Vector supersymmetry of the superstring in the super-Beltrami parametrization

A. Boresch#, M. Schweda*
and S.P. Sorella*

#CERN, Theory Division
CH-1211, Geneva 23 (Switzerland)

*Institut für Theoretische Physik
Technische Universität Wien
Wiedner Hauptstraße 8-10
1040 Wien (Austria)

Abstract

The vector supersymmetry recently found for the bosonic string is generalized to superstring theories quantized in the super-Beltrami parametrization.

1Supported in part by the "Fonds zur Förderung der Wissenschaftlichen Forschung", M008-Lise Meitner Fellowship
1 Introduction

Topological models \cite{1, 2, 3} are known to be characterized by a supersymmetric algebra of the Wess-Zumino type whose generators are identified respectively with the BRS symmetry and with the vector supersymmetry carrying a Lorentz index \cite{4}.

This supersymmetric structure turns out to be very useful in order to discuss the perturbative renormalization of these models. Indeed, as shown in \cite{5}, it provides a simple way for solving the descent equations corresponding to the integrated BRS cohomology, yielding then an algebraic characterization of anomalies and invariant counterterms for both Schwarz \cite{6} and Witten’s type \cite{7} topological models.

This paper is the continuation of a recent work \cite{8} where the vector supersymmetry has been established also in the case of the bosonic string quantized in the Beltrami parametrization \cite{11, 12}. In particular we show that the vector supersymmetry can be generalized to the (1, 1) and (1, 0) superstring theories and that, as in the bosonic case, it can be used for an algebraic characterization of the superdiffeomorphism anomaly.

In the following we shall adopt the so called super-Beltrami parametrization. We will make use of the results of \cite{9, 10} where this parametrization has been developed in a superfield as well as in a component field formalism. Let us recall that the Beltrami parametrization, introduced by \cite{11}, has been shown \cite{12} to be the most natural parametrization which exhibits the holomorphic factorization of the Green functions according to the Belavin-Polyakov-Zamolodchikov scheme \cite{13}. This holds true also for the super-Beltrami parametrization \cite{10}.

The work is organized as follows. Sect.2 is devoted to a brief summary of the BRS quantization of the superstring in the super-Beltrami parametrization. In Sect.3 we introduce the vector supersymmetry and we study the related Ward identities. Finally, in Sect.4 we solve the descent equations for the superdiffeomorphism anomaly. All the calculations are done for the case of the (1, 1) superstring, the (1, 0) case being easily recovered by means of a simple truncation procedure, as explained in Sect.2.
2 BRS-quantization of the superstring

Following [10], the superstring action in component fields and in the Wess-Zumino (WZ) gauge is given by:

\[-2iS_{inv} = \int dzd\bar{z} \frac{1}{1-\mu\bar{\mu}} \left( (\partial - \bar{\mu}\partial)X(\bar{\partial} - \mu\partial)X \right) + \int dzd\bar{z} \frac{1}{1-\mu\bar{\mu}} \left( \alpha\lambda(\partial - \bar{\mu}\partial)X + \alpha\bar{\lambda}(\bar{\partial} - \mu\partial)X + \frac{1}{2}\alpha\lambda\bar{\alpha}\bar{\lambda} \right) + \int dzd\bar{z} \left( \lambda(\partial - \mu\partial)\lambda + \bar{\lambda}(\bar{\partial} - \bar{\mu}\partial)\bar{\lambda} + (1 - \mu\bar{\mu})F^2 \right), \]  

with

\[\partial = \partial_z, \quad \bar{\partial} = \partial_{\bar{z}}.\]  

Here the \(\{X\}\) are the string coordinates, the \(\{\lambda\}\) denote the fermionic superpartners and the \(\{F\}\) are auxiliary fields. \(\mu\) stands for the conventional (i.e. bosonic) Beltrami differential and \(\alpha\), also called "Beltramino", is its fermionic superpartner. The doublets \((\mu, \alpha)\) and \((\bar{\mu}, \bar{\alpha})\) characterize the super-Beltrami parametrization which is a parametrization of the metric in the 2d-superspace [10]. For the two dimensional world sheet the line element takes the form

\[ds^2 \propto |dz + \mu d\bar{z}|^2.\]  

The quadratic part in the fields \(\{X\}\) of the expression (2.1) identifies the bosonic string whereas the quadratic terms in \(\{\lambda\}\) are the corresponding fermionic counterpart. Looking at the \((\lambda X)\)-terms one can see that the Beltramino \(\alpha\) plays the role of the gravitino or Rarita-Schwinger field which is present in other formulations of the superstring action.

Expression (2.1) represents the full \((1,1)\) superstring theory. According to [10], the corresponding \((1,0)\) version is easily obtained by means of the truncation \((\bar{\alpha} = \bar{\lambda} = F = 0)\), meaning that the supersymmetry is present only in one sector while in the other one has just the bosonic string. In the following we will always refer to the \((1,1)\) model, all the results being easily extended to the \((1,0)\) case by using the above truncation procedure.

The classical action (2.1) turns out to be invariant under the superdiffeomorphism transformations which are expressed by the following BRS transformations [10]

\[sX = \frac{1}{1-\mu\bar{\mu}} [(c - \mu\bar{c})\partial X + (\bar{c} - \bar{\mu}\tilde{c})\bar{\partial}X + \frac{1}{2}\alpha(\bar{c} - \bar{\mu}\tilde{c})\lambda] + \frac{1}{2} \epsilon\lambda, \]  

\[s\lambda = \frac{1}{1-\mu\bar{\mu}} [(c - \mu\bar{c})\partial\lambda + (\bar{c} - \bar{\mu}\tilde{c})\bar{\partial}\lambda] + \frac{1}{2} \left[ \partial c - \frac{(\bar{c} - \bar{\mu}\tilde{c})}{(1 - \mu\bar{\mu})}\partial\mu \right] \lambda \]

\[+ \frac{1}{2} \left[ \frac{\alpha(\bar{c} - \bar{\mu}\tilde{c})}{(1 - \mu\bar{\mu})} + \epsilon \right] DzX - \frac{i}{2} \left[ \frac{\bar{\alpha}(c - \mu\bar{c})}{(1 - \mu\bar{\mu})} + \bar{\epsilon} \right] F, \]

where \(c\) and \(\tilde{c}\) denote the conventional and the fermionic superpartners of the string coordinates, respectively.
\( s\bar{\lambda} = \text{c.c.} \), \hspace{1cm} \text{(2.6)}

\[
s F = \frac{1}{1-\mu\bar{\mu}} \left[ (c-\mu\bar{c})\partial F + (\bar{c}-\bar{\mu}c)\bar{\partial}F \right] + \frac{1}{2} \left[ \partial c - \frac{(\bar{c} - \bar{\mu}c)}{(1-\mu\bar{\mu})} \bar{\partial} \mu \right] F \\
+ \frac{1}{2} \left[ \bar{\partial} \bar{c} - \frac{(c - \mu\bar{c})}{(1-\mu\bar{\mu})} \partial \bar{\mu} \right] F \\
- \frac{i}{2} \left[ \left( \frac{\alpha(c - \mu\bar{c})}{1-\mu\bar{\mu}} + \epsilon \right) D_z\bar{\lambda} - \left( \frac{\bar{\alpha}(c - \mu\bar{c})}{1-\mu\bar{\mu}} + \bar{\epsilon} \right) D_z\lambda \right], \hspace{1cm} \text{(2.7)}
\]

\[
s\mu = (\bar{\partial} - \mu\bar{\partial} + \partial \mu)c + \frac{1}{2} \alpha \epsilon, \hspace{1cm} \text{(2.8)}
\]

\[
s\bar{\mu} = (\partial - \bar{\mu}\partial + \bar{\partial} \bar{\mu})\bar{c} + \frac{1}{2} \bar{\alpha} \bar{\epsilon}, \hspace{1cm} \text{(2.9)}
\]

\[
s\alpha = (\bar{\partial} - \mu\bar{\partial} + \frac{1}{2} \partial \mu)\epsilon + c\partial \alpha + \frac{1}{2} \alpha \partial c, \hspace{1cm} \text{(2.10)}
\]

\[
s\bar{\alpha} = (\partial - \bar{\mu}\partial + \frac{1}{2} \bar{\partial} \bar{\mu})\bar{\epsilon} + \bar{c}\partial \bar{\alpha} + \frac{1}{2} \bar{\alpha} \bar{\partial} \bar{c}, \hspace{1cm} \text{(2.11)}
\]

and

\[
s c = c\partial c - \frac{1}{4} \epsilon \bar{\epsilon}, \hspace{1cm} \text{(2.12)}
\]

\[
s \bar{c} = \bar{c}\partial \bar{c} - \frac{1}{4} \bar{\epsilon} \epsilon, \hspace{1cm} \text{(2.13)}
\]

\[
s \epsilon = c\partial \epsilon - \frac{1}{2} \epsilon \partial c, \hspace{1cm} \text{(2.14)}
\]

\[
s \bar{\epsilon} = \bar{c}\partial \bar{\epsilon} - \frac{1}{2} \bar{\epsilon} \bar{\partial} \bar{c}, \hspace{1cm} \text{(2.15)}
\]

so that

\[
s^2 = 0, \hspace{1cm} \text{(2.16)}
\]

with the supercovariant derivatives \( D_zX \) and \( D_z\bar{\lambda} \) defined as \[10\]

\[
D_zX \equiv \frac{1}{1-\mu\bar{\mu}} \left[ (\partial - \bar{\mu}\bar{\partial})X + \frac{1}{2} \bar{\mu}\alpha \lambda - \frac{1}{2} \bar{\alpha} \bar{\lambda} \right], \hspace{1cm} \text{(2.17)}
\]

\[
D_z\bar{\lambda} \equiv \frac{1}{1-\mu\bar{\mu}} \left[ (\bar{\partial} - \mu\bar{\partial} + \frac{1}{2} \partial \mu)\bar{\lambda} - \frac{1}{2} \bar{\alpha}(D_zX) + \frac{i}{2} \bar{\mu} \alpha F \right]. \hspace{1cm} \text{(2.18)}
\]

The fields \((c, \bar{c})\) have been introduced by Becchi \[12\] and are related to the usual diffeomorphism ghosts \((\xi, \bar{\xi})\) by

\[
c = \xi + \mu \bar{\xi}, \hspace{1cm} \text{(2.19)}
\]

\[
\bar{c} = \bar{\xi} + \bar{\mu} \xi. \hspace{1cm} \text{(2.19)}
\]

Their superpartners \((\epsilon, \bar{\epsilon})\) have bosonic statistic and are related in a similar way to the superdiffeomorphism ghosts \((\xi^\theta, \bar{\xi}^\theta)\) by \[10\]

\[
\epsilon = \xi^\theta + \bar{\xi} \alpha, \hspace{1cm} \text{(2.20)}
\]

\[
\bar{\epsilon} = \bar{\xi}^\theta + \xi \bar{\alpha}. \hspace{1cm} \text{(2.20)}
\]
In order to fix the superdiffeomorphism invariance of (2.1), we use a superconformal gauge. For the corresponding gauge fixing term one has

\[-2iS_{gf} = \int \! dzd\bar{z} \left( bs\mu + \bar{b}s\bar{\mu} - b^a s\alpha - \bar{b}^\alpha s\bar{\alpha} \right), \tag{2.21}\]

where \((b, \bar{b})\) are the usual antighost fields, whereas \((b^\alpha, \bar{b}^\alpha)\) are the corresponding supersymmetric partners possessing now a bosonic statistics and

\[sb = s\bar{b} = s b^\alpha = s\bar{b}^\alpha = 0. \tag{2.22}\]

To translate the BRS invariance of the gauge fixed action \((S_{\text{inv}} + S_{gf})\) into a Slavnov identity we introduce a set of invariant external sources \((J^X, J^\lambda, \bar{J}^\lambda, J^F, \bar{J}^c, \bar{J}^e, \bar{J}^e, \bar{J}^\epsilon)\) coupled to the nonlinear BRS variations of the fields, i.e.

\[-2iS_{\text{ext}} = \int \! dzd\bar{z} \left( J^X sX + J^\lambda s\lambda + \bar{J}^\lambda s\bar{\lambda} + J^F sF + J^c s\bar{c} + J^e s\epsilon + \bar{J}^e s\bar{\epsilon} \right). \tag{2.23}\]

The complete action

\[\Sigma = S_{\text{inv}} + S_{gf} + S_{\text{ext}}, \tag{2.24}\]

obeys then to the Slavnov identity

\[S(\Sigma) = 0, \tag{2.25}\]

with

\[S(\Sigma) = \int \! dzd\bar{z} \left( \frac{\delta \Sigma}{\delta J^X} \frac{\delta \Sigma}{\delta \bar{X}} + \frac{\delta \Sigma}{\delta J^\lambda} \frac{\delta \Sigma}{\delta \bar{\lambda}} + \frac{\delta \Sigma}{\delta J^F} \frac{\delta \Sigma}{\delta \bar{F}} + \frac{\delta \Sigma}{\delta J^c} \frac{\delta \Sigma}{\delta \bar{c}} + \frac{\delta \Sigma}{\delta J^e} \frac{\delta \Sigma}{\delta \bar{\epsilon}} \right) \tag{2.26}\]

From the above identity one can read off the linearized Slavnov operator \(\mathcal{B}\)

\[\mathcal{B} = \int \! dzd\bar{z} \left( \frac{\delta \Sigma}{\delta J^X} \frac{\delta \Sigma}{\delta \bar{X}} + \frac{\delta \Sigma}{\delta J^F} \frac{\delta \Sigma}{\delta \bar{F}} + \frac{\delta \Sigma}{\delta J^c} \frac{\delta \Sigma}{\delta \bar{c}} + \frac{\delta \Sigma}{\delta J^e} \frac{\delta \Sigma}{\delta \bar{\epsilon}} \right) \tag{2.27}\]

which, as a consequence of the Slavnov identity, turns out to be nilpotent

\[\mathcal{B} \mathcal{B} = 0. \tag{2.28}\]
As explained in [11, 10] the Beltrami(no) parameters are treated as unquantized external fields. In particular \((\mu, \bar{\mu})\) are the sources for the components \((T_{zz}, T_{\bar{z}\bar{z}})\) of the energy-momentum tensor, i.e.

\[
T_{zz} = \frac{\delta \Sigma}{\delta \mu}, \quad T_{\bar{z}\bar{z}} = \frac{\delta \Sigma}{\delta \bar{\mu}}.
\] (2.29)

In the same way the two Beltramino fields \((\alpha, \bar{\alpha})\) describe the fermionic counterpart of (2.29):

\[
T_{\alpha} = \frac{\delta \Sigma}{\delta \alpha}, \quad T_{\bar{\alpha}} = \frac{\delta \Sigma}{\delta \bar{\alpha}}.
\] (2.30)

One sees then that the classical Slavnov identity (2.25) describes the current algebra of the components of the energy-momentum tensor. In particular, from the expression for the linearized Slavnov operator \(B\) it follows

\[
T_{zz} = B\delta, \quad T_{\bar{z}\bar{z}} = \bar{B}\delta, \quad T_{\alpha} = B\alpha, \quad T_{\bar{\alpha}} = \bar{B}\alpha.
\] (2.31)

meaning that the energy momentum tensor is cohomologically trivial. In addition, it is easily verified that the complete action \(\Sigma\) is itself a \(B\)-variation

\[-2i\Sigma = B\int dzd\bar{z}\left(\frac{1}{2}J^X X - \frac{1}{2}J^\lambda \lambda - \frac{1}{2}J^{\bar{\lambda}} \bar{\lambda} + \frac{1}{2}J^\epsilon \epsilon + J^{\bar{\epsilon}} \bar{\epsilon} + J^c c + J^{\bar{c}} \bar{c}\right).\] (2.32)

Equations (2.31) and (2.32) are the generalization to superstring of the algebraic properties already found in the bosonic string \(\mathbb{F}\). They suggest that, as in the bosonic case, one can interpret the superstring as a topological model of the Witten’s type \(\mathbb{I}\).

### 3 The vector supersymmetry

In order to show the existence of the vector supersymmetry we introduce the two functional operators \(W\) and \(\bar{W}\)

\[
W = \int dzd\bar{z}\left(\frac{\mu}{\delta c} + \frac{\delta}{\delta \alpha} + J^c \frac{\delta}{\delta b} + \bar{\alpha} \frac{\delta}{\delta \bar{\epsilon}} + \bar{J}^c \frac{\delta}{\delta \bar{b}}\right),
\] (3.1)

\[
\bar{W} = \int dzd\bar{z}\left(\frac{\bar{\mu}}{\delta \bar{c}} + \frac{\delta}{\delta \bar{\alpha}} + J^c \frac{\delta}{\delta \bar{b}} + \alpha \frac{\delta}{\delta \epsilon} + J^c \frac{\delta}{\delta b}\right).
\] (3.2)

These operators are a direct extension of the vector supersymmetry operators introduced for the bosonic string \(\mathbb{F}\) and, together with the linearized operator \(B\), give rise to the following algebraic relations

\[
\{\mathcal{B}, W\} = \partial, \quad \{\mathcal{B}, \bar{W}\} = \bar{\partial}, \quad \{W, W\} = \{W, \bar{W}\} = \{\bar{W}, \bar{W}\} = 0.
\] (3.3)

The algebra (3.3) closes on the space-time translations, it this then a supersymmetric algebra of the Wess-Zumino type. Eqs.(3.3) represent a typical feature of
the topological models \([5, 6]\) and, as we will see in the next section, give a simple method for solving the descent equations corresponding to the superdiffeomorphism anomaly.

Let us observe also that, in complete analogy with the bosonic case, the operators \(W\) and \(\bar{W}\) when applied to the classical action \(\Sigma\) yield the linearly broken Ward identities

\[
W \Sigma = \Delta , \quad \bar{W} \Sigma = \bar{\Delta}
\]

(3.4)

with

\[
\Delta = \frac{i}{2} \int dzd\bar{z} \left( J^c \partial c + J^\epsilon \partial \bar{c} - J^\epsilon \partial c - J^\epsilon \partial \bar{c} - J^X \partial X + J^\lambda \partial \lambda \\
+ J^\lambda \partial \bar{\lambda} - J^F \partial F - b \partial \mu - \bar{b} \partial \bar{\mu} + b^\alpha \partial \alpha + \bar{b}^\alpha \partial \bar{\alpha} \right),
\]

(3.5)

\[
\bar{\Delta} = \frac{i}{2} \int dzd\bar{z} \left( J^c \bar{\partial} c + J^\bar{\epsilon} \bar{\partial} \bar{c} - J^\bar{\epsilon} \bar{\partial} c - J^\bar{\epsilon} \bar{\partial} \bar{c} - J^X \bar{\partial} X + J^\bar{\lambda} \bar{\partial} \bar{\lambda} \\
+ J^\bar{\lambda} \bar{\partial} \bar{\lambda} - J^F \bar{\partial} F - b \bar{\partial} \mu - \bar{b} \bar{\partial} \bar{\mu} + b^\alpha \bar{\partial} \alpha + \bar{b}^\alpha \bar{\partial} \bar{\alpha} \right).
\]

(3.6)

Notice that \(\Delta\) and \(\bar{\Delta}\), being linear in the quantum fields, are classical breakings \([13]\).

4 The superdiffeomorphism anomaly

In this section we apply the supersymmetric structure \((3.3)\) in order to solve the BRS consistency conditions for the Slavnov anomaly. In the following we identify the string world sheet with the complex plane \(C\), the result can be generalized to any Riemann surface by introducing an appropriate projective connection, as done in \([16, 17]\).

At the quantum level the classical action \(\Sigma\) gives rise to a one-loop effective action

\[
\Gamma = \Sigma + h \Gamma^{(1)},
\]

(4.1)

which obeys the anomalous Slavnov identity

\[
S(\Gamma) = h \mathcal{A}.
\]

(4.2)

The anomaly \(\mathcal{A}\) is an integrated local two form with ghost number one \([\bar{\mathbb{Q}}]\)

\[
\mathcal{A} = \int \mathcal{A}_2^1
\]

(4.3)

which has to fulfill the BRS consistency condition

\[
\mathcal{B} \mathcal{A} = 0.
\]

(4.4)

\(\bar{\mathbb{Q}}\) We adopt here the usual convention of denoting with \(\mathcal{A}_q^p\) a \(q\)-form with ghost number equal to \(p\).
This condition, when translated to the nonintegrated level, yields a tower of descent equations

\[
\begin{align*}
\mathcal{B}\mathcal{A}_1^2 + d\mathcal{A}_1^2 &= 0, \\
\mathcal{B}\mathcal{A}_2^2 + d\mathcal{A}_0^3 &= 0, \\
\mathcal{B}\mathcal{A}_0^3 &= 0,
\end{align*}
\] (4.5)

where \(d\) denotes the exterior space-time differential

\[
d = dz\partial + dz\bar{\partial}, \quad d^2 = 0, \quad \{\mathcal{B}, d\} = \{\mathcal{W}, d\} = \{\bar{\mathcal{W}}, d\} = 0. \quad (4.6)
\]

As shown in [8], in order to find a solution of the ladder (4.5) it is sufficient to know the nontrivial solution of the last equation of the tower (4.5). It is easy to check indeed that, once a nontrivial solution for \(\mathcal{A}_0^3\) has been found, the cocycles \(\mathcal{A}_1^2\) and \(\mathcal{A}_2^2\) are obtained by successive applications of the operators \(\mathcal{W}\) and \(\bar{\mathcal{W}}\) on \(\mathcal{A}_0^3\), i.e.

\[
\begin{align*}
\mathcal{A}_1^2 &= (\mathcal{W}\mathcal{A}_3^0)dz + (\bar{\mathcal{W}}\mathcal{A}_3^0)d\bar{z}, \\
\mathcal{A}_2^2 &= (\bar{\mathcal{W}}W\mathcal{A}_3^0)dz \wedge d\bar{z}.
\end{align*}
\] (4.7, 4.8)

For what concerns the local cohomology of the linearized operator \(\mathcal{B}\) in the zero form sector with ghost number three it turns out that the relevant cocycle \(\mathcal{A}_0^3\) can be identified, modulo a \(\mathcal{B}\) coboundary, with

\[
\mathcal{A}_0^3 = c\partial c\partial^2 c - c\partial\epsilon\partial\epsilon + \frac{1}{2}\partial c\partial\epsilon\partial\epsilon - c\partial\bar{c}\partial^2\bar{c} + c\partial\epsilon\partial\bar{c} - \frac{1}{2}\partial\bar{c}\partial\bar{c}\epsilon\epsilon. \quad (4.9)
\]

This expression is the supersymmetric extension of the well known term \(\mathcal{A}^3_{0(bos.)}\)

\[
\mathcal{A}^3_{0(bos.)} = c\partial c\partial^2 c - c\partial\bar{c}\partial^2\bar{c}, \quad (4.10)
\]

which is at the origin of the diffeomorphism anomaly for the bosonic string [11, 8, 17].

Applying now the formula (4.8) to \(\mathcal{A}_0^3\), for the integrated Slavnov anomaly \(\mathcal{A}\) one gets the expression

\[
\mathcal{A} = \int d\bar{z}d\bar{\bar{z}}\mathcal{A}_2^1 = 2 \int d\bar{z}d\bar{\bar{z}}(c\partial^3 \mu + \epsilon\partial^2 \alpha) + 2 \int d\bar{z}d\bar{\bar{z}}(c\partial^3 \bar{\mu} + \bar{\epsilon}\bar{\partial}^2 \bar{\alpha}). \quad (4.11)
\]

This is the superdiffeomorphism anomaly (in the Wess-Zumino gauge) of the (1,1) superstring theory [10]. Using the truncation procedure of Sect. 2, for the (1,0) case one has

\[
\mathcal{A}_{(1,0)} = 2 \int d\bar{z}d\bar{\bar{z}}(c\partial^3 \mu + \epsilon\partial^2 \alpha) + 2 \int d\bar{z}d\bar{\bar{z}}(c\partial^3 \bar{\mu}). \quad (4.12)
\]

Expressions (4.11) fixes, through the numerical coefficient of the corresponding Feynman diagrams, the critical dimension of the (1,1) superstring to be equal to 10, while in the (1,0) case Eq. (4.12) leads to a critical dimension 10 in the supersymmetric sector and 26 in the nonsupersymmetric one.
References

[1] E. Witten, *Comm. Math. Phys.* **117** (1988) 353;  
   *Comm. Math. Phys.* **121** (1989) 351;

[2] A. S. Schwarz, *Lett. Math. Phys.* **2** (1978) 247;  
   A. S. Schwarz, Baku International Topological Conference, Abstract, vol. 2, p. 345 (87);

[3] D. Birmingham, M. Blau, M. Rakowski, G. Thompson, *Phys. Rep.* **209** (1991) 129;

[4] F. Delduc, F. Gieres, S. P. Sorella, *Phys. Lett.* B**225** (1989) 367;  
   F. Delduc, C. Lucchesi, O. Piguet, S. P. Sorella, *Nucl. Phys.* B**346** (1990) 313;

[5] E. Guadagnini, N. Maggiore, S. P. Sorella, *Phys. Lett.* B**255** (1991) 65;  
   S. P. Sorella, *Comm. Math. Phys.* **157** (1993) 231;

[6] C. Lucchesi, O. Piguet, S. P. Sorella, *Nucl. Phys.* B**395** (1993) 325;

[7] D. Birmingham, M. Rakowski, *Phys. Lett.* B**269** (1991) 103;  
   *Phys. Lett.* B**272** (1991) 217;  
   *Phys. Lett.* B**275** (1992) 289;  
   *Phys. Lett.* B**289** (1992) 271;

[8] G. Bandelloni and S. Lazzarini, *Diffeomorphism Cohomology in Beltrami parametrization*, PAR-LPTM-93, GEF-TH/93;  
   M. Werneck de Oliveira, M. Schweda and S. P. Sorella, *Phys. Lett.* B**315** (1993) 93;

[9] L. Baulieu, M. Bellon and R. Grimm, *Phys.Lett.* B**198** (1987) 343;  
   *Nucl. Phys.* B**321** (1989) 697;

[10] F. Delduc, F. Gieres, *Class. Quant. Grav.* **7** (1990) 1907;

[11] L. Baulieu, C. Becchi, R. Stora, *Phys. Lett.* B**180** (1986) 55;  
   L. Baulieu, M. Bellon, *Phys. Lett.* B**196** (1987) 142;

[12] C. Becchi, *Nucl. Phys.* B**304** (1988) 513;

[13] A. A. Belavin, A. M. Polyakov, A. B. Zamolodchikov, *Nucl. Phys.* B**241** (1984) 333;

[14] J. Wess, J. Bagger, *Supersymmetry and Supergravity*, Princeton University Press, Princeton, NJ (1983);

[15] O. Piguet, S. P. Sorella, *Helv. Phys. Acta* **63** (1990) 683;
[16] R. Stora, in Non Perturbative Quantum Field Theory, G. 't Hooft and al. eds., Nato ASI series B, vol. 185, Plenum Press (88);
S. Lazzarini and R. Stora, Ward Identities for Lagrangian Conformal Models, in Knots, Topology and Quantum Field Theory, 13th John Hopkins Workshop, L. Lusanna ed., World Scientific (89);

[17] S. Lazzarini, On Bidimensional Lagrangian Conformal Models, Thesis LAPP-Annecy, France, (90) unpublished;
R.Zucchini, Comm. Math. Phys. 152 (1993) 269;