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OFF-SHELL CENTRAL CHARGES

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We show that there are only two kinds of (classical) central charges which can occur in the off-shell continuation of a supersymmetry algebra: mass dependent and interaction dependent. Both can be obtained by dimensional reduction, but only the mass-dependent kind allows self-interactions (with certain exceptions: see added note). As an example, we derive both kinds for a massive three-dimensional O(2) chiral multiplet coupled to an O(2) vector multiplet. We also discuss four-dimensional extended multiplets.

1. Continuation off-shell

Central charges Z in the on-shell supersymmetry algebra [1] are just masses m times dimensionless internal-symmetry generators G which commute with the supersymmetry generators Q_α. There are two ways to extend the relation $Z = mG$ off-shell (in a classical field theory):

(i) choose $Z = mG$ identically off-shell, so that the supersymmetry generators will have m-dependence off-shell also;

(ii) choose Z independent of m off-shell, so that $Z = mG$ is only an on-shell relation (i.e., a field equation), causing Z and Q_α to have explicit dependence on the coupling constants and fields off-shell.

On-shell central charges can be considered as momentum generators of extra spatial dimensions [2]. This suggests that superfield formalisms which include central charges [3] can most easily be derived by dimensional reduction of massless higher-dimensional superfields [4–7]. The most natural way to introduce masses in dimensional reduction is to choose the fields to have the dependence $\psi(y) = e^{-imGy}\psi(0)$ on one of the higher-dimensional coordinates y [8]. In order to avoid spontaneous breakdown of supersymmetry [8], we must have $[G, Q_\alpha] = 0$. Then $(i\partial/\partial y)\psi = Z\psi$, with $Z = mG$ a central charge, and Z will naturally occur in the supersymmetry algebra $\{Q, Q\} \sim P$. We thus obtain the m-dependent form of central charges described above. In this case there are no central charges in the off-shell supersymmetry algebra when $m = 0$. Also, since the Z-dependent ($m \neq 0$)

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theory as derived from the same higher-dimensional action as the Z-free ($m=0$) theory, the introduction of Z does not alter the auxiliary field structure.

Sohnius [3] has also used extra commuting coordinates to introduce central charges, and his notation can easily be translated into the language of higher dimensions. His method can then be seen to be equivalent to *imposing the higher-dimensional massless field equations as constraints*, and solving them to determine the y-dependence of the superfields. For example, if the explicit y-dependence of the superfield equation can be written in the form $(i\partial/\partial y - Z)\psi(y) = 0$, we have $\psi(y) = e^{-iZy}\psi(0)$. Instead of using the higher-dimensional action, we use an action which gives $(Z - mG)\psi(0) = 0$ as the field equation. Z is thus the off-shell central charge, and $Z = mG$ only on-shell. In this method, Z has explicit coupling-constant and field dependence. Note that in order for $(i\partial/\partial y - Z)\psi(y)$ to have a solution $\psi(y)$ whose y-dependence cancels in an action of the form $\int dx\, d\theta \bar{\psi}(y)(i\overleftrightarrow{\partial}/\partial y - 2mG)\psi(y)$ (see sect. 3) to preserve supersymmetry, $Z\psi$ must be linear in ψ. Therefore, *self-interactions are not allowed for fields with interaction-dependent central charges*. However, the mass-dependent form of Z has no such restriction. Another difference of the interaction-dependent method is the auxiliary field structure. For this method *the number of auxiliary components always equals the number of physical components*. In fact, the auxiliary field structure of the dimensionally reduced theory is in this case *totally unrelated* to the auxiliary fields of the higher-dimensional theory: since the higher-dimensional field equations were imposed as constraints, *all higher-dimensional auxiliary fields have been eliminated* (since their equations $\psi = 0$ have no $i\partial/\partial y$'s). Commuting physical component fields have equations of the form $(\partial^2/\partial y^2 + Z^2)\psi(y) = 0$, and so are doubled as $\psi(y) = \cos(Zy)\psi(0) + (\sin(Zy)/Z)\psi'(0)$, with $\psi(0)$ becoming a physical field and $\psi'(0)$ becoming an auxiliary field. Anticommuting component fields, which consist of an equal number of physical and auxiliary components (due to their field equations being linear in derivatives) solve $(i\partial/\partial y - Z)\psi(y) = 0$ as $\psi(y) = e^{-iZy}\psi(0)$, and are not doubled. Furthermore, this method can also be applied to theories without a suitable G: after dimensional reduction, one obtains a massless theory which may have fewer auxiliary fields but cannot have self-interactions.

2. Mass-dependent Z

In order to illustrate these features in more detail, we will first consider the example of the three-dimensional O(2) chiral multiplet, as obtained by dimensional reduction from the four-dimensional O(1) chiral multiplet. In the chiral representation, the action for a massless 4D O(1) chiral multiplet coupled to abelian vector multiplets (which will couple to the central charges) can be written as [9]

$$S = \int d^4x\, d^4\theta \bar{\phi} e^{2(U+VG)}\phi + \frac{1}{256g^2}\int d^4x\, d^2\theta W^\alpha W_\alpha ,$$

$$\bar{\partial}_\alpha \phi = 0, \qquad V = \bar{V}, \qquad U = \theta^\alpha \bar{\theta}^\beta \sigma^a_{\alpha\beta} i\partial_a , \qquad W_\alpha G = \bar{\partial}^2 e^{-2(U+VG)}\partial_\alpha e^{2(U+VG)} . \tag{2.1}$$

Here $VG = V^i G_i$, $W_\alpha G = W_\alpha^i G_i$, etc., with generators $G_i = G_i^\dagger$ which act as $[G_i, G_j] = G_i V_j = 0$, $G_i \phi_j = (G_i)_j{}^k \phi_k$. [We will generally suppress these indices to avoid confusion with O(N) indices.]

The m-dependent form of dimensional reduction is performed by choosing the explicit x^3 dependence (for a general superfield ψ)

$$\psi(x,\theta,x^3) = \mathrm{e}^{-imGx^3}\psi(x,\theta,0), \qquad \left(mG = m^i G_i\right) \tag{2.2a}$$

(where now $x = (x^0, x^1, x^2)$), or explicitly

$$\phi(x,\theta,x^3) = \mathrm{e}^{-imGx^3}\phi(x,\theta,0), \qquad V(x,\theta,x^3) = V(x,\theta,0). \tag{2.2b}$$

We then drop $\int \mathrm{d}x^3$ in S, which is already x^3 independent after $\int \mathrm{d}x^0 \mathrm{d}x^1 \mathrm{d}x^2 \mathrm{d}\theta$ (which is necessary to preserve supersymmetry: $\{Q, Q\} \sim \partial_3 + \dots$). The effects of the m-dependence can most easily be seen by performing the transformation

$$A' = \mathrm{e}^{+imGx^3} A \mathrm{e}^{-imGx^3} \tag{2.3}$$

on all fields and operators, thus removing all x^3 dependence. The net effect of eqs. (2.2) and (2.3) is thus

$$\begin{aligned} &\phi(x,\theta,x^3) \to \phi(x,\theta), \\ &U + VG \to \theta^\alpha \bar{\theta}^\beta \gamma^a_{\alpha\beta} i\partial_a + \theta^\alpha \bar{\theta}^\beta C_{\beta\alpha} mG + V(x,\theta)G, \\ &\left(Q_\alpha, \bar{Q}_\alpha\right) \to \left(\partial_\alpha, \bar{\partial}_\alpha - 2\theta^\beta \gamma^a_{\beta\alpha} i\partial_a - 2\theta^\beta C_{\alpha\beta} mG\right), \end{aligned} \tag{2.4}$$

where $\sigma^a_{\alpha\dot{\beta}} \to (\gamma^a_{\alpha\beta}, C_{\beta\alpha})$ in terms of the 3D Dirac and charge-conjugation matrices. (For notation, see ref. [6].) In terms of the 3D O(2) supersymmetry generators

$$Q_{1\alpha} = \sqrt{\tfrac{1}{2}}\left(Q_\alpha + \bar{Q}_\alpha\right), \qquad Q_{2\alpha} = \sqrt{\tfrac{1}{2}}\, i\left(Q_\alpha - \bar{Q}_\alpha\right), \tag{2.5a}$$

we thus have

$$\{Q_{i\alpha}, Q_{j\beta}\} = -2\delta_{ij}\gamma^a_{\alpha\beta} i\partial_a + 2\delta_{ij} C_{\alpha\beta} mG, \tag{2.5b}$$

which is merely the result of $i\partial_3 \to mG$ under (2.3). We thus have $Z = mG$ as a central charge, on- and off-shell.

To study the modification of the action it is convenient to reduce the 3D O(2) superfields to 3D O(1) superfields [5]. In order for the O(1) supersymmetry to take its usual form, we modify the transformation (2.3) to

$$\begin{aligned} &M = \exp(-\zeta^\alpha i\partial_\alpha)\exp(imGx^3)\exp\left(-\theta^2_{(4)} i\partial_3\right), \\ &\phi \to M\phi, \qquad \mathrm{e}^{2(U+VG)} \to M^{\dagger -1}\mathrm{e}^{2(U+VG)}M^{-1}, \qquad Q \to MQM^{-1}. \end{aligned} \tag{2.6}$$

Here we have written the 4D $\theta^\alpha_{(4)} = \sqrt{\tfrac{1}{2}}\,(\theta^\alpha + i\zeta^\alpha)$ in terms of the 3D Majorana θ^α which is the argument of 3D O(1) superfields, and the parameter ζ^α for expansion of 3D O(2) superfields in terms of 3D O(1) superfields. The factor $\exp(-i\zeta^\alpha\partial_\alpha)$ turns the argument $\theta^\alpha_{(4)}$ of ϕ into $\sqrt{\tfrac{1}{2}}\;\theta^\alpha$, so that it automatically becomes an O(1) superfield, as described in ref. [5]. The factor $\exp(-\theta^2_{(4)} i\partial_3)$ has been included to remove explicit θ^α dependence from the O(1) operators. After partially choosing a gauge for V [5], the O(2)$\rightarrow$O(1) reduction becomes (with only slight modification of the methods of ref. [5])

$$
\begin{aligned}
&\phi \rightarrow \phi\,,\\
&U + V \rightarrow \zeta^\alpha i \nabla_\alpha + \zeta^2(\chi + m)G\,, \qquad (\nabla_\alpha = D_\alpha - i\Gamma_\alpha G)\,,\\
&Q_{i\alpha} \rightarrow (Q_\alpha, \delta_\alpha + i\nabla_\alpha - 2i\theta_\alpha m G)\,, \qquad (\delta_\alpha = \partial/\partial\zeta^\alpha)\,. \qquad (2.7)
\end{aligned}
$$

ϕ is now a complex 3D O(1) superfield describing a 3D O(2) chiral multiplet; Γ_α and χ are Majorana spinor and real scalar 3D O(1) superfields describing a 3D O(2) vector multiplet. The action of eq. (2.1) becomes

$$
\begin{aligned}
S = \int \mathrm{d}^3x\,\mathrm{d}^2\theta \Big[&(\nabla^\alpha\bar\phi)(\nabla_\alpha\phi) - 2\bar\phi(\chi + m)G\phi \\
&+ \frac{1}{4g^2}(D^\alpha\chi)(D_\alpha\chi) - \frac{1}{144g^2} W^\alpha W_\alpha \Big]\,,
\end{aligned}
$$

$$
W_\alpha G = i\big[\nabla^\beta, \{\nabla_\beta, \nabla_\alpha\}\big] = 3D^\beta D_\alpha \Gamma_\beta G\,. \qquad (2.8)
$$

The supersymmetry transformations are

$$
\begin{aligned}
&\delta_1(\phi, \Gamma_\alpha, \chi) = \varepsilon_1^\beta Q_\beta(\phi, \Gamma_\alpha, \chi)\,,\\
&\delta_2(\phi, \Gamma_\alpha, \chi) = \big(i\varepsilon_2^\alpha(\nabla_\alpha - 2mG\theta_\alpha)\phi, -2\varepsilon_{2\alpha}\chi, -\tfrac{1}{6}\varepsilon_2^\alpha W_\alpha\big)\,. \qquad (2.9)
\end{aligned}
$$

The central charge of eq. (2.5b) is due to the $-2imG\varepsilon_2^\alpha\theta_\alpha\phi$ term in $\delta_2\phi$. Due to our partial choice of gauge, there is also a field-dependent gauge transformation in $\{Q, Q\}$, as manifested in Q by the gauge covariantization of $\delta_2\phi$ and $\delta_2\chi$. However, in an arbitrary gauge $Q_{2\alpha}$ would be totally field independent [see eq. (2.4)], which is not true for the interaction-dependent type of Z, as we will see in sect. 3. Note that it is also possible to introduce ϕ self-interactions with this mass-dependent type of Z: e.g., if we use a doublet ϕ_i, with $G\phi_1 = \phi_1$ and $G\phi_2 = -2\phi_2$, then a gauge-invariant self-interaction term $\lambda \int \mathrm{d}^4x\,\mathrm{d}^2\theta(\phi_1)^2\phi_2 + \text{h.c.}$ can be introduced into (2.1), which will become $2\int \mathrm{d}^3x\,\mathrm{d}^2\theta[\lambda(\phi_1)^2\phi_2 + \text{h.c.}]$ upon dimensional reduction.

3. Interaction-dependent Z

We will now obtain the interaction-dependent type of Z for the example of sect. 2. We again start with the action of eq. (2.1), but instead of solving the equation $(i\partial_3 - mG)\psi = 0$ [with solution eq. (2.2a)], we solve the field equations (*only* for the fields which are to be given central charges) which follow from eq. (2.1). We have thus reversed the roles of constraints and field equations: in both cases we have the final field equation $(A - mG)\psi = 0$ (for some operator A), but for the mass-dependent case $(i\partial_3 - mG)\psi = 0$ is a constraint and $(i\partial_3 - A)\psi = 0$ a field equation, whereas the reverse is true for the interaction-dependent case. (However, for the field V, which will not obtain a central charge, we constrain $i\partial_3 V = 0$ in *both* cases.)

For the interaction-dependent Z we thus have the constraints

$$\partial^2 e^{2(U+VG)}\phi = 0\,, \qquad i\partial_3 V = 0\,. \tag{3.1}$$

In order to solve the ϕ constraint we will go directly to 3D O(1) superfields, making the $m = 0$ case of the transformation of eq. (2.6). Eq. (2.7) then becomes

$$\begin{aligned}
&\phi\big(x, \theta^{\alpha}_{(4)}, x^3\big) \to \phi(x, \theta^{\alpha}, x^3)\,,\\
&U + V \to \zeta^{\alpha} i\nabla_{\alpha} + \zeta^2(\chi G + i\partial_3)\,,\\
&Q_{i\alpha} \to (Q_{\alpha}, \delta_{\alpha} + i\nabla_{\alpha} + 2\theta_{\alpha}\partial_3)\,.
\end{aligned} \tag{3.2}$$

The solution of eq. (3.1) is now found to be

$$\phi(x, \theta^{\alpha}, x^3) = \exp\left[\tfrac{1}{2} i x^3 (\nabla^2 + 2\chi G)\right]\phi(x, \theta^{\alpha})\,. \tag{3.3}$$

$[V(x, \theta^{\alpha}, x^3) \to V(x, \theta^{\alpha}, 0)$ as given by eq. (3.2).]

In order to obtain the field equation $(i\partial_3 - mG)\phi = 0$, we choose the action

$$S = \int d^3x\, d^2\theta \left[\bar{\phi}(i\overleftrightarrow{\partial}_3 - 2mG)\phi + \frac{1}{4g^2}(D^{\alpha}\chi)(D_{\alpha}\chi) - \frac{1}{144g^2} W^{\alpha}W_{\alpha}\right], \tag{3.4}$$

where the V part of the action is found from the V part of eq. (2.1). (Remember that fields without central charges are treated the same in this method as in the method of sect. 2.) The ϕ part of the action is independent of x^3 (which would not be true if $i\partial_3\phi$ were non-linear in ϕ):

$$\begin{aligned}
&\int d^3x\, d^2\theta\, \bar{\phi}(x, \theta, x^3)(i\overleftrightarrow{\partial}_3 - 2mG)\phi(x, \theta, x^3)\\
&\quad = \int d^3x\, d^2\theta\, \bar{\phi}(x, \theta, 0)\left[-\nabla^2 - 2(\chi + m)G\right]\phi(x, \theta, 0)\,,
\end{aligned} \tag{3.5}$$

so that eq. (3.4) gives the same action as eq. (2.8). In general, the action and the unconstrained fields obtained for the two different types of Z will not be directly related: more complicated cases will be discussed in sect. 4. The supersymmetry transformations are [from eqs. (3.2) and (3.3)]

$$\delta_1(\phi,\Gamma_\alpha,\chi)=\varepsilon_1^\beta Q_\beta(\phi,\Gamma_\alpha,\chi),$$

$$\delta_2(\phi,\Gamma_\alpha,\chi)=\left(i\varepsilon_2^\alpha[\nabla_\alpha+\theta_\alpha(\nabla^2+2\chi G)]\phi,-2\varepsilon_{2\alpha}\chi,-\tfrac{1}{6}\varepsilon_2^\alpha W_\alpha\right). \tag{3.6}$$

Note that since $Q_{2\alpha}$ is now field dependent (even in an arbitrary gauge), $\{Q,Q\}$ also includes the effect of one Q on the fields contained in the other Q. We now obtain eq. (2.5b) with mG replaced by $-\frac{1}{2}(\nabla^2+2\chi G)$ (when acting on ϕ; 0 on Γ_α and χ).

For the theory treated in this and the previous sections, the interaction-dependent-Z formalism has obtained results similar to those for the mass-dependent Z, but with the following drawbacks:

(i) fields with $Z\neq 0$ are treated differently from those with $Z=0$, with regard to both constraints and field equations;

(ii) the supersymmetry generators are field dependent, even in an arbitrary gauge;

(iii) self-interactions are not allowed for fields with $Z\neq 0$. In general, the two kinds of Z also produce different auxiliary field structures, as will be illustrated in sect. 4.

4. Four and five dimensions

By applying the methods of sects. 2, 3, massless six-dimensional multiplets can be reduced to five-dimensional multiplets with central charges, which can, in turn, be reduced to four dimensions by ordinary dimensional reduction. The only 4D O(2) multiplets whose auxiliary fields are presently known which can be derived by ordinary dimensional reduction (no Z's) from 6D O(1) multiplets are the massless vector multiplet [10–12, 4, 5] and the massless tensor multiplet [10, 13]. (The $m=0$ case of the action of ref. [10] describes the latter multiplet, with the gauge transformation $\delta\phi=\bar{D}^4 D^{(i\alpha}D^{j)}{}_\alpha K_{(ij)}, K_{(ij)}=\bar{K}_{(ij)}$, as can be shown using the projection operator of ref. [12]. This is due to the "duality", in terms of physical components$\leftrightarrow$auxiliary components, of the O(2) vector and tensor multiplets, as is also true in O(1) [14].) It is interesting to note that the 6D O(1) tensor multiplet, consisting of a fourth-rank antisymmetric-tensor gauge field, a scalar isovector, and a spinor isospinor, has sufficient auxiliary components in its physical fields, and needs no additional auxiliary fields, just as is true for the 4D O(1) tensor multiplet [14]. These two 6D O(1) multiplets do not have G's to allow construction of on-shell central charges, except by using more than one of the same multiplet and choosing a G which mixes them. (Also, a mass term cannot be constructed for either of these multiplets by itself, due to the fact that Weyl spinor indices cannot be raised or lowered in six dimensions. However, the two can be coupled together with a mass term, since one has a contravariant Weyl spinor field and the other a covariant one.)

More general cases can now be treated with interaction-dependent Z's, since for this type of Z it is only necessary to know the *on-shell* 6D (or 10D) theory, with superfields [4] or without (since the field equations will be applied as constraints). The 6D O(1) chiral (scalar) multiplet gives a 4D O(2) multiplet closely analogous to the example of sects. 2, 3. (Sohnius [3] has essentially done the reduction of the 5D O(1) multiplet, but 6D O(1) can be treated by a slight modification, and the results are the same.) One can also reduce the 6D O(1) and 6D O(2) (= 10D O(1): reduce to 9D O(1), then by ordinary reduction to 4D O(4)) vector multiplets to 4D multiplets with $Z = 0$ on-shell and $Z \neq 0$ off shell, since this case does not require a G. (The corresponding 5D O(1) and 9D O(1) multiplets have no auxiliary fields: the Bose fields of these multiplets are a vector and a $(D-2)$-rank antisymmetric-tensor gauge field.) One then obtains the abelian 4D O(2) and O(4) vector multiplets of Sohnius, Stelle, and West [15]. (The USp(4) = O(5) symmetry of the latter model is due to ordinary reduction from 9 dimensions.) However, as explained in sect. 3, such Z's do not allow self-interactions, so these theories are apparently not generalizable to the non-abelian case. In any case, they would not be related to 4D O(2) and 4D O(4) supersymmetric Yang-Mills, since the former theories contain an antisymmetric-tensor gauge field, which cannot have dimensionless couplings [16] (as can directly be seen from the form of its duality transformation which relates it to an equivalent scalar field [17]), whereas the latter theories contain only dimensionless couplings. Also, the auxiliary field structure of theories with interaction-dependent Z's is not related to those without Z's: in particular, the 4D O(2) chiral multiplet without Z's must have an auxiliary field with spin 1 (the corresponding tensor multiplet has spin-1 auxiliary components in the tensor gauge field), and the 4D O(4) vector multiplet an auxiliary field with spin 2. (This can be seen by applying an off-shell extension of the Salam-Strathdee state counting [18] to the case $m = Z = 0$: off-shell components for $m = 0$ but $p^2 \neq 0$ have the same counting as on-shell massive states with $p^2 = m^2 \neq 0$.)

Note added

After this work was completed, we became aware of an interesting preprint by M.F. Sohnius, K.S. Stelle and P.C. West (Dimensional reduction by Legendre transformation generates off-shell supersymmetric Yang-Mills theories, Imperial College preprint ICTP/79-80/22, (March, 1980)). This paper, by use of a hamiltonian approach, shows that actions of the form of our (3.4) are in fact correct even in the case of self-interactions. However, their method produces a non-local lagrangian when applied to Yang-Mills theories. The solution to the constraint on the 5D gauge field has the following form in terms of the 4D vector and antisymmetric-tensor gauge fields:

$$A_a^{(5)}(x^4) = A_a + x^4\left(\nabla^b{}^*A_{ab} - \tfrac{1}{2}i\nabla_a(\nabla^2)^{-1}\left[F^{bc}, {}^*A_{bc}\right]\right) + \mathrm{O}\left((x^4)^2\right)$$

(ignoring 5D fields other than $A_a^{(5)}$), due to the constraint $\nabla^{(5)a}(\partial/\partial x^4)A_a^{(5)} = 0$. The $A_a - A_{ab}$ part of the 4D lagrangian thus becomes

$$L = -\tfrac{1}{4}F_{ab}^2 - \tfrac{1}{2}\left(\nabla^b{}^*A_{ab}\right)^2 + \tfrac{1}{8}\left[F^{ab}, {}^*A_{ab}\right](\nabla^2)^{-1}\left[F^{ab}, {}^*A_{ab}\right].$$

In a local lagrangian, the dimensions of the A_{ab} couplings (A_{ab} couples only through its field strength) require coupling constants with dimensions of inverse mass. Since the only coupling here is dimensionless, the non-local operator $(\nabla^2)^{-1}$ is required to couple A_{ab} to A_a. Therefore, local self-interactions cannot generally be introduced for interaction-dependent Z's (with the exception of self-interactions for scalar multiplets).

References

[1] R. Haag, J. Łopuszanski, and M.F. Sohnius, Nucl. Phys. B88 (1975) 257
[2] W. Nahm, Nucl. Phys. B135 (1978) 149
[3] M.F. Sohnius, Nucl Phys. B138 (1978) 109
[4] W. Siegel, Phys. Lett. 80B (1979) 220
[5] W. Siegel, Nucl. Phys. B156 (1979) 135
[6] S.J. Gates, Jr., and W. Siegel, Nucl. Phys. B163 (1980) 519
[7] W. Siegel, Constraints on superspace, Supergravity, Stony Brook, 1979 (North-Holland, Amsterdam, 1979) p. 71
[8] J. Scherk and J H. Schwartz, Phys Lett 82B (1979) 60; Nucl. Phys. B153 (1979) 61;
E. Cremmer, J. Scherk and J.H Schwarz, Phys. Lett. 84B (1979) 83
[9] W. Siegel, Supergravity superfields without a supermetric, Harvard preprint HUTP-77/A068 (November, 1977);
W. Siegel and S.J. Gates, Jr., Nucl. Phys. B147 (1979) 77
[10] J. Wess, Acta Phys. Austriaca 41 (1975) 209
[11] R. Grimm, M. Sohnius and J Wess, Nucl. Phys. B133 (1978) 275
[12] L. Mezincescu, On the superfield formulation of O(2) supersymmetry, JINR preprint P2-12572 (June, 1979), in Russian
[13] B. de Wit and J.W. van Holten, Nucl. Phys. B155 (1979) 530
[14] W. Siegel, Phys. Lett. 85B (1979) 333
[15] M. Sohnius, K.S. Stelle and P.C. West, Off-mass-shell formulation of extended supersymmetric gauge theories, Imperial College preprint ICTP/79-80/17 (February, 1980)
[16] V.I. Ogievetsky and I V. Polubarinov, Sov J. Nucl. Phys. 4 (1967) 156
[17] E. Cremmer and J. Scherk, Nucl Phys. B72 (1974) 117;
P.K. Townsend, Local supersymmetry without gravity, Supergravity, Stony Brook, 1979 (North-Holland, Amsterdam, 1979) p. 37
[18] A. Salam and J. Strathdee, Nucl. Phys. B80 (1974) 499