possibility, we can have the Standard Model superfields y_{vis} outside the logarithm. This is called Case II in eq. (3.2). For small b , where inflation works, we find that the soft supersymmetry breaking parameters are

$$\begin{aligned} m_0 &= m_{3/2}, \\ \frac{A_0}{m_{3/2}} &= -6\alpha - 3(4 + \alpha)b^2 + \mathcal{O}(b^3), \\ \frac{B_0}{m_{3/2}} &= 2(1 - \beta) - (1 + \beta)b^2 + \mathcal{O}(b^3), \end{aligned} \quad \text{Case II} \tag{4.8}$$

where the relations for A_0 and B_0 are kept the same as in eq. (4.7) while the universal soft scalar mass term m_0 is now different from zero and equal to the gravitino mass $m_{3/2}$. Note that the soft scalar mass is equal to the gravitino mass, while the soft gaugino mass will be assumed to be much smaller, as in anomaly mediated SUSY breaking. As before, the main effect of switching on b is to increase the gravitino mass, since the terms proportional to b^2 are negligible.

For the same reasons as before, the general cases $\alpha \neq 0$ and $\beta \neq 1$ are not viable phenomenological choices. As mentioned, we impose $m_{3/2} \gg M_{1/2}$, meaning $A_0 \gg M_{1/2}$ (see eq. (4.8)). This will lead to tachyonic sfermion masses. Similarly, for $\beta \neq 1$ we cannot find a value of $\tan \beta$ which reproduces the correct Higgs Mass. Therefore, we will focus in the case $\alpha = 0$ and $\beta = 1$ on the following, corresponding to $m_0 = m_{3/2}$ and $A_0 = B_0 = 0$.

The special case $\alpha = 0$ and $\beta = 1$, meaning $A_0 \approx B_0 \approx 0$, is called pure gravity-mediated (PGM) models [71–79] where the gaugino masses, A and B terms are determined by anomaly mediation [80–84] leaving only the gravitino mass $m_{3/2} = m_0$ as a free parameter. Note that in this case since $b \lesssim 10^{-4}$, eq. (4.8) gives $A_0 = B_0 = 0$ to excellent approximation. The main phenomenological effect of the Polonyi term is then to set a limit on the gravitino mass from the requirement of successful inflation. These PGM models with GUT scale universality [74] are phenomenologically viable when including a Giudice-Masiero term in the Kähler potential [85, 86], which allows to choose $\tan \beta$ as a second free parameter. For successful electroweak symmetry breaking (EWSB),

$\tan\beta$ is restricted to a narrow range from about 1.7 – 2.5 [75, 79]. Then, the Higgs mass $m_H \sim 125$ GeV is obtained when the gravitino mass is in the range 300 – 1500 TeV. In our case we have the additional constraint from inflation requiring $m_{3/2} < 1000$ TeV, see figure 2. Therefore, the Case II, in which $m_0 = m_{3/2}$, for $\alpha = 0$ and $\beta = 1$ has an allowed parameter space when including the Giudice-Masiero term which is also compatible with inflation and corresponds to $b \sim 10^{-4}$.

Recently, PGM models have been studied in an SU(5) GUT model [87], allowing two new parameters, a high energy scale above the GUT scale, M_{in} and a new coupling $\lambda\bar{H}\Sigma H$, where Σ is an SU(5) adjoint Higgs. They find a viable parameter space as long as the Higgs coupling λ is relatively large, for a relaxed $\tan\beta$, away from $\tan\beta \sim 2$, opening up the parameter space considerably. In the case where higher dimensional operators involving are included, proton decay can be within reach of future experiments. Moreover, in some regions of parameter space the bino can be degenerate with the wino or gluino, giving an acceptable dark matter relic density.

5 Conclusion

In this paper we have studied a version of Starobinsky-like inflation in no-scale SUGRA where a Polonyi term in the hidden sector breaks SUSY after inflation, providing a link between the gravitino mass and inflation. The linear Polonyi term provides a simple way to break SUSY after inflation, with the requirement of successful inflation leading to an upper bound on the gravitino mass $m_{3/2} < 10^3$ TeV, with gaugino masses considerably less than this.

The progress made in this paper is to extend the previous hidden sector Polonyi theory (proposed by one of us) to include the visible sector and to calculate the resulting soft-SUSY breaking parameters in terms of the modular weights α and β for the bilinear and trilinear parts of the superpotential and a choice of Kähler potential. In this case, we can write the soft-SUSY breaking parameters in terms of α , β and the parameter b accounting for the Polonyi term. In case I we included the MSSM matter and Higgs superfields inside the Log of the Kähler potential following a no-scale approach. In the second case, Case II, these superfields are added via minimal kinetic terms.

We were thereby led to familiar phenomenological possibilities for supersymmetry (SUSY) breaking, based on no-scale SUSY breaking and pure gravity mediated SUSY breaking, but with the strict bound coming from our Polonyi model that $m_{3/2} < 10^3$ TeV. However we found that the only phenomenological viable choices for both cases I and II, are the ones with $\alpha = 0$ and $\beta = 1$, corresponding to $A_0 = B_0 = 0$.

In case I, in addition to $A_0 = B_0 = 0$, the scalar soft mass term is zero $m_0 = 0$, as in the no-scale SUSY breaking approach, in which case SUSY breaking at low energies is produced via non-zero gaugino masses $M_{1/2} \neq 0$. Models like this one had been discussed by ENO and have been shown to be compatible with LHC data, with all superpartners potentially observable at LHC or the FCC.

In case II, in addition to $A_0 = B_0 = 0$, the scalar mass is equal to the gravitino mass $m_0 = m_{3/2}$, with $M_{1/2} \ll m_{3/2}$ as in anomaly mediation models, with such a scenario

referred to as PGM. Such PGM models are viable models after the inclusion of a Giudice-Masiero term in the Kähler potential, for a large gravitino mass of order 100–1000 TeV, compatible with our inflation bound, and can evade LHC searches while still providing a good dark matter candidate and gauge coupling unification, with squarks and sleptons very heavy, while the gauginos remain light and observable at the LHC or the FCC.

In conclusion, Starobinsky-like inflation in no-scale SUGRA models with a Polonyi term provides a promising setting for both inflation and SUSY breaking within a well motivated particle physics framework. The Polonyi term provides a link between the gravitino mass and inflation leading to a strict upper bound on the gravitino mass $m_{3/2} < 10^3$ TeV. We have seen that under reasonable assumptions, the soft-SUSY breaking parameters can be calculated, leading to either no-scale SUGRA or PGM patterns, where the results presented here could form the basis of future phenomenological studies. In particular, since gauginos are significantly lighter than $m_{3/2}$, this suggests that SUSY could be discovered at the LHC or FCC.

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A Soft-supersymmetry breaking parameters

We present the analytic results for the gravitino mass and the soft-supersymmetry breaking parameters as a function of the modular weights α and β and the parameter b , which takes into account the contribution from the Polonyi term in the superpotential. For both cases I and II, in which we include the Standard Model superfields inside or outside the logarithm in the Kähler potential, see eqs. (3.1) and (3.2), the bilinear, B_0 and trilinear A_0 parameters are the same and the only difference is in the soft scalar mass term m_0 , vanishing in case I while being equal to the gravitino mass in case II.

$$\begin{aligned}
 m_{3/2} &= -\mu \frac{-1 + \sqrt{1 + 4b} + b(-6 + 4\sqrt{1 + 4b})}{\sqrt{2}(1 - 2b + \sqrt{1 + 4b})^{3/2}} \\
 m_0 &= 0 \text{ (Case I) } \quad \text{or} \quad m_0 = m_{3/2} \text{ (Case II) } , \\
 A_0 &= -\mu \frac{3(\sqrt{4b + 1} - 1)^2(-8b + \sqrt{4b + 1} - 1)}{\sqrt{2}(-2b + \sqrt{4b + 1} + 1)^3} \\
 &\quad \times \frac{(b(\alpha - 8(\alpha - 2)b + 5\alpha\sqrt{4b + 1} - 10\sqrt{4b + 1} + 14) - 2\sqrt{4b + 1} + 2)}{\sqrt{(-\sqrt{4b + 1} + 2b(-5\sqrt{4b + 1} + b(16b - 12\sqrt{4b + 1} + 21) + 6) + 1)}}, \\
 B_0 &= -\mu \frac{(\sqrt{4b + 1} - 1)^2(-8b + \sqrt{4b + 1} - 1)}{(-2b + \sqrt{4b + 1} + 1)^3} \\
 &\quad \times \frac{(b(-8b(\beta - 2) + 5\beta\sqrt{4b + 1} + \beta - 10\sqrt{4b + 1} + 6) - \sqrt{4b + 1} + 1)}{\sqrt{-2\sqrt{4b + 1} + 4b(-5\sqrt{4b + 1} + b(16b - 12\sqrt{4b + 1} + 21) + 6) + 2}}. \tag{A.1}
 \end{aligned}$$

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