Wess-Zumino-Novikov-Witten Models Based on Lie Superalgebras

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Abstract

The affine current algebra for Lie superalgebras is examined. The bilinear invariant forms of the Lie superalgebra can be either degenerate or non-degenerate. We give the conditions for a Virasoro construction, in which the currents are primary fields of weight one, to exist. In certain cases, the Virasoro central charge is an integer equal to the super dimension of the group supermanifold. A Wess-Zumino-Novikov-Witten action based on these Lie superalgebras is also found.

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1 Introduction

In this note we explore from a mathematical point of view a class of exact conformal field theories based on Wess-Zumino-Novikov-Witten (WZNW) built on Lie superalgebras. The group supermanifold corresponding to these Lie superalgebras could be neither compact nor semi-simple. In these models the Virasoro central charge does, in general, depend on the level of the affine Lie superalgebra. However, for certain cases the central charge is equal to the super dimension (the dimension of the bosonic part of the Lie superalgebra minus the dimension of the fermionic part) of the group supermanifold.

The motivation behind this study stems from the recent spate of interest in constructing WZNW models based on non-semi-simple groups [1–8]. This is because these models lead to exact string backgrounds having a target space dimension equal to the integer Virasoro central charge of the affine non-semi-simple algebra [3,4,5,6]. The first of these models, constructed by Nappi and Witten [1], was based on a central extension of the two-dimensional euclidean group and describes a string propagating in a homogeneous four-dimensional spacetime in the background of a gravitational plane wave. This construction of WZNW models for non-semi-simple groups was subsequently extended to other groups and generalised to higher dimensions [3,4,7,8]. Furthermore, all these models describe the propagation of strings on a target space whose metric possess a covariantly null Killing vector [9–14].

In ordinary WZNW models the invertibility of the Cartan-Killing invariant bilinear form is crucial for the Sugawara construction of the corresponding stress tensor. But since for non-semi-simple groups this bilinear form is degenerate, the Sugawara construction (that is a Virasoro construction in which the currents are primary fields of weight one) does not exist. However, for all the non-semi-simple algebras studied so far [1,3,4,6,7] it has been possible to find non-degenerate invariant bilinear forms. These invertible bilinear forms allow for a Sugawara-type construction leading to integer Virasoro central charges.

In [5] we gave the conditions for an affine Lie algebra (semi-simple or non-semi-simple) to admit a Sugawara-type construction with respect to which the currents are primary fields of conformal dimension one. An expression for the Virasoro central charge was also found and from which one sees that this central charge is not necessarily an integer number. A search, carried out in [6], for non-semi-simple Sugawara-type construction having non-integral value of the central charge shows that this construction necessarily factorises into a semi-simple standard Sugawara construction and a non-semi-simple one with integral
central charge.

In this paper we extend the analyses and constructions of ref.\[5\] to current algebras based on Lie superalgebras. In an affine Lie superalgebra one encounters two invariant bilinear forms; one symmetric and the other antisymmetric. In the literature super Sugawara constructions corresponding to affine Lie superalgebras were extensively studied in [15]. The existence of these super Sugawara construction relies on the invertibility of the two bilinear forms entering the affine Lie superalgebra. The aim of this paper is to explore the possibility of finding super Sugawara-type constructions even when the two bilinear forms of the affine Lie superalgebra are degenerate. In other words, using other invertible bilinear forms, we construct a solution of the super master equation [16] with the additional property that all the currents are primary field of conformal dimension one.

The paper is organised as follows: In section two we review the Lie supercurrent algebra and construct the energy-momentum tensor. The conditions under which the generators of the Lie supercurrent algebra are primary fields of conformal dimension one, are spelled out. Section three is dedicated to the construction of a WZWN action based on Lie superalgebras.

## 2 The Lie Supercurrent Algebra

A Lie superalgebra, $\mathcal{G}$, is expressed in terms of a set of bosonic generators, $t^a \in \mathcal{G}_0$, together with a set of fermionic generators, $S^\alpha \in \mathcal{G}_1$. Its affine extension (for which $t^a$ and $S^\alpha$ are the zero modes) is given by the commutation relations [15]

\[
\begin{align*}
[t_m^a, t_n^b] &= f_{c}^{ab} t_{m+n}^c + mg^{ab} \delta_{m+n,0} \\
[t_m^a, S^\alpha_r] &= N_{\beta}^{\alpha} S^\beta_{m+r} \\
\{S^\alpha_r, S^\beta_s\} &= R_{a}^{\alpha\beta} t_{r+s}^a + r\varphi^{\alpha\beta} \delta_{r+s,0} .
\end{align*}
\]

(2.1)

Here $g^{ab}$ and $\varphi^{\alpha\beta}$ are two invariant bilinear forms of the Lie superalgebra and they are, respectively, symmetric and antisymmetric. We do not assume these two bilinear forms to be invertible.

The structure constants of the above algebra are related by the super Jacobi identities. Firstly, we have the usual relations that one encounters in bosonic Lie algebras, namely

\[
\begin{align*}
f_{c}^{ab} f_{e}^{cd} + f_{c}^{da} f_{e}^{cb} + f_{c}^{bd} f_{e}^{ca} &= 0 \\
g^{ab} f_{b}^{cd} + g^{cb} f_{b}^{da} &= 0 .
\end{align*}
\]

(2.2)
Secondly, the symplectic form $\varphi^{\alpha\beta}$ satisfies
\[ N^{\alpha}_\beta \varphi^{\beta\gamma} - N^{\alpha\gamma}_\beta \varphi^{\beta\alpha} = 0 \quad (2.3) \]

Thirdly, the structure constants $N^{\alpha}_\beta$ and $R^\alpha_\beta$ have to fulfill
\[ N^{\alpha}_\beta R^\beta_\gamma + N^{\alpha\beta}_\gamma = f^{a\gamma}_{bc} R^\alpha_\beta \]
\[ N^{\alpha}_\beta N^\beta_\gamma - N^{\alpha\gamma}_\beta = -f^{ab}_c N^\alpha_\beta \quad (2.4) \]

The last equation means that the matrices $(-N^\alpha_\beta)$ define a representation of the bosonic part of the Lie superalgebra. The structure constants $R^\alpha_\beta$ are related to $N^{\alpha}_\beta$ by the relation
\[ R^\alpha_\beta g^{ab} - N^b_\alpha \varphi^{\tau\beta} = 0 \quad (2.5) \]

Notice that $R^\alpha_\beta$ is indeed symmetric in $\alpha$ and $\beta$. Finally, we have
\[ R^\alpha_\beta N^{\alpha\gamma}_\tau + R^\gamma_\alpha N^{\alpha\beta}_\tau + R^{\beta\gamma}_\alpha N^{\alpha\alpha}_\tau = 0 \quad (2.6) \]

In terms of operator product expansions, the above commutation relations are equivalent to
\[ J^a(z)J^b(w) = \frac{g^{ab}}{(z-w)^2} + f^{ab}_{c} \frac{J^c(w)}{(z-w)} \]
\[ J^a(z)S^\alpha(w) = N^{\alpha\alpha}_{\tau} \frac{S^\tau(w)}{(z-w)} \]
\[ S^\alpha(z)S^\beta(w) = \frac{\psi^{\alpha\beta}_\tau}{(z-w)^2} + R^\alpha_\beta \frac{J^c(w)}{(z-w)} \quad (2.7) \]

Here $J^a(z) = \sum_n t^a_n z^{-n-1}$ are the bosonic currents of the even part of the Lie superalgebra, and $S^\alpha(z) = \sum_r S^\alpha_r z^{-r-1}$ constitute the fermionic currents of the odd part of the Lie superalgebra. The label $n$ is an integer while $r$ can be either integer or half-integer.

Let us assume now that we have two bilinear forms $\Omega^{ab}$ and $\psi^{\alpha\beta}$ which are respectively symmetric and antisymmetric and invertible, namely $\Omega^{ab}\Omega_{bc} = \delta^a_c$ and $\psi^{\alpha\beta}\psi_{\beta\tau} = \delta^\alpha_\tau$. These two invertible bilinear forms are also assumed to satisfy
\[ \Omega^{ab} f^{cd}_b + \Omega^{cb} f^{ad}_b = 0 \]
\[ N^{\alpha\alpha}_\beta \psi^{\beta\gamma} - N^{\alpha\gamma}_\beta \psi^{\beta\alpha} = 0 \]
\[ R^\alpha_\beta \Omega^{ab} - N^{\alpha\alpha}_{\tau} \psi^{\tau\beta} = 0 \quad (2.8) \]

\[ ^1\text{To get the usual level } k, \text{ one simply scales } g^{ab} \text{ and } \varphi^{\alpha\beta} \text{ by a factor } k. \]
It is then easy to verify that the Casimir operator of the Lie superalgebra (the algebra of the zero modes \( t^a \) and \( S^\alpha \) in (2.1)) is given by

\[
Q = \Omega_{ab} t^a t^b + \psi_{\alpha\beta} S^\alpha S^\beta .
\]  

(2.9)

The energy-momentum tensor is taken to be given by the following quadratic combination of bosonic and fermionic currents \[16\]

\[
T(z) = C_{ab} : J^a J^b : (z) + D_{\alpha\beta} : S^\alpha S^\beta : (z) ,
\]  

(2.10)

where \( C_{ab} \) is symmetric and \( D_{\alpha\beta} \) is antisymmetric. We would like now to determine these last two tensors by requiring that both \( J^a(z) \) and \( S^\alpha(z) \) are primary fields of conformal dimension equal to one with respect to the above energy-momentum tensor.

In order to perform the different operator product expansions, we define the normal ordered product of two operators \( A(z) \) and \( B(z) \) at coincident points as \[17\]

\[
\langle AB : (z) \rangle = \frac{1}{2\pi i} \oint_{C_z} \frac{dx}{x - z} A(x)B(z) ,
\]  

(2.11)

where the contour of integration, \( C_z \), surrounds the point \( z \). This definition leads to the following form of Wick’s theorem for calculating the product expansion of \( A(z) \) with a composite field : \( BC : (w) \)

\[
\langle A(z) : BC : (w) \rangle = \frac{1}{2\pi i} \oint_{C_w} \frac{dx}{x - w} \{ A(z)B(x)C(w) + (-1)^{BC} A(z)C(w)B(x) \} ,
\]  

(2.12)

where \((-1)^{BC} = -1 \) iff both \( B \) and \( C \) are fermionic fields and the contraction ( \( \langle \rangle \) ) stands for the singular part in the expansion of the product of two operators at distinct points.

By requiring that the operator product expansion of \( T(z) \) with \( J^a(z) \) should be given by

\[
T(z)J^a(w) = \frac{J^a(w)}{(z - w)^2} + \frac{\partial J^a(w)}{(z - w)} \]

we get the following equations

\[
C_{cb} f^b_c + C_{eb} f^b_c = 0
\]

\[
D_{\alpha\tau} N_{\rho}^{\alpha\tau} - D_{\rho\tau} N_{\alpha}^{\alpha\tau} = 0
\]

\[
2C_{cb} g_{ba} + C_{bd} f^{ed} f^e_d + D_{\alpha\tau} N_{\rho}^{\alpha\alpha} R^{\rho\tau}_c = \delta^a_c .
\]  

(2.14)

On the other hand by demanding that the operator product expansion of \( T(z) \) with \( S^\alpha(z) \) should be given by

\[
T(z)S^\alpha(w) = \frac{S^\alpha(w)}{(z - w)^2} + \frac{\partial S^\alpha(w)}{(z - w)}
\]  

(2.15)
leads to

\[ D_{\alpha\beta}R^\beta_\alpha + C_{ab}N^b_\alpha N^a_\gamma = 0 \]
\[ 2D_{\alpha\tau}\varphi^{\tau\rho} + D_{\tau\gamma}R^\tau_\alpha N^a_\alpha + C_{ab}N^\alpha_\tau N^b_\alpha = \delta^\rho_\alpha \]  

(2.16)

The first two equations in (2.14) and the first equation in (2.16) are equivalent to the set of equations in (2.8) and are uniquely solved by

\[ C_{ab} = k\Omega_{ab} , \quad D_{\alpha\beta} = k\psi_{\alpha\beta} , \]  

(2.17)

where \(k\) is a constant. Using these solutions in the last equation in (2.14), leads to

\[ g^{ab} = -\frac{1}{2k} \left[ k \left( \gamma^{ab} - \theta^{ab} \right) - \Omega^{ab} \right] , \]  

(2.18)

where \(\gamma^{ab}\) and \(\theta^{ab}\) are defined by

\[ \gamma^{ab} = f^{eb}f^{ca} , \quad \theta^{ab} = N^\alpha_\tau N^b_\alpha . \]  

(2.19)

On the other hand, the second equation in (2.16) yields

\[ \varphi^{\alpha\beta} = -\frac{1}{2k} \left( 2k\Omega_{ab}N^\alpha_\tau N^b_\rho \psi^{\rho\sigma} - \psi^{\alpha\sigma} \right) . \]  

(2.20)

A straightforward calculation shows that \(g^{ab}\) and \(\varphi^{\alpha\beta}\), as given above, satisfy the relations dictated by the super Jacobi identities.

With these expressions for \(C_{ab}, g^{ab}, D_{\alpha\beta}\) and \(\varphi^{\alpha\beta}\), the above energy-momentum tensor satisfies the Virasoro algebra

\[ T(z)T(w) = \frac{c}{2(z-w)^2} + \frac{2T(w)}{(z-w)^2} + \frac{\partial T(w)}{(z-w)} . \]  

(2.21)

The central charge of this algebra is given by

\[ c = \dim(\mathcal{G}_0) - \dim(\mathcal{G}_1) - k\Omega_{ab} \left( \gamma^{ab} - 3\theta^{ab} \right) . \]  

(2.22)

Notice that if \(\Omega_{ab}\gamma^{ab} = \Omega_{ab}\theta^{ab} = 0\), then the central charge is given by an integer number equal to the super dimension of the Lie superalgebra. Before leaving this section, we would like to mention that equivalent results have been independently reached in [18].

3 The WZNW Action for Super Lie Algebras

An element in the Lie supergroup is expressed as the exponential

\[ g = e^{x+a}e^{y+S^a} , \]  

(3.1)
where $x_a$ are bosonic variables and $y_\alpha$ are fermionic ones. These variables are interpreted as coordinates on the group supermanifold. Notice that $y_\alpha S^\alpha = -S^\alpha y_\alpha$ since they are fermions. The WZNW action for Lie superalgebras is constructed out of the bosonic “gauge fields” $A_{ia}$ and the fermionic “gauge fields” $B_{i\alpha}$, defined for an element $g$ in the Lie supergroup via

$$g^{-1}\partial_i g = A_{ia} t^a + B_{i\alpha} S^\alpha .$$  \hfill (3.2)

Here $i, j, k, \ldots$ are world-sheet indices.

Under an infinitesimal transformation of the form

$$g \to g + hg + gl ,$$ \hfill (3.3)

where $h$ and $l$ are elements in the Lie supergroup and are written as

$$h = \omega_a t^a + \tilde{\omega}_\alpha S^\alpha$$

$$l = \theta_a t^a + \tilde{\theta}_\alpha S^\alpha .$$ \hfill (3.4)

The variation of the gauge fields is then given by

$$A_{ia} \to A_{ia} + \partial_i (\theta_a + \lambda_a) + A_{ib} (\theta_c + \lambda_c) f^{bc}_a - B_{i\alpha} (\theta_\beta + \tilde{\lambda}_\beta) R^\alpha_{\beta\beta}$$

$$B_{i\alpha} \to B_{i\alpha} + \partial_i (\tilde{\theta}_\alpha + \tilde{\lambda}_\alpha) + A_{ia} (\tilde{\theta}_\beta + \tilde{\lambda}_\beta) N^\alpha_{\beta\beta} - B_{i\alpha} (\theta_a + \lambda_a) N^a_{\beta\beta} ,$$ \hfill (3.5)

with $\lambda_a$ and $\tilde{\lambda}_\alpha$ defined by

$$g^{-1}hg = \lambda_a t^a + \tilde{\lambda}_\alpha S^\alpha .$$ \hfill (3.6)

An important property of the gauge field $A_{ia}$ and $B_{i\alpha}$ is that their curvatures vanish

$$F_{ija} = \partial_i A_{ja} - \partial_j A_{ia} + A_{ic} A_{jd} f^{cd}_a - B_{i\alpha} B_{j\beta} R^\alpha_{\beta\beta} = 0$$

$$H_{ija} = \partial_i B_{j\alpha} - \partial_j B_{i\alpha} + A_{ia} B_{j\beta} N^\alpha_{\beta\beta} - A_{ja} B_{i\beta} N^a_{\beta\beta} = 0 .$$ \hfill (3.7)

These last equations will be crucial in demonstrating the invariance of the action.

Let us also define the following quantities, which will also enter in determining the Noether currents associated to the above transformations

$$g^{-1}t^a g = V^a_b t^b + \tilde{V}^a_\alpha S^\alpha$$

$$g^{-1}S^\alpha g = W^\alpha_b t^b + \tilde{W}^\alpha_\beta S^\beta .$$ \hfill (3.8)

\textsuperscript{2}If $y_\alpha$ and $S^\alpha$ are matrices then their product is defined such that $y_\alpha S^\alpha = -S^\alpha y_\alpha$. See for example [19] for a review on operation on super matrices.
Here \((V^a_b, \tilde{W}^\alpha_\beta)\) are bosonic quantities while \((\tilde{V}^a_\alpha, W^\alpha_a)\) are fermionic ones. They satisfy

\[
\begin{align*}
\partial_i V^a_b &= V^a_c A^I_{cd} f^I_{bd} - \tilde{V}^a_\beta B^\alpha_b R^\alpha_\beta \\
\partial_i \tilde{V}^a_\alpha &= V^a_c B^\beta_\alpha N^\alpha_c - \tilde{V}^a_\alpha A^\alpha_{ic} N^c_\alpha \\
\partial_i W^\alpha_a &= W^\alpha_c A^I_{cd} F^I_{ca} - \tilde{W}^\alpha_\tau B^\beta_\tau R^\beta_a \\
\partial_i \tilde{W}^\alpha_\beta &= W^\alpha_c B^\nu_\beta N^\nu_c - \tilde{W}^\alpha_\nu A^\nu_{ic} N^c_\nu .
\end{align*}
\]

These last four equations have been used in proving the invariance of the action.

Let us now turn our attention to the gauge invariant action. This is given by

\[
I(g) = -\frac{k}{8\pi} \int_{\partial B} d^2 x \sqrt{-\eta} \epsilon^{ij} \left( g^{ab} A_{ia} A_{jb} - \varphi^{ab} B_{ia} B_{jb} \right)
+ \frac{i k}{12\pi} \int_B d^3 y \epsilon^{ijk} \left( g^{ab} (\partial_i A_{ja}) A_{kb} - \varphi^{ab} (\partial_i B_{ja}) B_{kb} \right),
\]

where \(B\) is a three-dimensional manifold whose boundary is the two-dimensional surface \(\partial B\). Here we do not require \(g^{ab}\) and \(\varphi^{ab}\) to be invertible. The only requirement on these two tensors is that they obey the relations given by the super Jacobi identities in the previous section. The above action has already appeared in [20,21] for non-degenerate \(g^{ab}\) and \(\varphi^{ab}\).

Using the fact that \(F_{ija}\) and \(H_{ija}\) vanish, and that \(g^{ab}\) and \(\varphi^{ab}\) are two invariants of the Lie superalgebra together with (3.9), the variation of the action is found to be given by

\[
\delta I(g) = -\frac{k}{4\pi} \int_{\partial B} d^2 x \left( \sqrt{-\eta} \epsilon^{ij} - i \epsilon^{ij} \right) \left( g^{ab} [\partial_i \omega_c A_{ja} V^c_b + \partial_i \bar{\omega}_a A_{ja} W^a_b + \partial_j \theta_a A_{ib}] - \varphi^{ab} \left[ \partial_i \omega_c B_{ja} \tilde{V}^c_\beta - \partial_i \bar{\omega}_a B_{ja} \tilde{W}^a_\beta + \partial_j \bar{\theta}_a B_{ib} \right] \right)
- \frac{i k}{2\pi} \int_{\partial B} d^2 x \epsilon^{ij} \left( g^{ab} \partial_i (\omega_c A_{ja} V^c_b + \bar{\omega}_a A_{ja} W^a_b) - \varphi^{ab} \partial_i \left( \omega_c B_{ja} \tilde{V}^c_\beta + \bar{\omega}_a B_{ja} \tilde{W}^a_\beta \right) \right).
\]

Neglecting the last term (surface term), we see that in complex coordinates \((z, \bar{z})\) such that \(\eta^{z \bar{z}} = 1\) and \(\epsilon^{z \bar{z}} = i\), the variation of the action vanishes if \(\theta_a = \theta_a(z)\), \(\bar{\theta}_a = \bar{\theta}_a(z)\) and \(\omega_a = \omega_a(z)\), \(\bar{\omega}_a = \bar{\omega}_a(z)\). The Noether currents associated to some Lie superalgebra elements, \(t^a\) and \(S^\alpha\) are given by

\[
J^a_z = -\frac{k}{2\pi} g^{ab} A_{zb} , \quad J^a_{\bar{z}} = -\frac{k}{2\pi} g^{bc} A_{cb} V^a_b + \frac{k}{2\pi} \varphi^{ab} B_{za} \tilde{V}^a_\beta \\
S^a_z = \frac{k}{2\pi} \varphi^{ab} B_{za} \tilde{B}^b_z , \quad S^a_{\bar{z}} = -\frac{k}{2\pi} g^{bc} A_{cb} W^a_b - \frac{k}{2\pi} \varphi^{ab} B_{za} \tilde{W}^a_\beta
\]

The sets of currents \((J^a_z, S^a_z)\) and \((J^a_{\bar{z}}, S^a_{\bar{z}})\) are, by virtue of the equations of motion, holomorphic and antiholomorphic currents, respectively. In the case when \(g^{ab}\) and \(\varphi^{ab}\) are invertible, these currents satisfy two commuting copies of the supercurrent algebra given in (2.7).
In summary, we have considered in this paper current algebras based on Lie superalgebras. We gave a systematic approach to constructing the energy-momentum tensor with respect to which the supercurrents are primary fields of conformal dimension one. The Virasoro central charge for certain cases is an integer number equal to the super dimension of the Lie superalgebra. A Wess-Zumino-Novikov-Witten action based on Lie superalgebras is also constructed.

The study carried out in this paper raises two questions. The first one is whether the WZNW action for Lie superalgebras would define some string backgrounds. In order to answer this question one has to determine the conformal invariance conditions (the beta functions) for a non-linear sigma model defined on a supermanifold. The second question concerns the connection between our construction and topological field theory where Lie superalgebras are of crucial importance [21,22]. It would also be very desirable to generalise the theorems of ref.[6] to Lie superalgebras. Work in this direction is already in progress in the first reference of [18].

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References

[1] C. Nappi and E. Witten, Phys. Rev. Lett. 71 (1993) 3751.

[2] E. Kiritsis and C. Kounnas, Phys. Lett. B320 (1994) 264;
E. Kiritsis, C. Kounnas and D. Lüst, “Superstring Gravitational Wave Backgrounds with Spacetime Supersymmetry”, preprint CERN-TH.7218/94, hep-th/9404114.

[3] K. Sfetsos, “Gauging a non-semi-simple WZW model”, preprint THU-93/30, hep-th/9311010.
K. Sfetsos, “Exact String Backgrounds from WZW Models Based on Non-semi-simple Groups”, preprint THU-93/31, hep-th/9311093.

[4] D. A. Olive, E. Rabinovici and A. Schwimmer, Phys. Lett. B321 (1994) 361.

[5] N. Mohammedi, Phys. Lett. B325 (1994) 371.
[6] J. M. Figueroa-O’Farrill and S. Stanciu, “Nonsemisimple Sugawara Constructions”, preprint QMW-PH-94-2, hep-th/9402033.

[7] A. Kehagias and P. A. A. Meessen, “Exact String backgrounds from a WZW Model Based on the Heisenberg Group”, preprint THEF-NYM-94.2, hep-th/9403041.

[8] I. Antoniadis and N. A. Obers, “Plane Gravitational Waves in String Theory”, preprint CPTH-A299.0494, hep-th/9403191.

[9] C. Klimčík and A. A. Tseytlin, “Duality invariant class of exact string backgrounds”, preprint CERN-TH.7069, hep-th/9311012.

[10] R. Guven, Phys. Lett. B191 (1987) 275.
[11] D. Amati and C. Klimčík, Phys. Lett. B219 (1989) 443.
[12] G. T. Horowitz and A. R. Steif, Phys. Rev. Lett. 64 (1990) 260.
[13] A. A. Tseytlin, Nucl. Phys. B390 (1993) 153; Phys. Rev. D47 (1990) 3421.
[14] C. Duval, Z. Horváth and P. A. Horváthy, Mod. Phys. Lett. A8 (1993) 3749; Phys. Lett. B313 (1993) 10; preprint hep-th/9404018.
[15] P. Goddard, D. Olive and G. Waterson, Comm. Math. Phys. 112 (1987) 591.
[16] A. Deckmyn and W. Troost, Nucl. Phys. B370 (1992) 231.
[17] F. A. Bais, P. Bouwkneget, K. S. Schoutens and M. Surridge, Nucl. Phys. B304 (1988) 371.
[18] J. M. Figueroa-O’Farrill and S. Stanciu, in preparation; A. Deckmyn, Ph.D. thesis, Katholieke Universiteit Leuven (1994).
[19] B. De Witt, “Supermanifolds”, Cambridge University Press, 1985; D. A. Leites, Russian Math. Surveys 35 (1980) 1.
[20] M. Henningson, Int. J. Mod. Phys. A6 (1991) 1137.
[21] J. M. Isidro and A. V. Ramallo, Nucl. Phys. B414 (1994) 715.
[22] M. Bershadsky and H. Ooguri, Phys. Lett. B229 (1989) 374.