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SUSY EFFECTIVE POTENTIAL USING SUPERFIELDS*

by

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ABSTRACT

We illustrate the use of superfields for computing the effective potential in supersymmetric theories by applying the superfield tadpole and vacuum bubble methods to the general renormalizable action of interacting chiral superfields. We work with the unconstrained chiral superfield potentials over which the functional integral for the chiral superfields may be easily formulated. The additional abelian gauge invariance thus introduced is taken care of by adding to the action a ghost-free gauge-fixing term. The superpropagators in a classical background are derived in a conveniently compact form to be useful for calculations.

Key-words: Supersymmetry; Effective potential; Superfields.

1 INTRODUCTION

For the globally supersymmetric theories[1] the superfield[2] formulation is very economical to calculate quantum corrections. The cancellation of the divergences associated with the bosonic loops with those of the fermionic loops is automatically taken care of through the superpropagators. The number of supergraphs required to be calculated is greatly reduced compared to that needed in the component formulation. The non-renormaliztion theorems may be shown directly. The superfield path integral formulation is a powerful calculational tool, for example, through the background field method[3] , the possibility of performing a change of variables with the corresponding evaluation of the Jacobi deteminant leading to Ward identities apart from comactness and other advantages. The superfield formulation has now been developed sufficiently[3),4)] so as to allow manageable calculations to higher order loops.

We first derive for the case of several ineracting chiral superfields the superpropagators of the unconstrained superpotentials in the presence of a classical background[5]. After a short comment on the background field method we discuss in detail the superfield tadpole[5),6)] and the vacuum bubble methods for computing the effective potential to two-loops.

2. THE GENERATING FUNCTIONAL FOR CHIRAL SUPERFIELD

2.1 Unconstrained Chiral Superfield Potential

The chiral superfield satisfies differential constraint $\bar{D}\Phi = 0$ and it seems difficult to formulate the functional integral over , which could take care of these constraints. We may, however, analogous to the case of the e.m. field introduce the unconstrained superfield potentials[7),5)] S, $\bar{S}$ and write

$$\Phi = -(1/4)\bar{D}^2 S \qquad , \qquad \bar{\Phi} = -(1/4)D^2 \bar{S} \tag{1}$$

It introduces in the theory an additional invariance under the abelian gauge transformations: $S \longrightarrow S + \bar{D}\bar{F}$, $\bar{S} \longrightarrow \bar{S} + DF$. We may take care of it by adding to the action the following ghost-free gauge-fixing term[8)]

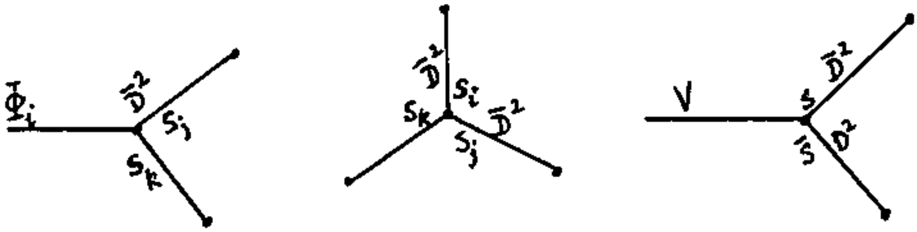
$$I_{G.F.} = \alpha^{-1}\int d^8 z \;\; \bar{S}_i (1 - P_1)\Box S_i \tag{2}$$

The functional integral may then be formulated[8)] over S, $\bar{S}$ and the perturbation theory done with the S, $\bar{S}$ propagators rather than Φ , $\bar{\Phi}$

ones results in the integrals over the full superspace. This is also clear from, say, if we write the cubic interaction term as the full superspace integral

$$I_{int.}=(1/3)g_{ijk}\int d^8z\ S_i(-1/4)\bar{D}^2S_j\ (-1/4)\bar{D}^2S_k\ + C.C. \tag{3}$$

and consider, for example, the following self-explanatory diagrams which also indicate the Feynman rules for the vertices following from the Wick's theorem



(a) (b) (c)

Fig.1: Feynman rules for the vertices (a) two internal S lines, (b) three internal S lines, (c) $\bar{\Phi}V\Phi$ vertex where V is the gauge superfield.

2.2 S, $\bar{S}$ Superpropagators For The Shifted Theory[5),8)]

It is well known that in order tocompute the effective action we are required to know the propagators of the 'shifted' theory obtained in our context by substituiting each chiral superfield Φ_i (and V when the gauge superfield is present) by $\Phi_i + C_i$ where C_i are background superfields satisfying $\bar{D}C_i = D\bar{C}_i = 0$. For the case of several interacting chiral superfields we obtain the following terms in the resulting action

$$I_b=\int d^8z\ \ \bar{C}_iC_i\ \ +[\int d^6s\ W(C) + C.C.] \tag{4}$$

$$I_{int}=\int d^6s\ [\ \lambda_i\Phi_i + m_{ij}C_i\Phi_j + g_{ijk}C_iC_j\Phi_k + \int d^2\bar{\theta}\ \bar{C}_i\Phi_i\ ]$$

$$+(1/3)g_{1jk}\int d^8z\ \Phi_1S_j(-1/4)\ \bar{D}^2S_k \tag{5}$$

$$I_0=\int d^8z\ [\ \bar{S}_i\ P_1\Box\ S_i\ -(1/8)\ C_{ij}S_i\bar{D}^2S_j-(1/8)\ \bar{C}_{ij}\bar{S}_iD^2\bar{S}_j\ ]$$

$$+\ I_{G.F.} + \int d^8z\ (J_iS_i + \bar{J}_i\bar{S}_i) \tag{6}$$

Here $W(\Phi) = \lambda_i \Phi_i + (1/2) m_{ij}\, \Phi_i \Phi_j + (1/3)\, g_{ijk}\, \Phi_i \Phi_j \Phi_k$ indicates the superpotential, $(C_{ij}) = (m_{ij} + 2\, g_{ijk}\, C_k)$ and to the free action I_o we have added a source term and the gauge-fixing term given in (2). The equations of motion following from (6) are

$$A \begin{pmatrix} S \\ \bar{S} \end{pmatrix} = \begin{pmatrix} -(1/4) C_{ij} \bar{D}^2 & \Box [P_2 + \alpha^{-1}(P_1+P_T)] \\ \Box [P_1 + \alpha^{-1}(P_2+P_T)] & -(1/4)\bar{C}_{ij} D^2 \end{pmatrix} \begin{pmatrix} S \\ \bar{S} \end{pmatrix} = - \begin{pmatrix} J \\ \bar{J} \end{pmatrix} \tag{7}$$

The free superpropagators of the shifted theory may be defined by

$$A \begin{pmatrix} \Delta^{SS} & \Delta^{S\bar{S}} \\ \Delta^{\bar{S}S} & \Delta^{\bar{S}\bar{S}} \end{pmatrix} = i\, \delta^8(z-z') \tag{8}$$

We find

$$\Box [P_1 + \alpha^{-1}(P_2+P_T)] \Delta^{S\bar{S}} - (1/4)\ \bar{C}\ D^2\, \Delta^{\bar{S}\bar{S}} = i\, \delta^8(z-z') \tag{9a}$$

$$-(1/4) C\ \bar{D}^2 \Delta^{S\bar{S}} + \Box\ [P_2 + \alpha^{-1}(P_1+P_T\ )\ ]\ \Delta^{\bar{S}\bar{S}} = 0 \tag{9b}$$

From the properties of the projection operators P_1, P_2 and P_T it is clear that (9b) is satisfied if

$$\Delta^{\bar{S}\bar{S}} = (1/4\,\Box)\ C\ \bar{D}^2\, \Delta^{S\bar{S}} \tag{10}$$

It follows then from (9a) that

$$[\ P_1(\Box - \bar{C}P_1 C) + \alpha^{-1}(P_2+P_T)\Box]\ \Delta^{S\bar{S}} = i\ \delta^8(z-z') \tag{11}$$

We may easily invert it if we note that $\bar{C}P_1 C = P_1 \bar{C} P_1 C = \bar{C} P_1 C P_1$ and find

$$\Delta^{S\bar{S}} = i\ [\ \alpha(P_2+P_T)/\Box + P_1(\Box - \bar{C}P_1 C)^{-1}]\ \delta^8(z-z') \tag{12}$$

The other propagators or Green's functions have analogous forms.

In the background field method[3)] the effective action is obtained by evaluating in the perturbation theory 1PI graphs with only internal S, $\bar{S}$ lines and external C_i legs. The terms which are linear in Φ or $\bar{\Phi}$ are dropped as they do not contribuite to 1PI graphs with no external Φ, $\bar{\Phi}$ legs. The term in (4) gives the background action; it is the classical contribuition to the effective action. The background vertices are read from (6) and the S, $\bar{S}$ propagators obtained from (10), (12) by replacing C_{ij} by m_{ij} so that $P_1(\Box - \bar{C}P_1 C)^{-1} \longrightarrow P_1(\Box - \bar{m}\, m)^{-1}$. We note that the background superfields C_i are not required to be expressed in terms of the unconstrained superfields in order to write the background vertex term in (6) as an integral over the full superspace. The non-renormalization theorem is then apparent. The Feynman rules derived above give rise to the contribuitions that involve only full superspace integrals and the couplings to the constrained background superfields C_i and $\bar{C}_i$. If we admit only local operators it is not possible to write counterterms depending only on C_i as full superspace integrals and consequently they are ruled out; the superpotential is not renormalized.

3 SUPERFIELD METHODS FOR COMPUTING EFFECTIVE POTENTIAL

3.1 Superfield Tadpole Method

The well known methods of calculating effective potential in ordinary field theory may be extended to supersymmetric theories formulated in terms of superfields. The number of supergraphs required to be evaluated is greatly reduced compared to that needed in the corresponding component formulation.

In the superfield tadpole method the background field is a constant chiral superfield with vanishing spinor components, $C_i = a_i + f_i\theta^2$. We remark that the auxiliary fields must also be shifted along with the physical fields to obtain for the initially supersymmetric theory a non-vanishing tadpole contribuition from the corresponding 'shifted' theory. A compact form for the superpropagators expliciting the poles and the θ, $\bar{\theta}$ dependence is easily derived[5)] for the case under consideration:

$$P_1(\Box - \bar{C}P_1 C)^{-1}\, \delta^8(z-z')$$
$$= P_1(\Box - \bar{M}M)^{-1}\delta^8(z-z') +$$
$$e^{-i\theta\sigma\cdot\partial\bar{\theta} - i\theta'\sigma\cdot\partial\bar{\theta}'} \left[\, A\bar{\theta}^2\theta'^2 + B + C\,\bar{\theta}^2 + E\,\theta'^2\,\right] \delta^4(x-x') \qquad (13)$$

where $(\partial_\ell \longrightarrow ik_\ell)$

$$
\begin{aligned}
A &= (k^2+\bar{M}M)^{-1} - [\,k^2 + \bar{M}M - \bar{f}\,(k^2 + M\bar{M})^{-1} f\,]^{-1} \,,\\
k^2 B &= \bar{M}\,(k^2 + M\bar{M})^{-1} f\, C \,,\\
k^2 C &= [\bar{f}\,(k^2+M\bar{M})^{-1} f - k^2 - \bar{M}M\,]^{-1}\,\bar{f}M(k^2+\bar{M}M)^{-1}\\
k^2 E &= (k^2+\bar{M}M)^{-1}\,\bar{M}f\,[\bar{f}(k^2+M\bar{M})^{-1} f - k^2 - \bar{M}M]^{-1}
\end{aligned} \tag{14}
$$

with $M=(M_{ij}) = m_{ij} + 2g_{ijk}a_k$ and $f=(f_{ij}) = 2g_{ijk}f_k$.

The one-point function to zero loop approximation is readily calculated from (5) and we find

$$
i\,\Gamma_o^{(1)} = i\int d^2\theta\,[\lambda_i + m_{ij}C_j(\theta) + g_{ijk}C_jC_k + \int d^2\bar{\theta}\,\bar{C}_i(\bar{\theta})]\,\tilde{\Phi}_i(0,\theta,\bar{\theta}) + \text{C.C.} \tag{15}
$$

where a tilde denotes the Fourier transform and

$$
\tilde{\Phi}_i(0,\theta,\bar{\theta}) = \tilde{A}_i(0) + \sqrt{2}\,\theta\,\tilde{\psi}_i(0) + \theta^2\,\tilde{F}_i(0) \tag{16}
$$

We find

$$
\Gamma_o^{(1)} = (\bar{f}_i + \frac{\partial W}{\partial a_i})\,\tilde{F}_i(0) + M_{ij}\,f_j\,\tilde{A}_i(0) + \text{C.C.} \tag{17}
$$

The tadpole contribuitions of the component fields may be read off[5] as the coefficients of $\tilde{F}_i(0)$, $\tilde{A}_i(0)$ etc. Hence we obtain[9),5)]

$$
-\frac{\partial V_o}{\partial f_i} = \bar{f}_i + \frac{\partial W}{\partial a_i} \quad , \quad -\frac{\partial V_o}{\partial a_i} = M_{ij}\,f_j \tag{18}
$$

plus their complex conjugates. On integrating (18) we find the tree level expression for the effective potential

$$
V_o = -(\bar{f}_i f_i + f_i\,\frac{\partial W}{\partial a_i} + \bar{f}_i\,\frac{\partial \bar{W}}{\partial \bar{a}_i}) \tag{19}
$$

where $a \longrightarrow A_c$ and $f \longrightarrow F_c$ is understood.

The computation of the one-loop effective potential requires the evaluation of a single superfield tadpole graph. From (5) we read its contribuition to be

$$
i\,\Gamma_1^{(1)} = i\,(1/3).3.\int d^4\theta\,\tilde{\Phi}_\ell(0,\theta,\bar{\theta})\,[\,\text{Tr}\,g_\ell\,(-1/4)\,D^2\,\Delta^{\bar{5}\bar{5}}(z,z')]_{z=z'} \tag{20a}
$$

where $(g_\ell)_{ij} = g_{\ell ij}$. Working out the $\theta,\bar{\theta}$ integrations in a straightforward fashionin view of (13) we find[5]

$$-\partial V_1/\partial \bar{f}_l = \mathrm{tr}\; g_l\, (k^2+M\bar{M})^{-1}\, f[k^2+\bar{M}M-\bar{f}(k^2+M\bar{M})^{-1}f]^{-1} \quad ,$$

$$-\partial V_1/\partial \bar{a}_l = \mathrm{tr}\; g_l\, M\{(k^2+\bar{M}M)^{-1}-[\,k^2+\bar{M}M-\bar{f}(k^2+M\bar{M})^{-1}f]^{-1}\}$$
$$+\; \mathrm{tr}Mg_l\,\{(k^2+M\bar{M})^{-1}-[k^2+M\bar{M}-\bar{f}(k^2+\bar{M}M)^{-1}f]^{-1}\} \qquad (20b)$$

where $\mathrm{tr}=-i\int d^4k\mathrm{Tr}/(2\pi)^4$. The expressions are valid even when the supersymmetry is spontaneously broken with evident modifications during the substituitions of a and f by A and F. A straightforward integration of these partial differential equations leads to[5),10)]

$$V_1 = (1/2)\;\mathrm{tr}\;\ln[\;I - (k^2+\bar{M}M)^{-1}\;\bar{f}\;(k^2+M\bar{M})^{-1}f] \qquad (21)$$

Analogous procedure may be followed to calculate the effective potential to two-loops. We must evaluate the following supergraphs

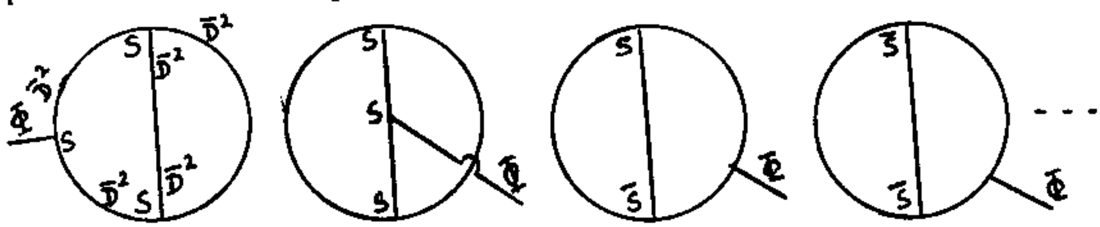


Fig.2: Two-loops superfield tadpole diagrams needed for effective potential in two-loops.

In practice we need to calculate only three diagrams; the others being obtained by inspection[11)].

3.2 Superfield Vacuum Bubble Method

The superfield tadpole method results in the partial derivatives of the effective potential and is very convenient to study the new minimum when the supersymmetry is spontaneously broken and imposing the renormalization constraints. On the contrary the vacuum bubble method gives directly the effective potential. We are simply required to calculate for the 'shifted' theory discussed in Sec.3.1 the vacuum-vacuum diagrams. For example, in the zero loop approximation we find from (4)

$$-i\;V_o = i\int d^4\theta\;\bar{C}_i C_i\;+[\int d^2\theta\;W(C)\;+\;\mathrm{C.C.}] \qquad (22)$$

On performing the integrations we find the tree level result. At the two-loops level we have simply to evaluate the three distinct diagrams shown in Fig.2 with the external lines removed. The calculation is much less tedious than using the component formulation and a computer program makes it quite manageable even in still higher loops except for the usual difficulties associated with renormalization of the effective potential.

4 RENORMALIZED EFFECTIVE POTENTIAL TO TWO-LOOPS FOR THE W-Z MODEL

We describe here briefly the procedure for obtaining the renormlized effective potential in the case of a single chiral superfield. From the comments made earlier on the background field method it is clear that we need to perform only the wave function renormalization. The action written in terms of renormalized superfields and parameters reads as

$$\int d^8z \quad Z\,\bar{\Phi}\Phi + [\int d^6s \left\{(1/2)m\Phi^2 + (1/3)g\Phi^3\right\} + \text{C.C.}] \tag{23}$$

where Z is the renormalization constant which is expressed as $Z=1+\hbar Z_1+\hbar^2 Z_2+..$ with $\hbar$ explicitly indicated to keep track of the order of the counterterms. We obtain at the zero-loop

$$V_o = V_o^{\,o} + V_o^{\,1} + V_o^{\,2} + \ldots \tag{24}$$

where $V_o^{\,o} = V_{oR}$ is the regular zero-loop contribuition to the effective potential while

$$V_o^{\,1} = -\hbar\, Z_1 \mid f \mid^2 \qquad , V_o^{\,2} = -\hbar^2\, Z_2 \mid f \mid^2 \tag{25}$$

and which act as counter terms in the cancellation of the divergences at the one and two-loops. Writing $\tilde{f} = 2gf$, $\tilde{a} = m + 2ga$ we obtain the one-loop correction

$$V_1 = -(\hbar/2)\int d\tilde{f}\,(f^*/Z^4)\int \frac{d^4k}{(2\pi)^4}\,[(k^2+|\tilde{a}|^2)^2 - |\tilde{f}|^2]^{-1} \tag{26}$$

where we make use of (20) and perform a change of variables $Z\tilde{f} \longrightarrow \tilde{f}$, $Zk \longrightarrow k$. We use the dimensional regularization and find

$$V_1 = V_1^{\,1} + V_1^{\,2} \qquad ,$$

$$V_1^{\,1} = (\hbar/64\pi^2)\ [\ -(-\frac{4}{\epsilon} + 3)|\tilde{f}|^2 + \ldots\ldots],$$

$$V_1^{\,2} = (\hbar^2 Z_1/16\pi^2)\ [\ 2(-\frac{1}{\epsilon} + 1)|\tilde{f}|^2 + \ldots\ldots] \tag{27}$$

The two-loop correction is calculated to be

$$V_2=(4g^2\hbar^2/(16\pi^2)^2)[\frac{1}{\epsilon^2}|f|^2-\frac{1}{\epsilon}\{(\frac{3+\alpha}{2})|f|^2-\tilde{a}_+^4\ln(\tilde{a}_+^2/\mu^2)-\tilde{a}_-^4\ln(\tilde{a}_-^2/\mu^2)$$
$$+(1/2)\tilde{a}_+^2|f|\ln(\tilde{a}_+^2/\mu^2)-(1/2)\tilde{a}_-^2|f|\ln(\tilde{a}_-^2/\mu^2)+2\tilde{a}^4\ln(\tilde{a}^2/\mu^2)\}+..] \quad (28)$$

where $\tilde{a}_{+(-)}^2=\tilde{a}^2{}_{(-)}^{+}|f|$ and only divergent terms are displayed.

An important point to note is that in the conventional theories the coupling constant is renormalized independent of the wave function and any renormalization prescription for the first on cannot influence the kinetic energy terms which are controlled by the second. In the supersymmetric theories it is clear that the kinetic term dictates the correct renormalization prescription so as to avoid negative kinetic terms[12)]. In the case under consideration we may determine Z_1 and Z_2 by imposing the following constraint

$$-\partial^2 V_{eff}/\partial f^*\partial f=1 \quad \text{for } f=f^*=0 \quad (29)$$

We find

$$Z_1=(4g^2/16\pi^2)\ [\ \frac{1}{\epsilon}+(1/2)\ln(\tilde{a}^2/\mu^2)\ ],$$
$$Z_2=-16(g^2/16\pi^2)^2\ [\frac{1}{\epsilon^2}+\frac{1}{\epsilon}\{\frac{\alpha-1}{2}+\ln(\tilde{a}^2/\mu^2)\}-(1/2)\ln(\tilde{a}^2/\mu^2)+(1+\gamma)/2] \quad (30)$$

The complete regularized expression for the effective potential hence obtained is given in the ref. 11 where its behavior with respect to the physical field A when the auxiliary field F is eliminated using its (corrected) equations of motion is also discussed.

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APPENDIX

POTENCIAL EFETIVO NAS TEORIAS SUPERSIMÉTRICAS USANDO SUPERCAMPOS*

Supersimetria[1] é uma simetria relativistica no espaço-tempo, diferente das simetrias internas, entre bosons e fermions. Esta é no momento atual a única maneira de obter uma unificação de simetrias do espaço-tempo e de simetrias internas da matriz S no contexto da teoria de partículas relativisticas. Os multipletos de supersimetria contém tanto as partículas(ou campos) bosonicas como as fermionicas. Tentou-se, antes da descoberta da supersimetria, formular de diversas maneiras simetrias 'superiores' cujos multipletos contenham todas as partículas ou com espins inteiros ou todas com espins semi-inteiros, por exemplo, a simetria SU(6). Demonstrou-se[2], entretanto, que estas simetrias 'superiores' não são relativisticas e a unificação acima mencionada é trivial. Os teoremas 'no go' se basearem sob a hipotese do que os geradores de simetrias são bosonicos e geram a bem conhecida algebra de Lie. Os geradores de supersimetria por outro lado são fermionicos e, portanto, obedecem relações de anti-comutação. Verifica-se que os geradores de grupo de Poincaré junto com os de supersimetria não geram uma algebra de Lie, mas, sim uma algebra[3] Lie-admissivel sob o seguinte comutador modificado ou gradado

$$\{M_1, M_2] = M_1 M_2 - (-1)^{m_1 m_2} M_2 M_1$$

onde 'm' indica o número total dos índices fermionicos no operador M. Devido a mudança em algebra os teoremas 'no go' acima não são mais válidos e a unificação relativistica agora torna-se não trivial[4].

A supersimetria torna sempre uma teoria de campo comum em uma teoria com alto grau de renormalizibilidade. As divergências quadraticas decorrentes dos loops bosonicos são milagrosamente canceladas pelas divergências correspondentes geradas em loops dos companheiros fermionicos. Elevando a supersimetria ao nível local a gravitção é introduzida automaticamente. A teoria de Supergravitação[5] é a única teoria no momento atual candidata para unificar as interações fortes, eletromagnéticas, fracas e gravitcionais bem como serve como ponto de ligação entre problemas não resolvidos de Cosmologia[6] e de física das partículas elementares

A introdução de noção de superespaço que contém além das coordenadas x tmabém as coordenadas anticomutativas θ_α , $\bar{\theta}_{\dot{\alpha}}$ e do supercampo[7] permitiu um rápido desenvolvimento da teoria de campos supersimetricos. É viável agora usando supercampos quantizados calcular ação efetiva, por exemplo, usando método de

*Apresentado no 5º Encontro Nacional de Física de Partículas Elementares e Campos (Itatiaia-1984).

'background field' em loops de ordens superiores. Contribuimos neste contexto recentemente com os seguinte trabalhos:

P.P. Srivastava, Superfield Tadpole Method For SUSY Effective Potential, Phys. Lett. 132B(1983)80

P.P. Srivastava, On The Functional Integral For Chiral Superfield, CERN preprint TH-3937/84 e Phys. Lett. 149B(1984)135

R. P. dos Santos e P.P. Srivastava, Two-Loop Effective Potential for Wess-Zumino and O' Raifeartaigh Model, CBPF preprint 1985.

Introduziu-se no primeiro trabalho supercampos $S,\bar{S}$ sem vínculos para supercampos quirais de matéria e demonstrou-se que para quantização estes e não os supercampos com vínculos diferenciais eram relevantes. Foram achados também os propagadores para a teoria 'deslocada' na qual a massa em vez de ser um constant torna-se um supercampo. Estes propagadores são necessários para computação de potencial efetivo usando supercampos. No trabalho seguinte formulou-se o integral funcional sobre os supercampos S, $\bar{S}$ achando-se inclusivo uma condição de fixação de 'gauge' sem fantasmas. Não é conhecido tal integral funcional em termos de supercampos quirais com vínculos onde sempre se interpreta o integral funcional em termos de campos em componentes. É possível agora trablhar diretamente com os supercampos também no caso de supercampo de matéria. Estudou-se também o teorema de não-renormalização usando supercampos. No último trablho estendemos o calculo de potencial efetivo feito no primeiro trablho até 1-loop para até dois -loops usando supercampos e métodos de 'tadpole' e de bolha de vácuo.

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