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

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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.

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

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[1] Taking into account vector fields is believed to be important after inflation, during reheating.

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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 $K=K(\Phi_i,\overline{\Phi}_i)$ 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 ($M_{\rm Pl}=1$)

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

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

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

$$2\mathcal{E}|=e,\quad \mathcal{D}\mathcal{D}(2\mathcal{E})|=4\overline{M}\,,$$
$$\mathcal{R}|=-\tfrac{1}{6}M,\quad \mathcal{D}\mathcal{D}\mathcal{R}|=-\tfrac{1}{3}R+\tfrac{4}{9}M\overline{M}+\tfrac{2}{9}b_mb^m-\tfrac{2}{3}i\mathcal{D}_mb^m\,,$$

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

$$\Phi_i|=A_i\quad \mathcal{D}_\alpha\mathcal{D}_\beta\Phi_i|=-2\varepsilon_{\alpha\beta}F_i\,,\quad \overline{\mathcal{D}}_{\dot\alpha}\mathcal{D}_\alpha\Phi_i|=-2\sigma_{\alpha\dot\alpha}{}^m\partial_mA_i\,,$$
$$\overline{\mathcal{D}}\overline{\mathcal{D}}\mathcal{D}\mathcal{D}\Phi_i|=16\Box A_i+\tfrac{32}{3}ib_a\partial^aA_i+\tfrac{32}{3}F_iM\,,$$
$$V|=C\quad \mathcal{D}_\alpha\mathcal{D}_\beta V|=\varepsilon_{\alpha\beta}X\,,\quad \overline{\mathcal{D}}_{\dot\alpha}\mathcal{D}_\alpha V|=\sigma_{\alpha\dot\alpha}{}^m(B_m-i\partial_mC)\,,$$
$$\begin{aligned}\mathcal{D}_\alpha W^\beta|&\equiv-\tfrac{1}{4}\mathcal{D}_\alpha(\overline{\mathcal{D}}\overline{\mathcal{D}}-8\mathcal{R})\mathcal{D}^\beta V\\&=\tfrac{1}{4}\sigma_{\alpha\dot\alpha}{}^m\bar\sigma^{\dot\alpha\beta n}(\mathcal{D}_m\partial_nC+2iF_{mn})+\delta_\alpha{}^\beta D\,,\end{aligned}$$
$$\overline{\mathcal{D}}\overline{\mathcal{D}}\mathcal{D}\mathcal{D}V|=\tfrac{16}{3}b^m(B_m-i\partial_mC)+6\Box C-\tfrac{16}{3}MX+8D\,,$$

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 $F_{mn}=\mathcal{D}_mB_n-\mathcal{D}_nB_m$ of B_m.

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

$$\begin{aligned}e^{-1}\mathcal{L}_{\rm kin.}=e^{-\frac{1}{3}(K+2J)}\Big\{&-\tfrac{1}{2}R+\tfrac{1}{2}K_i\Box A_i+\tfrac{1}{2}K_{i^*}\Box\bar A_i\\&-\tfrac{1}{6}K_iK_j\partial_mA_i\partial^mA_j-\tfrac{1}{6}K_{i^*}K_{j^*}\partial_m\bar A_i\partial^m\bar A_j\\&-\left(\tfrac{1}{3}J'^2-\tfrac{1}{2}J''\right)\partial_mC\partial^mC+\left(\tfrac{1}{3}J'^2-\tfrac{1}{2}J''\right)B_mB^m\\&+\tfrac{3}{4}J'\Box C+\tfrac{i}{3}J'B_m(K_{i^*}\partial^m\bar A_i-K_i\partial^mA_i)\\&-\tfrac{1}{3}J'\partial_mC(K_{i^*}\partial^m\bar A_i+K_i\partial^mA_i)\Big\}-\tfrac{1}{4}F_{mn}F^{mn}\,,\end{aligned}\tag{2}$$

while its auxiliary part reads

$$\begin{aligned}e^{-1}\mathcal{L}_{\rm aux.}=e^{-\frac{1}{3}(K+2J)}\Big\{&\tfrac{1}{3}b_mb^m+\tfrac{i}{3}b_m(K_{i^*}\partial^m\bar A_i-K_i\partial^mA_i)\\&+\tfrac{2}{3}J'b_mB^m+J'D+K_{ij^*}F_i\overline{F}_j-\left(\tfrac{1}{3}J'^2-\tfrac{1}{2}J''\right)X\overline{X}\\&-\tfrac{1}{3}(M\overline{M}+K_iK_{j^*}F_i\overline{F}_j-J'K_{i^*}\overline{F}_iX-J'K_iF_i\overline{X}\\&+K_{i^*}\overline{F}_i\overline{M}+K_iF_iM-J'MX-J'\overline{M}\overline{X})\Big\}-\tfrac{1}{4}D\Box C\\&+\tfrac{1}{2}D^2+F_i\mathcal{W}_i+\overline{F}_i\overline{\mathcal{W}}_i-\overline{M}\mathcal{W}-M\overline{\mathcal{W}}\,.\end{aligned}\tag{3}$$

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 $K_i\equiv\frac{\partial K}{\partial A_i}$, $K_{i^*}\equiv\frac{\partial K}{\partial\bar A_i}$, $K_{ij^*}\equiv\frac{\partial^2K}{\partial A_i\partial\bar A_j}$, $J'\equiv\frac{\partial J}{\partial C}$, $\mathcal{W}_i\equiv\frac{\partial\mathcal{W}}{\partial A_i}$, $\overline{\mathcal{W}}_i\equiv\frac{\partial\overline{\mathcal{W}}}{\partial\bar A_i}$.

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,

$$M=N+J'\overline{X}-K_{i^*}\overline{F}_i\,, \tag{4}$$
$$\overline{M}=\overline{N}+J'X-K_iF_i\,. \tag{5}$$

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

[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.