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SUSY breaking after inflation in supergravity with inflaton in a massive vector supermultiplet

Yermek Aldabergenov^a, Sergei V. Ketov^{a,b,c,d,*}^a Department of Physics, Tokyo Metropolitan University, Minami-ohsawa 1-1, Hachioji-shi, Tokyo 192-0397, Japan^b Kavli Institute for the Physics and Mathematics of the Universe (IPMU), The University of Tokyo, Chiba 277-8568, Japan^c Department of Physics, Faculty of Science, Chulalongkorn University, Thanon Phayathai, Pathumwan, Bangkok 10330, Thailand^d Institute of Physics and Technology, Tomsk Polytechnic University, 30 Lenin Ave., Tomsk 634050, Russian Federation

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ABSTRACT

We propose a limited class of models, describing interacting chiral multiplets with a non-minimal coupling to a vector multiplet, in curved superspace of $N = 1$ supergravity. Those models are suitable for the inflationary model building in supergravity with inflaton assigned to a massive vector multiplet and spontaneous SUSY breaking in Minkowski vacuum after inflation, for any values of the inflationary parameters n_s and r , and any scale of SUSY breaking.

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1. Introduction

Success of the inflationary scenario for early Universe is, on the one hand, due to overcoming the theoretical problems (horizon, flatness, structure formation) of the standard (Einstein–Friedmann) cosmology and, on the other hand, due to its remarkable agreement with the CMB observational data (COBE, WMAP, PLANCK). For instance, the observed breaking of CMB scale invariance is measured by the scalar tilt, $n_s - 0.9666 = \pm 0.0062$ [1,2], and the relative magnitude of primordial gravity waves is parametrized by the tensor-to-scalar ratio $r < 0.07$ [3]. Those observations favour chaotic slow-roll inflation in its single-field realization, i.e. the large-field inflation driven by a single scalar called *inflaton* with an approximately flat scalar potential.

Embedding a single-field inflation into $N = 1$ four-dimensional supergravity is needed to connect inflationary models to particle physics beyond the Standard Model, and towards their ultimate embedding into string theory. It requires inflaton to belong to a massive $N = 1$ multiplet that can be either a chiral multiplet (of the highest spin 1/2) or a real vector multiplet (of the highest spin 1). Most of the literature about inflation in supergravity uses the first option – see e.g., the reviews [4,5] – since it is usually assumed that vector fields do not play any role during

inflation.¹ However, assuming inflaton to be in a chiral multiplet also causes some problems. First, the scalar component of a chiral multiplet is *complex*, which implies the need to stabilize another (non-inflaton) scalar during inflation. Second, there is also the so-called *η -problem* caused by the presence of the exponential factor e^K in the scalar potential of supergravity with chiral superfields, which generically prevents slow roll. Third, there are problems also with ensuring the inflaton scalar potential to be bounded from below, and with getting SUSY breaking in a Minkowski vacuum after inflation too. Of course, the inflationary model building in supergravity now has many models that overcome some of those problems – see e.g., [6–10] and references therein. However, it often comes at the price of having more matter superfields together with a need to invent the dynamics for them. The *minimal* inflationary models with a single inflaton chiral superfield, with or without SUSY breaking after inflation, are also possible [11–13] but require tuning both Kähler potential and a superpotential. Yet another approach, based on the use of non-linear realizations of SUSY and nilpotent chiral superfields, was introduced to the supergravity-based inflationary model building in [14].

When inflaton is assigned to a massive *vector* multiplet, there is no need of its complexification, because the scalar field component of a real massive $N = 1$ vector multiplet is *real*. Accordingly, there is no need for other scalars and their stabilization during

* Corresponding author.

E-mail addresses: aldabergenov-yermek@ed.tmu.ac.jp (Y. Aldabergenov), ketov@tmu.ac.jp (S.V. Ketov).¹ Taking into account vector fields is believed to be important after inflation, during reheating.

inflation, in the minimal supergravity setup. The η -problem also does not arise because the scalar potential of a vector multiplet in supergravity has a different structure (of the D -type instead of the F -type). Actually, the corresponding minimal inflationary models were already constructed by Ferrara, Kallosh, Linde and Porrati in [15] by exploiting the non-minimal self-coupling of a vector multiplet to supergravity, found by Van Proeyen in [16].

The supergravity inflationary models of [15] have the single-field scalar potential given by an arbitrary real function squared. Those scalar potentials are always bounded from below and allow any desired values of n_s and r . However, the minima of the scalar potentials of [15] have the vanishing cosmological constant and the vanishing VEV of the auxiliary field D , so that they only have Minkowski vacua where supersymmetry is always restored after inflation. It is desirable to have more theoretical flexibility, as regards SUSY breaking, for phenomenological purposes.

In this paper we propose a simple extension of the inflationary models [15] by adding a *Polonyi* (chiral) superfield [17]. Our models also can accommodate arbitrary values of n_s and r , but have a Minkowski vacuum after inflation, with spontaneously broken supersymmetry (SUSY).

Our paper is organized as follows. In Sec. 2 we propose a new class of supergravity models in curved superspace of $N = 1$ old-minimal supergravity. Our models can be considered as the extensions of those in [15] via adding a chiral superfield and its coupling to a vector (inflaton) superfield in supergravity. We also compute the bosonic kinetic terms and the scalar potential in our models. In Sect. 3 we identify the chiral sector with the Polonyi model, and find a Minkowski vacuum with spontaneously broken supersymmetry after inflation that is not affected by the Polonyi superfield. Sect. 4 is our Conclusion.

2. A vector multiplet non-minimally coupled to a chiral multiplet in supergravity

Let us consider some chiral superfields Φ_i with arbitrary Kähler potential $K = K(\Phi_i, \bar{\Phi}_i)$ and a chiral superpotential $\mathcal{W} = \mathcal{W}(\Phi_i)$, interacting with a real superfield V whose arbitrary potential is described by a real function $J = J(V)$. The real vector superfield V is supposed to describe a massive vector multiplet, while the chiral superfields are supposed to be (gauge) singlets in our construction.

We employ the curved superspace formalism of $N = 1$ supergravity [18]. Our notation and conventions coincide with the standard ones in [18], including the spacetime signature $(-, +, +, +)$.²

Our models are defined by the Lagrangian ($M_{\text{Pl}} = 1$)

$$\mathcal{L} = \int d^2\theta 2\mathcal{E} \left\{ \frac{3}{8}(\bar{D}\bar{D} - 8\mathcal{R})e^{-\frac{1}{3}(K+2J)} + \frac{1}{4}W^\alpha W_\alpha + \mathcal{W} \right\} + \text{h.c.}, \quad (1)$$

where we have introduced the chiral density superfield $2\mathcal{E}$, the chiral scalar curvature superfield \mathcal{R} , and the chiral vector superfield strength $W_\alpha \equiv -\frac{1}{4}(\bar{D}\bar{D} - 8\mathcal{R})D_\alpha V$.

In order to calculate the bosonic part of our models, we set all fermions to zero, and define the bosonic field components of the relevant superfields. As regards the supergravity multiplet, we have

$$2\mathcal{E}| = e, \quad D\mathcal{D}(2\mathcal{E})| = 4\bar{M},$$

$$\mathcal{R}| = -\frac{1}{6}M, \quad D\mathcal{D}\mathcal{R}| = -\frac{1}{3}R + \frac{4}{9}M\bar{M} + \frac{2}{9}b_m b^m - \frac{2}{3}iD_m b^m,$$

² The $N = 1$ superconformal calculus used in [15,16] is equivalent to the curved superspace description [18] of $N = 1$ Poincaré supergravity after the superconformal gauge fixing.

where we have introduced the vierbein determinant $e \equiv \det e_m^a$, the spacetime scalar curvature R , and the old-minimal set of the supergravity auxiliary fields, the complex scalar M and the real vector b_m . The vertical bars denote the leading field components of a superfield at $\theta = \bar{\theta} = 0$.

The field components of Φ_i and V are defined by

$$\Phi_i| = A_i \quad D_\alpha D_\beta \Phi_i| = -2\varepsilon_{\alpha\beta} F_i, \quad \bar{D}_{\dot{\alpha}} D_\alpha \Phi_i| = -2\sigma_{\alpha\dot{\alpha}}{}^m \partial_m A_i,$$

$$\bar{D}\bar{D}D\mathcal{D}\Phi_i| = 16\Box A_i + \frac{32}{3}ib_a \partial^a A_i + \frac{32}{3}F_i M,$$

$$V| = C \quad D_\alpha D_\beta V| = \varepsilon_{\alpha\beta} X, \quad \bar{D}_{\dot{\alpha}} D_\alpha V| = \sigma_{\alpha\dot{\alpha}}{}^m (B_m - i\partial_m C),$$

$$D_\alpha W^\beta| \equiv -\frac{1}{4}D_\alpha(\bar{D}\bar{D} - 8\mathcal{R})D^\beta V \\ = \frac{1}{4}\sigma_{\alpha\dot{\alpha}}{}^m \bar{\sigma}^{\dot{\alpha}\beta n} (D_m \partial_n C + 2iF_{mn}) + \delta_{\alpha}{}^\beta D,$$

$$\bar{D}\bar{D}D\mathcal{D}V| = \frac{16}{3}b^m (B_m - i\partial_m C) + 6\Box C - \frac{16}{3}MX + 8D,$$

in terms of the physical fields A_i , C , B_m as complex scalars, a real scalar, and a real vector respectively, the chiral auxiliary fields F_i and X as complex scalars, the real auxiliary field D as a real scalar, and the vector field strength $F_{mn} = D_m B_n - D_n B_m$ of B_m .

Using those definitions, we find by a straightforward calculation that the kinetic part of our Lagrangian is given by

$$e^{-1}\mathcal{L}_{\text{kin.}} = e^{-\frac{1}{3}(K+2J)} \left\{ -\frac{1}{2}R + \frac{1}{2}K_i \Box A_i + \frac{1}{2}K_i{}^* \Box \bar{A}_i \right. \\ - \frac{1}{6}K_i K_j \partial_m A_i \partial^m A_j - \frac{1}{6}K_i{}^* K_j{}^* \partial_m \bar{A}_i \partial^m \bar{A}_j \\ - \left(\frac{1}{3}J'^2 - \frac{1}{2}J'' \right) \partial_m C \partial^m C + \left(\frac{1}{3}J'^2 - \frac{1}{2}J'' \right) B_m B^m \\ + \frac{3}{4}J' \Box C + \frac{i}{3}J' B_m (K_i{}^* \partial^m \bar{A}_i - K_i \partial^m A_i) \\ \left. - \frac{1}{3}J' \partial_m C (K_i{}^* \partial^m \bar{A}_i + K_i \partial^m A_i) \right\} - \frac{1}{4}F_{mn} F^{mn}, \quad (2)$$

while its auxiliary part reads

$$e^{-1}\mathcal{L}_{\text{aux.}} = e^{-\frac{1}{3}(K+2J)} \left\{ \frac{1}{3}b_m b^m + \frac{i}{3}b_m (K_i{}^* \partial^m \bar{A}_i - K_i \partial^m A_i) \right. \\ + \frac{2}{3}J' b_m B^m + J'D + K_{ij}{}^* F_i \bar{F}_j - \left(\frac{1}{3}J'^2 - \frac{1}{2}J'' \right) X \bar{X} \\ - \frac{1}{3}(M\bar{M} + K_i K_j{}^* F_i \bar{F}_j - J'K_i{}^* \bar{F}_i X - J'K_i F_i \bar{X} \\ + K_i{}^* \bar{F}_i \bar{M} + K_i F_i M - J'MX - J'\bar{M}\bar{X}) \\ \left. - \frac{1}{4}D\Box C + \frac{1}{2}D^2 + F_i \mathcal{W}_i + \bar{F}_i \bar{\mathcal{W}}_i - \bar{M}\mathcal{W} - M\bar{\mathcal{W}} \right\}. \quad (3)$$

In our equations above, the K , J and \mathcal{W} now represent the lowest components of the corresponding superfields, being functions of the scalar fields A_i and C . As regards their derivatives, we have used the notation $K_i \equiv \frac{\partial K}{\partial A_i}$, $K_i{}^* \equiv \frac{\partial K}{\partial \bar{A}_i}$, $K_{ij}{}^* \equiv \frac{\partial^2 K}{\partial A_i \partial \bar{A}_j}$, $J' \equiv \frac{\partial J}{\partial C}$,

$$\mathcal{W}_i \equiv \frac{\partial \mathcal{W}}{\partial A_i}, \quad \bar{\mathcal{W}}_i \equiv \frac{\partial \bar{\mathcal{W}}}{\partial \bar{A}_i}.$$

In order to eliminate the auxiliary fields in accordance to their algebraic equations of motion, we first separate M , F_i and X from each other via a substitution,

$$M = N + J'\bar{X} - K_i{}^* \bar{F}_i, \quad (4)$$

$$\bar{M} = \bar{N} + J'X - K_i F_i. \quad (5)$$

In terms of the new auxiliary fields N and \bar{N} , the auxiliary part of the Lagrangian takes the form

$$e^{-1}\mathcal{L}_{\text{aux.}} = e^{-\frac{1}{3}(K+2J)} \left\{ \frac{1}{3}b_m b^m + \frac{i}{3}b_m(K_i^* \partial^m \bar{A}_i - K_i \partial^m A_i) \right. \\ \left. + \frac{2}{3}J' b_m B^m + J'D + K_{ij} F_i \bar{F}_j + \frac{1}{2}J'' X \bar{X} - \frac{1}{3}N \bar{N} \right\} \\ - \frac{1}{4}D \square C + \frac{1}{2}D^2 + F_i \mathcal{W}_i + \bar{F}_i \bar{\mathcal{W}}_i \\ - \mathcal{W}(\bar{N} + J'X - K_i F_i) - \bar{\mathcal{W}}(N + J'\bar{X} - K_i^* \bar{F}_i), \quad (6)$$

so that Euler–Lagrange equations of the auxiliary fields are easily solved as

$$b_m = -J' B_m - \frac{i}{2}(K_i^* \partial_m \bar{A}_i - K_i \partial_m A_i), \\ D = \frac{1}{4}\square C - J'e^{-\frac{1}{3}(K+2J)}, \quad N = -3e^{\frac{1}{3}(K+2J)}\mathcal{W}, \\ F_i = -e^{\frac{1}{3}(K+2J)}K_{ij}^{-1}(\bar{\mathcal{W}}_j + K_j^* \bar{\mathcal{W}}), \quad X = 2\frac{J'}{J''}e^{\frac{1}{3}(K+2J)}\bar{\mathcal{W}}.$$

After a substitution of those solutions back into the Lagrangian, we find

$$e^{-1}\mathcal{L} = e^{-\frac{1}{3}(K+2J)} \left\{ -\frac{1}{2}R + \frac{1}{2}K_i \square A_i + \frac{1}{2}K_i^* \square \bar{A}_i \right. \\ \left. - \frac{1}{6}K_i K_j^* \partial_m A_i \partial^m \bar{A}_j - \frac{1}{12}K_i K_j \partial_m A_i \partial^m A_j \right. \\ \left. - \frac{1}{12}K_i^* K_j^* \partial_m \bar{A}_i \partial^m \bar{A}_j - \left(\frac{1}{3}J'^2 - \frac{1}{2}J'' \right) \partial_m C \partial^m C \right. \\ \left. + J' \square C - \frac{1}{3}J' \partial_m C (K_i^* \partial^m \bar{A}_i + K_i \partial^m A_i) - \frac{1}{2}J'' B_m B^m \right\} \\ - \frac{1}{4}F_{mn} F^{mn} - \frac{1}{2}e^{-\frac{2}{3}(K+2J)} J'^2 \\ - e^{\frac{1}{3}(K+2J)} \left[K_{ij}^{-1} (\mathcal{W}_i + K_i \mathcal{W}) (\bar{\mathcal{W}}_j + K_j^* \bar{\mathcal{W}}) \right. \\ \left. - \left(3 - 2\frac{J'^2}{J''} \right) \mathcal{W} \bar{\mathcal{W}} \right]. \quad (7)$$

A transition from Jordan to Einstein frame is achieved by Weyl rescaling of spacetime metric,

$$g_{mn} \rightarrow e^\Lambda g_{mn}, \quad e \rightarrow e^{2\Lambda} e, \quad \text{with } \Lambda = \frac{1}{3}(K+2J).$$

Then the scalar curvature term transforms as

$$-\frac{1}{2}e e^{-\frac{1}{3}(K+2J)} R \rightarrow -\frac{1}{2}e R + \frac{1}{12}(\partial_m K + 2\partial_m J)^2. \quad (8)$$

It gives rise to the Lagrangian

$$e^{-1}\mathcal{L} = -\frac{1}{2}R + \frac{1}{2}K_i \square A_i + \frac{1}{2}K_i^* \square \bar{A}_i - \frac{1}{4}F_{mn} F^{mn} \\ - \frac{1}{2}J'' \partial_m C \partial^m C - \frac{1}{2}J'' B_m B^m - \mathcal{V}, \quad (9)$$

with the scalar potential

$$\mathcal{V} = \frac{1}{2}J'^2 + e^{K+2J} \left[K_{ij}^{-1} (\mathcal{W}_i + K_i \mathcal{W}) (\bar{\mathcal{W}}_j + K_j^* \bar{\mathcal{W}}) \right. \\ \left. - \left(3 - 2\frac{J'^2}{J''} \right) \mathcal{W} \bar{\mathcal{W}} \right]. \quad (10)$$

Equations (1), (9) and (10) are our main results in this Section. When the real superfield V is dropped ($J=0$), our result coincides with the standard Lagrangian and the scalar potential of chiral superfields in $N=1$ supergravity [19]. When all the chiral superfields Φ_i are dropped ($K=\mathcal{W}=0$), our results coincide with those in [15,16].³

As is clear from (9), the absence of ghosts requires $J''(C) > 0$.

3. Vacuum solution

In this Section we restrict ourselves to a single chiral superfield Φ having the canonical Kähler potential and the superpotential given by a sum of a linear term and a constant,

$$K = \Phi \bar{\Phi}, \quad \mathcal{W} = \mu(\Phi + \beta). \quad (11)$$

This particular choice is known in the literature as *Polonyi model* [17].⁴

In accordance to the previous Section, it gives rise to the Lagrangian

$$e^{-1}\mathcal{L} = -\frac{1}{2}R - \partial_m A \partial^m \bar{A} - \frac{1}{4}F_{mn} F^{mn} - \frac{1}{2}J'' \partial_m C \partial^m C \\ - \frac{1}{2}J'' B_m B^m - \frac{1}{2}J'^2 - \mu^2 e^{A\bar{A}+2J} \\ \times \left[|1 + A\beta + A\bar{A}|^2 - \left(3 - 2\frac{J'^2}{J''} \right) |A + \beta|^2 \right]. \quad (12)$$

The (Minkowski) vacuum conditions in this model are given by

$$V = \frac{1}{2}J'^2 + \mu^2 e^{A\bar{A}+2J} \\ \times \left[|1 + A\beta + A\bar{A}|^2 - \left(3 - 2\frac{J'^2}{J''} \right) |A + \beta|^2 \right] = 0, \quad (13)$$

$$\partial_{\bar{A}} V = A \tilde{V}_F + \mu^2 e^{A\bar{A}+2J} \left[A(1 + \bar{A}\beta + A\bar{A}) + (A + \beta) \right. \\ \left. \times (1 + A\beta + A\bar{A}) - \left(3 - 2\frac{J'^2}{J''} \right) (A + \beta) \right] = 0, \quad (14)$$

$$\partial_C V = J' \left\{ J'' + 2\mu^2 e^{A\bar{A}+2J} \left[|1 + A\beta + A\bar{A}|^2 \right. \right. \\ \left. \left. - \left(1 - 2\frac{J'^2}{J''} + \frac{J'J'''}{J''^2} \right) |A + \beta|^2 \right] \right\} = 0, \quad (15)$$

where we have introduced \tilde{V}_F as the F-type scalar potential with the additional J -dependent term as

$$\tilde{V}_F = \mu^2 e^{A\bar{A}+2J} \left[|1 + A\beta + A\bar{A}|^2 - \left(3 - 2\frac{J'^2}{J''} \right) |A + \beta|^2 \right]. \quad (16)$$

A simple solution to those equations exist when $J'=0$, which separates the Polonyi multiplet from the vector multiplet. The remaining vacuum equations allow a solution with the VEV $\langle A \rangle \equiv \alpha = (\sqrt{3}-1)$ and $\beta = 2 - \sqrt{3}$ [17]. This celebrated (Polonyi) solution describes a stable Minkowski vacuum with spontaneously broken SUSY since $\langle F \rangle = \mu$. Hence, the parameter μ defines the scale of SUSY breaking, which is *arbitrary* in this model. The related gravitino mass is given by $m_{3/2} = \mu e^{2-\sqrt{3}}$. There is also a massive scalar of mass $2m_{3/2}$ and a massless fermion in the physical spectrum.

It should be emphasized that the Polonyi field does not affect inflation associated with the scalar C as the inflaton belonging to the massive vector multiplet, and having the D -type scalar potential $V(C) = \frac{1}{2}J'^2$ with arbitrary real J -function. Of course, the true inflaton field should be canonically normalized via the appropriate field redefinition of C .

When trying to get other patterns of SUSY breaking after inflation by demanding $J' \neq 0$ and $\alpha = \beta = 0$, we get two conditions on the J -function,

⁴ It is worth mentioning that this choice is most natural for a nilpotent (Akulov–Volkov) superfield, $\Phi^2 = 0$.

³ Our notation for J differs by the sign from that of [15,16].

$$J'^2 = J'' , \quad (17)$$

$$J'' = -2\mu^2 e^{2J} . \quad (18)$$

The first equation is solved by $J = -\log C + \text{const.}$, then the second condition yields the consistency relation $\text{const.} = -\frac{1}{2} \log(-2\mu^2)$. Since both J and μ should be real, there is no solution. However, when allowing $\beta \neq 0$, the second equation (18) gets modified as

$$J'' = -2\mu^2 e^{2J} (1 - \beta^2) , \quad (19)$$

so that the reality of J and μ requires $\beta > 1$. Then (17) reads $J'^2 = C^{-2}$ and is easily solvable. However, such scalar potential is not suitable for inflation (no slow roll). More general vacuum solutions with $J' \neq 0$ will be investigated elsewhere.

4. Conclusion

Our basic equations (1), (9) and (10) supply new theoretical tools for the inflationary model building in supergravity. They can be further generalized e.g., by including an extra function $g(\Phi)$ of the chiral superfields in front of the vector multiplet kinetic term in (1), and/or replacing the Maxwell-type kinetic term of the vector multiplet by the Born–Infeld-type action, like e.g., in [20].

Our models have three arbitrary (input) potentials K , \mathcal{W} and J , providing more flexibility to the inflationary model building and, perhaps, being derivable from a more fundamental theory, like string theory.

Our construction does not have an R-symmetry, and is apparently unrelated to (the dual version of) matter-coupled ($R + R^2$) supergravity in its ‘new-minimal’ formulation [21].

In particular, as was demonstrated in Sec. 3, our construction easily supplies spontaneous SUSY breaking after inflation to the supergravity-based inflationary models whose inflaton belongs to a massive vector multiplet, via their coupling to Polonyi multiplet. Those models are *limited* in the sense that they provide the *minimal* extension of the inflationary models proposed in [15] for the sake of spontaneous SUSY breaking in Minkowski vacuum after inflation. The Polonyi multiplet itself can be assigned to the hidden sector needed for SUSY breaking and its gravitational mediation to the visible sector in more general field-theoretical models of particle physics beyond the Standard Model in the context of supergravity.

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