SUPERFIELD EFFECTIVE POTENTIAL

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1 Introduction

This paper is a brief review of recent work carried out by our group [1-3] on definition and calculation of effective potential in N=1 supersymmetric field theory. As is well known, many problems of quantum field theory lead to necessity of studying of effective action or, in certain cases, effective potential. It is quite natural to expect that in theories processing some invariance group there should exist manifestly covariant methods for determing these objects in an explicitly symmetric form. The most adequate formulation for the N=1 supersymmetric field theories is based on the use of the N=1 superspace and corresponding superfields (see for example [4-6]). Any physical quantity of interest arising in such theories may be described in superfield terms. As to the effective potential in supersymmetric field theories, a problem of its calculation was studied by many authors [7-14] but practically all results have been obtained at the component level. The final form of superfield effective potential was unknown till nowadays. An essence of problem from our point of view looks like as follows. In conventional field theory the effective potential is an effective lagrangian evaluated at constant values of scalar fields, so its calculated can be implemented with the help of standard methods [15]. If we try to define the superfield effective potential as the superfield effective action at constant chiral and antichiral scalar superfields we will obtain that the effective action vanishes because of the known properties of the Berezin integral. Thus we must consider the superfield effective action for the scalar superfields preserving the arbitrary dependence on Grassmann coordinates $\theta$ and $\bar{\theta}$. It means, for calculation of superfield effective potential it is necessary to use a propagator of the theory in external superfields nontrivially depending on $\theta$ and $\bar{\theta}$. It is not clear from the very beginning that such a propagator can be found in explicit form. As shown in our papers [1, 2, 3] to calculate the superfield effective potential we should evaluate the superfield effective action at special (supersymmetric) conditions $\partial_{\theta} \Phi = \partial_{\bar{\theta}} \bar{\Phi} = 0$ where $\Phi(\bar{\Phi})$ is chiral (antichiral) scalar superfield. Dependence of the $\Phi$ and $\bar{\Phi}$ on $\theta$ and $\bar{\theta}$ is arbitrary. Thus we have a theory in the external superfields $\Phi$ and $\bar{\Phi}$ of above type. In thus case the standard super graph technique which is
very convenient for counterterm computations turns out to be badly adapted for treating the superfield effective potential. The more efficient approach should be developed. We have shown that the problem under consideration can be solved very efficiently in framework of superfield proper time technique. As a result we have found that the superfield effective potential is determined by three effects: kaehlerian effective potential, chiral effective potential and auxiliary fields effective potential. The main role for calculation of all these objects plays the superfield analog of heat kernel which can be solved explicitly for the arbitrary $\theta$, $\bar{\theta}$ - dependent background superfields. The question about chiral effective potential is worthy of special discussion. According to nonrenormalization theorem \cite{1,2,16}, all loop corrections to the effective action are expressed in terms of integrals over whole superspace but not over its chiral subspace. It was commonly believed as a result that the super symmetric effective potential does not get chiral-line quantum corrections, but it is not time. It was shown by West \cite{17} that in massless supersymmetric theories can arise contributions to the effective action of the form

$$\int d^4x d^2 \theta g(\Phi) (\frac{D^2}{4\Box}) f(\Phi)$$

which can be written due to the chirality of $\Phi$ in the form

$$\int d^4x d^2 \theta g(\Phi) f(\Phi)$$

and presents a contribution to chiral potential. This observation was further confirmed in ref \cite{3,8,13,20}.
2 The structure of effective action for Wess-Zumino model.

In this section we are going to discuss the general properties of superfield effective action of the Wess-Zumino (WZ) model. Its action has the form

$$S[\Phi, \bar{\Phi}] = \int d^8z \Phi \bar{\Phi} + \left( \int d^6z L_c(\Phi) + h.c. \right)$$

$$L_c(\Phi) = \frac{m}{2} \Phi^2 + \frac{\lambda}{3!} \Phi^3, \bar{D}_\alpha \Phi = 0$$

i.e this is theory of chiral superfield.

The effective action of WZ model $\Gamma[\Phi, \bar{\Phi}]$ has the common definition as the Legendre transform of the generating functional for connected Green function $W[J, \bar{J}]$:

$$\exp(i \bar{\hbar} W[J, \bar{J}]) = \int D\Phi D\bar{\Phi} \exp(i \bar{\hbar} S[\Phi, \bar{\Phi}] + J\Phi + \bar{J}\bar{\Phi})$$

$$\Gamma[\Phi, \bar{\Phi}] = W[J, \bar{J}] - J\Phi - \bar{J}\bar{\Phi}$$

where $\Phi$ is the vacuum expectation value of $\phi$ in theory with action $S[\Phi, \bar{\Phi}] + J\phi + \bar{J}\bar{\phi}$. After extracting of quantum fields $\chi, \bar{\chi}$ we will have

$$\exp(i \frac{\bar{\hbar}}{\hbar} \bar{\Gamma}[\Phi, \bar{\Phi}]) = \int D\chi D\bar{\chi} \exp[iS^{(\psi)}[\chi, \bar{\chi}]
+ i \frac{\hbar}{\sqrt{\hbar}} \frac{\lambda}{3!} \int d^6z \chi^3 - \chi \frac{\delta \bar{\Gamma}}{\delta \Phi} + h.c.)]$$

where $\bar{\Gamma}[\Phi, \bar{\Phi}] = \Gamma[\Phi, \bar{\Phi}] - S[\Phi, \bar{\Phi}]$ and

$$S^{(\psi)}[\chi, \bar{\chi}] = \int d^8z \chi \bar{\chi} - \left( \frac{1}{2} \int d^6z \psi \chi^2 + h.c. \right)$$

$$- \psi(z) = m + \lambda \Phi(z) = L_c''(\Phi), \bar{D}_\alpha \Phi = 0.$$ (7)

We will carry out our calculations in the framework of loop expansion

$$\Gamma[\Phi, \bar{\Phi}] = \sum_{n=1}^{\infty} h^n \Gamma^{(n)}[\Phi, \bar{\Phi}]$$

At the one-loop level, one gets

$$\Gamma^{(1)}[\Phi, \bar{\Phi}] = -\frac{i}{2} Tr \ln G^{(\psi)}$$

(9)
At the two-loop level one gets

$$\Gamma^{(2)}[\Phi, \bar{\Phi}] = -\frac{i}{2} \int D\chi D\bar{\chi} e^{iS[\chi, \bar{\chi}]} S_{\text{int}}^{2}[\chi, \bar{\chi}]$$ (10)

$G^{(\psi)}$ is a matrix superpropagator for theory (7),

$$G^{(\psi)}(z_1, z_2) = \left( \begin{array}{cc} G_{++}(z_1, z_2) & G_{+-}(z_1, z_2) \\ G_{-+}(z_1, z_2) & G_{--}(z_1, z_2) \end{array} \right)$$ (11)

it satisfies the equation

$$-\left( \frac{\psi}{4D^2} \right) \left( \begin{array}{cc} G_{++} & G_{+-} \\ G_{-+} & G_{--} \end{array} \right) = \left( \begin{array}{cc} \delta_+ & 0 \\ 0 & \delta_- \end{array} \right)$$ (12)

Calculation of $\Gamma^{(1)}$ is a very difficult problem due to very complicated structure of $G^{(\psi)}$. But using the results of [1-2] we can obtain the expression for $\Gamma^{(1)}[\Phi, \bar{\Phi}]$ in terms of the Green function $G^{(\psi)}_v$ satisfying the equation

$$\Delta G^{(\psi)}_v(z_1, z_2) = -\delta^8(z_1 - z_2)$$

$$\Delta = \Box + \frac{1}{4} \tilde{\psi} \tilde{D}^2 + \frac{1}{4} \tilde{D}^2 \psi$$ (13)

We can rewrite $\Gamma^{(1)}[\Phi, \bar{\Phi}] = -\frac{i}{2} tr \ln G^{(\psi)}_v$. The expression for $G^{(\psi)}_v$ has the form

$$G^{(\psi)}(z_1, z_2) = \frac{1}{16} \left[ \tilde{D}_1^2 \tilde{D}_2^2 G_{++}(z_1, z_2) \tilde{D}_1^2 \tilde{D}_2^2 G_{++}(z_1, z_2) \right]$$ (14)

For calculation of $G^{(\psi)}_v$ we can use proper-time representation

$$G^{(\psi)}_v = i \int_0^\infty ds U^{(\psi)}_v(z_1, z_2; s)$$ (15)

where the kernel $U^{(\psi)}_v$ is expressed as

$$U^{(\psi)}_v(z_1, z_2; s) = \exp(is\Delta)\delta^8(z_1 - z_2)$$ (16)

$$U_v(s) = U^{(0)}_v(s) = -\frac{i}{(4\pi s)^2} \delta_{12} \exp\left(\frac{i}{4} \frac{(x - x')^2}{s}\right)$$

Then can be rewritten in the form

$$\Gamma^{(1)}[\Phi, \bar{\Phi}] = -\frac{i}{2} \int_0^\infty ds U^{(\psi)}_v(s)$$ (17)

We also can read off the general structure of the effective action

$$\Gamma[\Phi, \bar{\Phi}] = \int d^8z L(\Phi, D_A \Phi, D_A D_B \Phi, \ldots \Phi, D_A \bar{\Phi}, D_A D_B \bar{\Phi}) + (\int d^8z L_c(\Phi) + h.c.)$$ (18)
L is effective super lagrangian

\[ L = K(\Phi, \bar{\Phi}) + F(D_\alpha \Phi, D^2 \Phi, \bar{D}_\dot{\alpha} \bar{\Phi}, \bar{D}^2 \bar{\Phi}, \Phi, \bar{\Phi}) + \ldots \]

\[ K = \Phi \bar{\Phi} + \sum_{n=1}^{\infty} \tilde{h}^n K^{(n)} \quad (19) \]

\[ F = \sum_{n=1}^{\infty} \tilde{h}^n F^{(n)} \]

\[ F|_{D_\alpha \Phi = \bar{D}_\dot{\alpha} = 0} \]

K is the kählerian effective potential (depending only on fields but not of their derivatives). F is the auxiliary fields potential (it is at least of third order in auxiliary fields of \( \Phi \) and \( \bar{\Phi} \)). \( L_c \) is the chiral effective potential (one-loop chiral contribution is equal to zero but higher corrections exist). The effective potential of WZ model is defined to be the effective lagrangian evaluated at the superfield conditions

\[ \partial_\alpha \Phi = \partial_\alpha \bar{\Phi} = 0 \quad (20) \]

To calculate K and F it is sufficient to evaluate the effective action for superfield satisfying eq.(20). As to \( L_c \), we can calculate \( \bar{\Psi} \)-independent part of \( G_c^{(\psi)} \) straightforward, and we have (18)

\[ G_c^{(\psi)}(z_1, z_2) = -\frac{1}{4} \delta^2(z_1 - z_2) + \frac{1}{4} \delta^2(\psi(z_1) \bar{D}^2 \delta^2(z_1 - z_2)) + O(\bar{\Psi}) \quad (21) \]

This ansatz may be used for loop calculations of \( L_c \).
3 One-loop approximation

To calculate one-loop effective potential we must calculate the heat kernel $U^{(\psi)}_v$ for special values of the background superfields:

$$\partial_a \Phi = \partial_{\bar{a}} \bar{\Phi} = 0 \quad (22)$$

(16) leads to

$$U^{(\psi)}_v(s) = \exp\left[i \int s(\bar{\psi}D^2 + \psi \bar{D}^2)\right]U_v(s) = \Omega(\psi|s)U_v(s) \quad (23)$$

To calculate $\Omega(\psi|s)$ we will use the Schwinger—De Witt method: really $\Omega(\psi|s)$ satisfies the equation

$$i \frac{\partial \Omega}{\partial s} + \frac{1}{4} \Omega(\psi \bar{D}^2 + \bar{\psi}D^2) = 0 \quad (24)$$

with initial condition

$$\Omega(0|s) = 1$$

We will represent $\Omega(\psi|s)$ in the form

$$\Omega(\psi|s) = 1 + \frac{1}{16} A(s)D^2 \bar{D}^2 + \frac{1}{16} \bar{A}(s)\bar{D}^2 D^2 + \frac{1}{8} B^a(s)D_a \bar{D}^2 + \frac{1}{8} \bar{B}_a(s)\bar{D}^a D^2 + \frac{1}{4} C(s)D^2 + \frac{1}{4} \bar{C}(s)\bar{D}^2 \quad (25)$$

and we will obtain equations for coefficients $A, \bar{A}, B, \bar{B}, C, \bar{C}$. Since we are interested to calculate the trace of $U^{(\psi)}_v$, we must to calculate essentially $A$ and $\bar{A}$, because

$$trU^{(\psi)}_v(s) = \int d^n z \int d^4 x' \delta^4(x - x') [A(s) + \bar{A}(s)]U(x, x'; s) \quad (26)$$

due to famous properties of trace and covariant derivatives. To solve the system of equations for coefficients $A, \bar{A}, B, \bar{B}, C, \bar{C}$ is a very complicated task, but for calculating kählerian potential (when $\psi$ and $\bar{\psi}$ are constant) the system is simplified and kernel has the form [1-2]

$$U^\psi_{v=const}(s) = (1 + \frac{1}{16\Box} \cosh(is\sqrt{\psi\bar{\psi}} - 1)] \times \{D^2, \bar{D}^2\} + \frac{1}{4\sqrt{\psi\bar{\psi}}} \sinh(is\sqrt{\psi\bar{\psi}})(\psi \bar{D}^2 + \bar{\psi}D^2)U_v(s) \quad (27)$$

To calculate regularized kählerian potential we must put term proportional to $\{D^2, \bar{D}^2\}$ and will obtain

$$K^{(1)}_{\text{reg}} = -i \int_L d^4 s \int \frac{d^4 x' \delta^4(x - x')}{s}[\cosh(is\sqrt{\psi\bar{\psi}} - 1)]U(x, x'; s) \quad (28)$$
We must use regularization by means of the proper-time cutoff because we have singularity at lower limit. Using famous properties of delta-function and kernel \( U(x, x'; s) \) we can obtain

\[
K^{(1)}_{\text{reg}} = -\frac{1}{32\pi^2} \bar{\psi}\psi(\ln \frac{\psi \bar{\psi}}{\mu^2} - \xi + \ln \mu^2 L^2) \quad (29)
\]

choosing the one-loop wave-function renormalization

\[
Z^{(1)} = \frac{\lambda^2}{32\pi^2} \ln \mu^2 L^2 \quad (30)
\]

cancels divergences, and we stay with one-loop effective kählerian potential

\[
K^{(1)}(\Phi, \bar{\Phi}) = -\frac{1}{32\pi^2} (m + \lambda \Phi)(m + \lambda \bar{\Phi}) \ln \left( \frac{(m + \lambda \Phi)(m + \lambda \bar{\Phi})}{\mu^2} \right) - \xi \quad (31)
\]

As to auxiliary fields' potential, it does not contain divergences, but its explicit calculation is very complicated due to highly cumbersome structure of the kernel \( U^{(v)}_v(s) \), and we can only to develop procedure for perturbative determining \( F^{(1)} \) (we have not yet succeeded in bringing result for auxiliary fields' potential in final form). For example, in lower (fourth) order we get

\[
F^{(1)} = \frac{\lambda^4 \zeta}{(4\pi)^2} (D^\alpha \Phi)(D_{\bar{\alpha}} \Phi)(\bar{D}^\bar{\alpha} \Phi)(\bar{D}_{\bar{\bar{\alpha}}} \Phi) |\mathcal{L}_\mu' |^{-4} + O^8(D, \bar{D}) \quad (32)
\]

where \( \zeta \) is a finite constant, defined in \[1\].
4 Two-loop approximation

Calculation of two-loop chiral and kählerian effective potential is based on properties of Green function \(G\), \(\bar{\psi}\), \(\Omega\) and presents no difficulties. For calculating two-loop chiral effective potential we must use \(\bar{\psi}\) and \(\Psi\) and will get \(\bar{\psi}\)-independent part of \(G^{(\psi)}\):

\[
G^{(\psi)}(z_1, z_2) = \frac{1}{16} \begin{bmatrix}
0 & -D_2^2 \cdot D_1^2 \cdot \delta^8(z_1 - z_2) & -D_2^2 \cdot D_1^2 \cdot \delta^8(z_1 - z_2) \\
-D_2^2 \cdot D_1^2 \cdot \delta^8(z_1 - z_2) & \delta^8(z_1 - z_2) & -D_2^2 \cdot D_1^2 \cdot \delta^8(z_1 - z_2) \\
-D_2^2 \cdot D_1^2 \cdot \delta^8(z_1 - z_2) & -D_2^2 \cdot D_1^2 \cdot \delta^8(z_1 - z_2) & \delta^8(z_1 - z_2)
\end{bmatrix}
\]  

(33)

It follows from \(\bar{\psi}, \psi, \Omega\) that the two-loop contribution to \(\Gamma[\Phi, \bar{\Phi}]\) depending only on \(\psi\) looks like

\[
\Gamma^{(2)} = \frac{\lambda^2}{12} \int d^6z_1 d^6z_2 [G^-(z_1, z_2)]^3
\]

\[
G^-(z_1, z_2) = \frac{D_2^2 \cdot D_1^2}{16 \Box_1} (|z_1|) \frac{1}{4 \Box_1} \delta^8(z_1 - z_2)
\]  

(34)

After quite obvious transformation we will get for massless case

\[
\Gamma^{(2)} = -\frac{\lambda^5}{12} \int \prod_{i=1}^5 \delta^8(z_i) \Phi(z_4) \Phi(z_5)
\]

\[
\delta^8(z_1 - z_3) \frac{D_2^2 \cdot D_1^2}{16 \Box_2} \delta^8(z_3 - z_2) \times \frac{1}{4 \Box_1} \delta^8(z_2 - z_4) \frac{D_2^2 \cdot D_1^2}{16 \Box_1} \delta^8(z_4 - z_1)
\]

\[
\frac{D_2^2 \cdot D_1^2}{16 \Box_1} \delta^8(z_1 - z_5) \frac{D_2^2 \cdot D_1^2}{16 \Box_2} \delta^8(z_2 - z_5)
\]  

(35)

After D-algebra transformations and converting to momentum representation we arrive at

\[
\Gamma^{(2)} = -\frac{\lambda^5}{12} \int d^4x d^4\theta \int \frac{d^4k_1 d^4k_2 d^4p_1 d^4p_2}{(2\pi)^{16}} \int d^4y_1 d^4y_2 e^{i p_1 (x - y_1) + i p_2 (x - y_2)} \Phi(x, \theta) \times
\]

\[
[k^2 \rho_1^2 (-D_2^2 \Phi(y_1, \theta)) \Phi(y_2, \theta) + (1 \leftrightarrow 2) + \frac{1}{2} (k_1, k_2) D^a \Phi(y_1, \theta) D_0 \Phi(y_2, \theta)]
\]

\[
\Omega^{-1}(k, p)
\]  

(36)

where

\[
\Omega(k, p) = k_1^2 k_2^2 (k_1 + k_2)^2 (k_1 - p_1)^2 (k_2 - p_2)^2 (k_1 + k_2 - p_1 - p_2)^2
\]  

(37)

Then use the rule \(\int d^4\theta = \int d^2\theta (-\frac{i}{2} D^2)\) we will get

\[
\Gamma^{(2)} = \frac{\lambda^5}{12} \int d^4x d^2\theta d^4y_1 d^4y_2 e^{i p_1 (x - y_1) + i p_2 (x - y_2)} \Phi(x, \theta) \Phi((y_1, \theta) \Phi(y_2, \theta) \times
\]

\[
J(p_1, p_2)
\]  

(38)
where
\[ J(p_1, p_2) = \int \frac{d^4k_1 d^4k_2 k_1^2 p_1^2 + k_2^2 p_2^2 - 2k_1 k_2 p_1 p_2}{\Omega(k, p)} \]  
(39)

To calculate effective potential we must put Φ to be slowly varying in spacetime, it implies
\[ \Phi(x, \theta) \Phi(y_1, \theta) \Phi(y_2, \theta) \simeq \Phi^3(x, \theta) \]  
(40)

Then (40) takes the form
\[ \Gamma^{(2)} = \lambda_5 \frac{12}{12} J(p_1 \to 0, p_2 \to 0) \int d^6z \Phi^3(z) \]  
(41)

It is famous that
\[ J(p, 0) = \frac{6}{(4\pi)^4}\zeta(3) \]  
(42)

and the two-loop chiral effective superpotential reads as [3]
\[ L_c(\Phi) = L_{c} + \bar{h}^2 L^{(2)}_{c} = \left( \frac{\lambda}{3!} + \frac{\lambda^5 \zeta(3)}{2(4\pi)^4} \right) \Phi^3 \]  
(43)

As for two-loop kählerian potential, we can calculate \( G(\psi) \) for constant values and obtain
\[ G^{(\psi)} = -\frac{1}{\Box - M^2} \left( \begin{array}{cc} \bar{\psi} & -\frac{i}{2} \bar{D}^2 \\ \frac{i}{2} D^2 & \psi \end{array} \right) \left( \begin{array}{cc} \delta_+ & 0 \\ 0 & \delta_- \end{array} \right) + \ldots \]  
(44)

where \( M^2 = \psi \bar{\psi} \), dots means all terms which give no contribution to \( K^{(2)} \). From [3] we can conclude that \( K^{(2)} \) is expressed by Green function as
\[ K^{(2)} = \frac{\lambda^2}{6} \int d^6z_1 d^6z_2 (G_+(z_1, z_2))^3 \]  
(45)

In momentum representation it leads to
\[ K^{(2)} = -\frac{\lambda^2}{6} \int d^4xd^4\theta_1 d^4\theta_2 \frac{d^4k d^4l}{(2\pi)^8} \frac{1}{(k^2 + M^2)(l^2 + M^2)(k^2 + l^2 + M^2)} \]  
\[ \frac{D^2_1}{4} \delta_{12} \frac{D^2_2}{16} \delta_{12} \frac{D^2_3}{4} \delta_{12} \]  
(46)

After D-algebra transformations, integration by momenta and regularization we have
\[ K_{REG}^{(2)} = \frac{\lambda^2}{6(4\pi)^4} \int d^4\theta \left[ \frac{6M^2}{\epsilon^2} + \frac{3M^2}{\epsilon} (3 - 2\gamma - 2 \ln \frac{M^2}{\mu^2}) + 3M^2 \ln^2 \frac{M^2}{\mu^2} - 3M^2 (3 - 2\gamma) \ln \frac{M^2}{\mu^2} + 9M^2 (1 - \gamma) \right] \]  
(47)
After subtraction of one-loop and two-loop counterterms we arrive to

\[ K_{\text{renorm}}^{(2)} = -\frac{\lambda^2}{(4\pi)^4} \int d^4\theta (m + \lambda \Phi)(m + \lambda \bar{\Phi}) \left[ -\frac{1}{4} \ln^2 \frac{(m + \lambda \Phi)(m + \lambda \bar{\Phi})}{\mu^2} - \frac{3}{2} \ln \frac{(m + \lambda \Phi)(m + \lambda \bar{\Phi})}{\mu^2} + \frac{3}{2}(\gamma - 1) + \frac{1}{4}(\gamma^2 + \zeta(2)) - b \right] \]

where \( b \) is finite part of two-loop counterterm. It will be determined by renormalization condition (for massive case)

\[ \frac{\partial^2}{\partial \Phi \partial \bar{\Phi}} \left( K^{(0)} + K^{(1)} + K^{(2)} \right)_{\Phi=\bar{\Phi}=0} = 1 \] (49)

where \( K^{(0)} = \Phi \bar{\Phi} \) and \( K^{(1)} \) is given in (29). After calculations we get

\[ K^{(2)} = -\frac{\lambda^2}{(4\pi)^4} \int d^4\theta (m + \lambda \Phi)(m + \lambda \bar{\Phi}) \left[ -\frac{1}{4} \ln^2 \frac{(m + \lambda \Phi)(m + \lambda \bar{\Phi})}{m^2} - \frac{5}{2} \ln \frac{(m + \lambda \Phi)(m + \lambda \bar{\Phi})}{m^2} + \frac{9}{2} \right] \]

and using this choice for \( \mu \) we have

\[ K^{(1)} = \frac{1}{2(4\pi)^2} \int d^4\theta (m + \lambda \Phi)(m + \lambda \bar{\Phi}) \left[ \ln \frac{(m + \lambda \Phi)(m + \lambda \bar{\Phi})}{m^2} - 2 \right] \] (51)
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