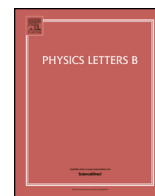


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On the one-loop effective potential in the higher-derivative four-dimensional chiral superfield theory with a nonconventional kinetic term

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ABSTRACT

We explicitly calculate the one-loop effective potential for a higher-derivative four-dimensional chiral superfield theory with a nonconventional kinetic term. We consider the cases of minimal and nonminimal general Lagrangians. In particular, we find that in the minimal case the divergent part of the one-loop effective potential vanishes by reason of the chirality.

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1. Introduction

Higher-derivative field theories have aroused great interest by many reasons. The main of them is certainly the fact that the presence of higher derivatives improves the renormalization behavior. This called the attention to these theories, in the gravity context where higher-derivative additive terms were used to construct renormalizable theories of gravity [1]. Further studies of gravity with additive higher-derivative terms, which naturally emerge within the anomaly context [2], led to the possibility of supersymmetric extensions. The first example of such theories is the dilaton supergravity [3] based on the results for the supertrace anomaly [4]. In its essence, the dilaton supergravity, in which the dynamics is completely described by a nontrivial prepotential, so-called chiral compensator [5], was the first higher-derivative superfield theory discussed in detail at the quantum level.

Further, the equivalence of higher-derivative (super)field theories to theories without higher derivatives but with additional ghost fields has been shown at the tree level for several types of models [6]. Therefore, the interest to the general properties of the effective action of higher-derivative superfield theories arose. Another important reason for the interest in higher-derivative field theories is the concept of the Horava–Lifshitz gravity [7] which involves higher derivatives, although only in the spatial sector, and without ghosts.

In the papers [8] a higher-derivative theory with a conventional kinetic term (in which the higher derivatives enter the general Lagrangian in the form $\int d^8z \Phi \square \bar{\Phi}$) has been studied, the one-loop low-energy effective action has been considered, and its superrenormalizability (and, in certain cases, finiteness) has been shown. Further, similar studies have been done for the four-dimensional supergauge theory [9] and for the three-dimensional superfield theories [10].

In this paper, we extend these studies by obtaining the one-loop effective potential for the chiral superfield theory with a nonconventional kinetic term. It is characterized by the presence of higher derivatives in a classical chiral Lagrangian. Up to now, such a theory was treated only in the paper [6] where it was studied on the tree level in the context of equivalence with theories without higher derivatives but with additional ghost fields. Its quantum properties never have been studied. This is just the problem we address in this paper.

For this study, we use the methodology of obtaining the superfield effective potential originally developed in [11] (this method has been successfully applied in [8,9,12,13] for different superfield models). Throughout this paper, we use notation and conventions from the book [14].

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2. Minimal general Lagrangian

Let us start with the following 4D higher-derivative superfield theory where the higher derivatives are implemented in a nonconventional way, that is, in the chiral Lagrangian:

$$S = \int d^8z \Phi \bar{\Phi} + \left[\int d^6z \left(-\frac{a}{2} \Phi \square \Phi + W(\Phi) \right) + \text{h.c.} \right]. \quad (1)$$

Here a is a constant with a negative (-1) mass dimension, and the chiral potential is given by an arbitrary holomorphic function $W(\Phi)$ which in some cases will be particularized as $W(\Phi) = \frac{1}{2}m\Phi^2 + \frac{\lambda}{3!}\Phi^3$. We remind that earlier, theories of such a form have been studied only at the tree level [6], whereas the quantum calculations were carried out only for higher derivatives in the general effective Lagrangian, that is, just the case discussed in details in [8].

Before performing any quantum calculation, it is convenient to determine the superficial degree of divergence ω for the model (1). First, we have to calculate the propagators for (1). They are given by

$$\langle \bar{\Phi}(1)\Phi(2) \rangle = \frac{i}{k^2 + (m + ak^2)^2} \delta_{12}, \quad (2)$$

$$\langle \Phi(1)\Phi(2) \rangle = i \frac{(m + ak^2)D_1^2}{4k^2[k^2 + (m + ak^2)^2]} \delta_{12}, \quad (3)$$

$$\langle \bar{\Phi}(1)\bar{\Phi}(2) \rangle = i \frac{(m + ak^2)\bar{D}_1^2}{4k^2[k^2 + (m + ak^2)^2]} \delta_{12}, \quad (4)$$

where $m = W''(\Phi)|_{\Phi=0}$. Second, we have to determine the number of covariant derivatives in each of the supergraphs. In our case, the number is the same as that of the Wess–Zumino model, namely $4V - 2E + 2C - 4L$, where V , E , P , C , and L denote the number of vertices, external lines, $\langle \bar{\Phi}(1)\Phi(2) \rangle$ -propagators, chiral propagators, and loops, respectively.

Finally, we assume that all covariant derivatives are converted into momenta via the D -algebra. Then, it follows that the maximal superficial degree of divergence is given by

$$\begin{aligned} \omega_{\max} &= 4L - 4P - 4C + \frac{1}{2}(4V - 2E + 2C - 4L) \\ &= 2 - 2P - C - E, \end{aligned} \quad (5)$$

where we used the topological identity $L + V - P - C = 1$. Of course, for any non-trivial supergraph contribution to the effective action we will have $E \geq 2$ and P or C are no less than 1. Therefore, the theory (1) is finite which is natural for higher-derivative models (cf. [3]).

Now, let us begin with the calculation of the one-loop effective potential. In order to find the one-loop effective action we use a trick originally introduced in [11] to transform the integral over chiral superfields to the integral over unconstrained superfields, that is, we introduce the effective action W for the free gauge superfield v :

$$e^{iW} = \int Dv \delta\left(\frac{1}{4}D^2v - \bar{\phi}\right) \delta\left(\frac{1}{4}\bar{D}^2v - \phi\right) \exp\left(-\frac{i}{16} \int d^8zv D^\alpha \bar{D}^2 D_\alpha v\right). \quad (6)$$

It is clear that W is a constant (indeed, it is gauge independent, so, it does not depend on the gauge fixing functions, and, thus, on ϕ and $\bar{\phi}$). Then, let us do a background-quantum splitting in the action (1), by the rule $\Phi \rightarrow \Phi + \phi$, $\bar{\Phi} \rightarrow \bar{\Phi} + \bar{\phi}$, where ϕ , $\bar{\phi}$ are from now on the background chiral and antichiral fields. The quadratic part in the quantum fields of the action (1) is

$$S_2 = \int d^8z \phi \bar{\phi} + \left[\frac{1}{2} \int d^6z (-a\phi \square \phi + \Psi \phi^2) + \text{h.c.} \right], \quad (7)$$

where $\Psi = W''(\Phi)$ is a background chiral field. The one-loop effective action corresponding to the theory (1) is $\Gamma^{(1)}$ defined by

$$e^{i\Gamma^{(1)}[\phi, \bar{\phi}]} = \int D\phi D\bar{\phi} \exp(iS_2) = \int D\phi D\bar{\phi} \exp\left(i \left[\int d^8z \phi \bar{\phi} + \left[\frac{1}{2} \int d^6z (-a\phi \square \phi + \Psi \phi^2) + \text{h.c.} \right] \right] \right). \quad (8)$$

Let us multiply this expression with Eq. (6). Using functional delta functions and disregarding irrelevant constants, we get

$$e^{i\Gamma^{(1)}[\phi, \bar{\phi}]} = \int D\phi D\bar{\phi} \exp(iS_2) = \int Dv \exp\left[i \frac{1}{2} v \int d^8z \left(\square - \frac{a}{4} \square (D^2 + \bar{D}^2) + \frac{1}{4} (\Psi \bar{D}^2 + \bar{\Psi} D^2) \right) v \right]. \quad (9)$$

So, we find that the desired one-loop effective action is

$$\Gamma^{(1)} = -\frac{i}{2} \text{Tr} \ln \left[\square + \frac{1}{4} (-a\square + \Psi) \bar{D}^2 + \frac{1}{4} (-a\square + \bar{\Psi}) D^2 \right]. \quad (10)$$

So, it remains to calculate this trace of the logarithm. To do it, we expand it in power series:

$$\Gamma^{(1)} = -\frac{i}{2} \text{Tr} \sum_{n=1}^{\infty} \frac{(-1)^{n+1}}{n} \left[\frac{(-a\square + \Psi) \bar{D}^2 + (-a\square + \bar{\Psi}) D^2}{4\square} \right]^n. \quad (11)$$

It is easy to see that, for A and \bar{A} constant and arbitrary (this is just the case sufficient for the calculation of the superfield effective potential where the background superfields are constant [11]),

$$\text{Tr} \left[\frac{1}{4} (\bar{A} D^2 + A \bar{D}^2) \right]^{2k} = \int d^8 z (A \bar{A} \square)^k \frac{2}{\square} \delta^4(x - x') \Big|_{x=x'}. \tag{12}$$

For an odd number of multipliers the trace is zero. Thus,

$$\Gamma^{(1)} = \frac{i}{2} \int d^8 z \sum_{k=0}^{\infty} \frac{1}{k} \left[\frac{(-a \square + \Psi)(-a \square + \bar{\Psi})}{\square} \right]^k \frac{1}{\square} \delta^4(x - x') \Big|_{x=x'}, \tag{13}$$

after Fourier transform, Wick rotation and summation, we arrive at the following contribution to the Kählerian effective potential

$$K^{(1)}(\bar{\Phi}, \Phi) = -\frac{1}{2} \int \frac{d^4 p}{(2\pi)^4} \frac{1}{p^2} \ln \left[1 + \frac{(ap^2 + \Psi)(ap^2 + \bar{\Psi})}{p^2} \right]. \tag{14}$$

It is clear that at $a = 0$, we restore the well-known expression $K^{(1)}(\bar{\Phi}, \Phi) = -\frac{1}{2} \int \frac{d^4 p}{(2\pi)^4} \frac{1}{p^2} \ln(1 + \frac{\Psi \bar{\Psi}}{p^2})$ (it was explicitly found in [11] where it was shown that in this case, after subtraction of the UV divergences, one has $K^{(1)}(\bar{\Phi}, \Phi) = -\frac{1}{32\pi^2} \Psi \bar{\Psi} \ln \frac{\Psi \bar{\Psi}}{\mu^2}$). So, it remains to do the integral in (14). This task can be done by splitting the logarithm in three parts,

$$K^{(1)}(\bar{\Phi}, \Phi) = -\frac{1}{2} \int \frac{d^4 p}{(2\pi)^4} \frac{1}{p^2} [\ln(p^2 + \Omega_+) + \ln(p^2 + \Omega_-) - \ln(p^2)], \tag{15}$$

where

$$\Omega_{\pm} = \frac{1 + a(\bar{\Psi} + \Psi) \pm \sqrt{[1 + a(\bar{\Psi} + \Psi)]^2 - 4a^2 \bar{\Psi} \Psi}}{2a^2}. \tag{16}$$

The last logarithm does not depend on the background superfields, then we can drop it by means of the normalization of the effective action. The integrals in (15) are well-known and can be computed by using the dimensional reduction prescription. Therefore, we regularize this integral by the formal replacement of $d^4 p$ by $\mu^{4-2\omega} d^{2\omega} p$, so, one has

$$K^{(1)}(\bar{\Phi}, \Phi) = -\frac{1}{2} \mu^{4-2\omega} \int \frac{d^{2\omega} p}{(2\pi)^{2\omega}} \frac{1}{p^2} [\ln(p^2 + \Omega_+) + \ln(p^2 + \Omega_-)]. \tag{17}$$

In the limit $\omega \rightarrow 2$ we find

$$K^{(1)}(\bar{\Phi}, \Phi) = \frac{1}{32\pi^2(2-\omega)} (\Omega_+ + \Omega_-) - \frac{1}{32\pi^2} \left[\Omega_+ \ln \left(\frac{\Omega_+}{\mu^2} \right) + \Omega_- \ln \left(\frac{\Omega_-}{\mu^2} \right) \right], \tag{18}$$

where dimensionless constants were removed by means of a redefinition of the parameter μ^2 .

Although the first term is divergent, we notice from (16) that $(\Omega_+ + \Omega_-) = \frac{1+a(\bar{\Psi}+\Psi)}{a^2}$. Due to the fact that the contribution to the effective action is obtained by the integration of the Kählerian potential over $d^8 z$, we find that the divergent term is annihilated by the Grassmann integration. Therefore, the final result is finite and equal to

$$\begin{aligned} K^{(1)}(\bar{\Phi}, \Phi) = & -\frac{1}{32\pi^2} \left[\frac{1 + a(\bar{W}'' + W'') + \sqrt{[1 + a(\bar{W}'' + W'')]^2 - 4a^2 \bar{W}'' W''}}{2a^2} \right. \\ & \times \ln \left(\frac{1 + a(\bar{W}'' + W'') + \sqrt{[1 + a(\bar{W}'' + W'')]^2 - 4a^2 \bar{W}'' W''}}{2\mu^2 a^2} \right) \\ & + \frac{1 + a(\bar{W}'' + W'') - \sqrt{[1 + a(\bar{W}'' + W'')]^2 - 4a^2 \bar{W}'' W''}}{2a^2} \\ & \left. \times \ln \left(\frac{1 + a(\bar{W}'' + W'') - \sqrt{[1 + a(\bar{W}'' + W'')]^2 - 4a^2 \bar{W}'' W''}}{2\mu^2 a^2} \right) \right]. \tag{19} \end{aligned}$$

It should be noticed that, in spite of the explicit dependence of this expression on μ , the corresponding low energy effective action (but not the effective Lagrangian) is scale invariant.

3. Nonminimal general Lagrangian

Now, let us consider the more generic case, that is, the arbitrary general Lagrangian $K(\bar{\Phi}, \Phi)$ (we will refer to this case as to the nonminimal one, as it was called in [12]):

$$S = \int d^8 z K(\bar{\Phi}, \Phi) + \left[\int d^6 z \left(-\frac{a}{2} \Phi \square \Phi + W(\Phi) \right) + \text{h.c.} \right]. \tag{20}$$

Repeating the calculations from the previous section, we find that the effective potential for this model is given by:

$$K^{(1)}(\bar{\Phi}, \Phi) = - \int \frac{d^4 p}{(2\pi)^4} \frac{1}{p^2} \ln(K_{\bar{\Phi}\Phi}) - \frac{1}{2} \int \frac{d^4 p}{(2\pi)^4} \frac{1}{p^2} \ln \left[1 + \frac{(ap^2 + \bar{W}'')(ap^2 + W'')}{K_{\bar{\Phi}\Phi}^2 p^2} \right]. \quad (21)$$

The first term in this expression vanishes within the dimensional reduction scheme. The second one can be simplified by use of the identity

$$1 + \frac{(\tilde{a}p^2 + \tilde{W}'')(\tilde{a}p^2 + \tilde{W}'')}{p^2} = \frac{(p^2 + \Omega_+)(p^2 + \Omega_-)}{p^2} \tilde{a}^2, \quad (22)$$

where we denoted $\tilde{Q} \equiv \frac{Q}{K_{\bar{\Phi}\Phi}}$ (with $Q = W''$ or $Q = \bar{W}''$ or $Q = a$), and

$$\Omega_{\pm} = \frac{1 + \tilde{a}(\tilde{W}'' + \bar{W}'') \pm \sqrt{[1 + \tilde{a}(\tilde{W}'' + \bar{W}'')]^2 - 4\tilde{a}^2 \tilde{W}'' \bar{W}''}}{2\tilde{a}^2}. \quad (23)$$

Thus, we can rewrite (21) as

$$K^{(1)}(\bar{\Phi}, \Phi) = -\frac{1}{2} \int \frac{d^4 p}{(2\pi)^4} \frac{1}{p^2} [\ln(p^2 + \Omega_+) + \ln(p^2 + \Omega_-) - \ln(p^2)]. \quad (24)$$

As before, the last term in this expression can be disregarded as it is field independent. The other integrals can be calculated by use of the dimensional reduction which again consists in the change $d^4 p \rightarrow \mu^{4-2\omega} d^{2\omega} p$. So, we have for $\omega \rightarrow 2$

$$K^{(1)}(\bar{\Phi}, \Phi) = \frac{1}{32\pi^2(2-\omega)} (\Omega_+ + \Omega_-) - \frac{1}{32\pi^2} \left[\Omega_+ \ln\left(\frac{\Omega_+}{\mu^2}\right) + \Omega_- \ln\left(\frac{\Omega_-}{\mu^2}\right) \right], \quad (25)$$

where some constants are reabsorbed into redefinition of μ^2 .

Finally, under appropriate substitutions, we find

$$\begin{aligned} K^{(1)}(\bar{\Phi}, \Phi) &= \frac{K_{\bar{\Phi}\Phi}^2}{32\pi^2 a^2 (2-\omega)} + \frac{(\bar{W}'' + W'')}{32\pi^2 a (2-\omega)} \\ &- \frac{1}{32\pi^2} \left[\frac{K_{\bar{\Phi}\Phi}^2 + a(\bar{W}'' + W'') + \sqrt{[K_{\bar{\Phi}\Phi}^2 + a(\bar{W}'' + W'')]^2 - 4a^2 \bar{W}'' W''}}{2a^2} \right. \\ &\times \ln \left(\frac{K_{\bar{\Phi}\Phi}^2 + a(\bar{W}'' + W'') + \sqrt{[K_{\bar{\Phi}\Phi}^2 + a(\bar{W}'' + W'')]^2 - 4a^2 \bar{W}'' W''}}{2\mu^2 a^2} \right) \\ &+ \frac{K_{\bar{\Phi}\Phi}^2 + a(\bar{W}'' + W'') - \sqrt{[K_{\bar{\Phi}\Phi}^2 + a(\bar{W}'' + W'')]^2 - 4a^2 \bar{W}'' W''}}{2a^2} \\ &\left. \times \ln \left(\frac{K_{\bar{\Phi}\Phi}^2 + a(\bar{W}'' + W'') - \sqrt{[K_{\bar{\Phi}\Phi}^2 + a(\bar{W}'' + W'')]^2 - 4a^2 \bar{W}'' W''}}{2\mu^2 a^2} \right) \right]. \quad (26) \end{aligned}$$

This result is explicitly divergent. However, we note that again, to obtain the contribution to the effective action, one must integrate the Kählerian effective potential over $d^8 z$. Thus, the second divergent term is annihilated via integration over Grassmannian coordinates, and the first one in the minimal case $K(\bar{\Phi}, \Phi) = \bar{\Phi}\Phi$ reduces to a field-independent constant whose integral over the superspace vanishes. At the same time, there is no other cases when the divergences vanish. Therefore, only the minimal models with $K(\bar{\Phi}, \Phi) = \bar{\Phi}\Phi$ yield the finite one-loop Kählerian effective potential.

It is interesting to discuss the behavior of this theory in the limit $a \rightarrow 0$, where the higher-derivative term in the action disappears. Taking the limit $a \rightarrow 0$ in (26), we find

$$\begin{aligned} K^{(1)}(\bar{\Phi}, \Phi) &= -\frac{1}{32\pi^2} \left\{ \frac{1}{a^2} \left[-\frac{K_{\bar{\Phi}\Phi}^2}{(2-\omega)} + K_{\bar{\Phi}\Phi}^2 \ln\left(\frac{K_{\bar{\Phi}\Phi}^2}{\mu^2 a^2}\right) \right] + \frac{1}{a} \left[W'' \ln\left(\frac{eK_{\bar{\Phi}\Phi}^2}{\mu^2 a^2}\right) + \bar{W}'' \ln\left(\frac{eK_{\bar{\Phi}\Phi}^2}{\mu^2 a^2}\right) \right] \right. \\ &\left. + \frac{1}{2K_{\bar{\Phi}\Phi}^2} (\bar{W}''^2 + W''^2) - \frac{\bar{W}'' W''}{K_{\bar{\Phi}\Phi}^2} \ln\left(\frac{K_{\bar{\Phi}\Phi}^2}{\mu^2 a^2}\right) + \frac{\bar{W}'' W''}{K_{\bar{\Phi}\Phi}^2} \ln\left(\frac{\bar{W}'' W''}{\mu^2 K_{\bar{\Phi}\Phi}^2}\right) + \mathcal{O}(a) \right\}. \quad (27) \end{aligned}$$

For the usual case, when $K(\bar{\Phi}, \Phi) = \bar{\Phi}\Phi$ and $W(\Phi) = \frac{m}{2}\Phi^2 + \frac{\lambda}{3!}\Phi^3$, the only terms surviving in the limit $a \rightarrow 0$ (when the higher derivatives are “switched off”) are

$$K^{(1)}(\bar{\Phi}, \Phi) = -\frac{1}{32\pi^2} \bar{\Psi} \Psi \ln(a^2 \bar{\Psi} \Psi), \quad (28)$$

where $\Psi \equiv m + \lambda\Phi$, that is, the one-loop effective potential for the Wess–Zumino model [11]. We conclude that the Wess–Zumino model can be treated as a limit of our higher-derivative theory.

4. Summary

We calculated the one-loop effective potential for nonconventional higher-derivative chiral superfield models with minimal and non-minimal kinetic terms. We find that in the minimal case it displays a nonpolynomial, logarithmic-like behavior while its finiteness is achieved in a very nontrivial manner – the divergent part vanishes due to properties of the integral over the Grassmann variables (in the nonminimal case the finiteness does not occur, actually the theory becomes nonrenormalizable). This finiteness of a formally divergent contribution is due to its chirality. The surviving of the nonpolynomial, logarithmic-like contributions, has never been achieved in superfield theories formulated in terms of $\mathcal{N} = 1$ superfields. It is easy to check that this contribution to the effective action is scale invariant.

Earlier, the arising of the logarithm-like finite contributions to the effective action (with the absence of divergences) has been shown to take place only in theories formulated in $\mathcal{N} = 2$ superspace, more precisely, in $\mathcal{N} = 4$ super-Yang–Mills theories [15]. So, we showed a new mechanism for achieving the finiteness for the $\mathcal{N} = 1$ superfield theories. We also saw that in a limit where the higher-derivative terms vanish in the classical Lagrangian, the one-loop effective potential for the Wess–Zumino model is reproduced, which confirms the consistency of the results. However, our theory is not restricted to be a simple regularization of the Wess–Zumino model, since we considered rather generic forms of chiral and general classical potentials. Furthermore, the significance of higher derivatives is not exhausted by their applications as a regularization since these theories can be treated on their own base and have been discussed within different situations [16]. But the presence of higher derivatives signalizes the possible occurrence of ghosts which jeopardize the physical interpretation of the theory. In this situation, modification on the properties of the asymptotic states, as in the Lee–Wick procedure [17], will become necessary or else our proposal should be considered as a low energy field theory.

For example, higher-derivative terms naturally emerge within string and supergravity contexts [18]. Therefore, it is natural to suppose that the presence of higher derivatives can result as a consequence of coupling with some extra fields. Indeed, it takes place in the dilaton supergravity when the integration over matter field is performed [4] and in theories where the light (massless) fields are coupled to heavy ones whose mass is compatible with the characteristic string mass, where the integration over these heavy fields gives a nontrivial modification in the effective action [13]. Therefore, we expect that the presence of higher derivatives can be treated as an efficient tool for the study of the low-energy effective behavior of fundamental field theory models which may arise within the context of strings and higher spins [18,19]. We expect that the study carried out in this paper can be useful to develop an answer to a key problem of the higher-derivative field theories – whether the equivalence between higher-derivative theories and models with ghosts discussed in [6] at the tree level also holds in the quantum contributions to the effective action.

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