Abstract

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Models on Higher-Ceens Super Riemann Surfaces

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2. The $N = 1$ Superstring

We follow the superspace formulation of the super Riemann surface proposed in Refs.10,11. The chiral $N = 1$ superstring field $X^\mu(Z)$ is a function of the superspace variables $Z = (z, \theta)$. In terms of component fields we get the expansion

$$X^\mu(Z) = x^\mu(z) + \theta \psi^\mu(z). \quad (\mu = 0, 1, 2, \ldots, d - 1)$$

We also introduce a superdilatation field

$$\Phi(Z) = \phi(z) + \theta \psi(z).$$

The chiral stress-energy tensor for these free matter fields $X^\mu$ and $\Phi$ is given by

$$T_m = T_X + T_\Phi,$$

$$T_X = -\frac{1}{2} i D X^\mu \partial X_\mu, \quad (\gamma^0 \gamma^0 = -1),$$

$$T_\Phi = -\frac{1}{2} D \Phi \partial \Phi + \alpha(D + 2\Gamma)(D + \Gamma) D \Phi,$$

where the covariant derivative $D$ is defined by

$$D = \partial / \partial z + \theta \partial / \partial \theta,$$

$\Gamma(Z)$ a superaffine connection and $\alpha$ a real parameter.

The stress-energy tensor for the superghosts $B(Z)$ and $C(Z)$ is given by

$$T_{gh} = -C \partial B + \frac{1}{2} D C \cdot D B - \frac{3}{2} \partial C \cdot B,$$

where $C(Z) = \sigma(z) + \theta \gamma(z)$ is a fermionic ghost with a weight $h = -1$ and $B(Z) = \beta(z) + \theta b(z)$ is a bosonic ghost with a weight $h = 3/2$. The BRST current $Q_0(Z)$ and the superghost number current $J_0(Z)$ are also known to be

$$Q_0 = C(T_m + \frac{1}{2} T_{gh}),$$

$$J_0 = CB.$$
Since $T$ is a real parameter, the minimum value of $\frac{Q}{Q}$ is $1$. Hence, from (3),(2),

$$\left(1 + \frac{Q}{Q}\right) \frac{Q}{Q} = 0$$

and we have from (1) and (2),

$$\left(\frac{Q}{Q} \right) = \frac{Q}{Q} + \frac{Q}{Q}$$

We have the total current change of the contact-resistor system (3),

$$\frac{Q}{Q} = \frac{Q}{Q} + \frac{Q}{Q}$$

Therefore, we need to consider constraints.

Where we need consider:

\begin{align*}
11 \text{ (2):} & \quad \frac{Q}{Q} = \frac{Q}{Q} + \frac{Q}{Q} \\
12 \text{ (2):} & \quad \frac{Q}{Q} = \frac{Q}{Q} + \frac{Q}{Q} \\
13 \text{ (2):} & \quad \frac{Q}{Q} = \frac{Q}{Q} + \frac{Q}{Q} \\
14 \text{ (2):} & \quad \frac{Q}{Q} = \frac{Q}{Q} + \frac{Q}{Q} \\
15 \text{ (2):} & \quad \frac{Q}{Q} = \frac{Q}{Q} + \frac{Q}{Q}
\end{align*}

We refer to the results of our previous work (4.1). The approach is not a full superposition formulation.

\begin{align*}
(11) \quad \frac{Q}{Q} + \frac{Q}{Q} = \frac{Q}{Q} + \frac{Q}{Q} \\
(12) \quad \frac{Q}{Q} + \frac{Q}{Q} = \frac{Q}{Q} + \frac{Q}{Q} \\
(13) \quad \frac{Q}{Q} + \frac{Q}{Q} = \frac{Q}{Q} + \frac{Q}{Q} \\
(14) \quad \frac{Q}{Q} + \frac{Q}{Q} = \frac{Q}{Q} + \frac{Q}{Q}
\end{align*}

The new operators are defined by

\begin{align*}
11 \text{ (2):} & \quad \frac{Q}{Q} = \frac{Q}{Q} + \frac{Q}{Q} \\
12 \text{ (2):} & \quad \frac{Q}{Q} = \frac{Q}{Q} + \frac{Q}{Q}
\end{align*}

By using the contractions (3),(2) \sim (4),(2) we can derive the \( I \) supertopes.
see
\[ \hat{c} = d + 1 \leq 2, \]  
(2.24)

which means that the dimensions of $X^\mu$ should be $d \leq 1$. The unknown functions $\Omega$ and $\Lambda$ are fixed to be
\[ \Omega(Z) = \frac{1}{4} \left[ \hat{c}(R_m - R) - 10(R_{2\hat{A}} - R) \right], \]  
(2.25)
\[ \Lambda(Z) = l(Z, Z) + \Gamma(Z), \]  
(2.26)

where the $R$'s are superprojective connections defined by
\[ R(Z) = \partial \Gamma + \Gamma D \Gamma, \]  
(2.27)
\[ R_m(Z) = 2[\partial_1 D_2 f(Z_1, Z_2)]_{Z_1 = Z_2 = Z}, \]  
(2.28)
\[ R_{2\hat{A}}(Z) = -\frac{1}{5}[2\partial_1 l(Z_1, Z_2) + 3\partial_2 l(Z_1, Z_2) - D_1 D_2 l(Z_1, Z_2)]_{Z_1 = Z_2 = Z}, \]  
(2.29)

and $-l(Z, Z)$ is a superaffine connection. The c-number functions $\Omega$ and $\Lambda$ are tensors with weights $3/2$ and $1/2$, respectively. We shall prove that these $R$'s and $-l(Z, Z)$ are actually subject to transformation rules of superprojective and superaffine connections.

Before we prove the algebra (2.19) with constraints (2.20) to (2.22), we would like to mention a reason why this algebra is topological, and also make a remark about a role of the superaffine connections in the algebra. By using formulas
\[ \frac{1}{2\pi i} \oint_{C_1} dZ_1 Z_{12}^{n-1} = 0, \quad \frac{1}{2\pi i} \oint_{C_2} dZ_1 \partial_1 Z_{12}^{n-1} = \delta_{n,0}, \]  
(2.30)

and from (2.19e) we find the relation
\[ T(Z) = [Q_B, G(Z)], \]  
(2.31)

where $Q_B$ is the super BRST charge defined by
\[ Q_B = \frac{1}{2\pi i} \oint dZ Q(Z). \]  
(2.32)

The relation (2.31) means that the stress-energy tensor $T(Z)$ is BRST-trivial. We also see from (2.19a) or (2.19i) that the total central charge $\hat{c}_{\text{tot}}$ is zero. The existence of the nilpotent BRST charge $Q_B$ and the BRST-trivial stress-energy tensor $T(Z)$ is the characteristic feature of topological theories.

A role of the affine connection $\Gamma$ is as follows: The affine connection guarantees coordinate-independence of the algebra. In order to see this more explicitly, we define the super Virasoro operator
\[ L_V = \frac{1}{2\pi i} \oint dZ V(Z) T(Z) \]  
(2.33)

and
\[ J_A = \frac{1}{2\pi i} \oint dZ A(Z) J(Z), \]  
(2.34)

where $V(Z)$ is a supervector field with weight $h = -1$ and $A(Z)$ is a scalar field. From (2.19d) we calculate their commutator to give
\[ [L_V, J_A] = J_{(V,A)} - \frac{\gamma^2}{2} \chi(V, A), \]  
(2.35)

where
\[ J_{(V,A)} = \frac{1}{2\pi i} \oint dZ (-\partial A \cdot V + \frac{1}{2} DA \cdot DV) J, \]  
(2.36)
\[ \chi(V, A) = \frac{1}{2\pi i} \oint dZ A D(D - \Gamma)(D - 2\Gamma)V. \]  
(2.37)

Comparing (2.19d) with (2.35) we see a correspondence
\[ R_{12}^{(13)} = \left[ -\frac{1}{Z_{12}} \frac{\partial_{12}}{Z_{12}} \Gamma_2 - \frac{1}{Z_{12}} \frac{\partial_{12}}{Z_{12}} \Gamma_2 - \frac{1}{Z_{12}} D \Gamma_2 \right] \longrightarrow D(D - \Gamma)(D - 2\Gamma). \]  
(2.38)

Namely, the affine connection $\Gamma$ appears in the super-covariant derivative $(D + 2h \Gamma)$ in the anomaly $\chi(V, A)$. Since, for $\phi_h$ with weight $h$, $(D + 2h \Gamma) \phi_h$ has a weight.
The notation but the condition that the local curl equation is given by

\[(A) \frac{\partial \phi}{\partial x} + \frac{\partial \phi}{\partial y} = 0 \]


is used when the following relation holds

\[(B) \frac{\partial \phi}{\partial x} + \frac{\partial \phi}{\partial y} = \nabla \times \mathbf{A} = 0 \]

where \( \mathbf{A} \) is the vector field.

\[\phi \quad \text{or} \quad \nabla \times \mathbf{A} = 0 \]

In the same manner as we did in the derivation of (9.19a) we calculate the

\[\text{function of } \left( \frac{\partial \phi}{\partial x} \right) \]

where we have used the relation (9.19c) and (9.19d). The notation \( \phi \) for general situation is defined by (9.19a). Similar calculation of (9.19b) gives the result (9.19c).

\[\phi = \frac{\partial \phi}{\partial x} + \frac{\partial \phi}{\partial y} \]

Similarly, the component of (9.19) are coordinate-independent. To be a scalar, we have

\[\phi \quad \text{or} \quad \nabla \times \mathbf{A} = 0 \]

We have used the notation \( \phi \) for general situation is defined by (9.19a). Similar calculation of (9.19b) gives the result (9.19c).
we have
\[
\frac{\partial_{12}}{2Z_{12}^2} [(l(2, 2) - \gamma^2 \Gamma_2 - \Lambda_2 - 2 \alpha \beta \Gamma_2) C_2 = \frac{\partial_{12}}{2Z_{12}^2} [l(2, 2) + \Gamma_2 - \Lambda_2] C_2, \tag{2.51}
\]
\[
\frac{1}{2Z_{12}} D_1 [l(1, 2) + l(2, 1) - \gamma^2 \Gamma_2 - \Lambda_2 - 2 \alpha \beta \Gamma_2] z_l = z_l C_2 \tag{2.52}
\]
\[
\frac{\partial_{12}}{Z_{12}} \{ \partial_1 [l(1, 2) + l(2, 1) - 2 \alpha \beta \Gamma_1] z_l = z_l C_2 \\
+ \frac{1}{2} l((2, 2) - \alpha \beta \Gamma_2) \partial_2 C_2 \\
+ \frac{1}{2} D_1 [l(1, 2) + l(2, 1) - 2 \alpha \beta \Gamma_1] z_l = z_l DC_2 \} + \cdots.
\]

Taking account of the constraints (2.21) and (2.22), we find
\[
\frac{1}{2} + \alpha \beta - \frac{1}{2} \beta \gamma = 0, \tag{2.44}
\]
\[
\frac{1}{2} - \gamma - \alpha \beta - \beta \gamma = 0. \tag{2.45}
\]

As a result we see that the second, third and fourth terms on the r.h.s. of (2.43) identically vanish, i.e.,
\[
\frac{1}{2} \partial_1 \left[ \left( \frac{1}{2} + \alpha \beta - \frac{1}{2} \beta \gamma \right) \partial_2 C_2 + \frac{1}{2} (\beta + \gamma) C_2 \partial \Phi_2 \right] = 0, \tag{2.46}
\]
\[
\frac{1}{2} \partial_1 \left[ \left( \frac{1}{2} - \gamma - \alpha \beta - \beta \gamma \right) C_2 = 0, \tag{2.47}
\right.
\]
\[
\frac{\partial_{12}}{Z_{12}} \left[ \left( \frac{1}{2} + \alpha \beta - \frac{1}{2} \beta \gamma \right) DC_2 - \frac{1}{2} (\beta + \gamma) C_2 D \Phi_2 \right] = 0. \tag{2.48}
\]

Substituting from (2.22) into the last three terms on the r.h.s. of (2.43) and noting
\[
D_1 [l(1, 2) + l(2, 1)] z_l = z_l = D_2 l(2, 2), \tag{2.49}
\]
\[
\partial_1 [l(1, 2) + l(2, 1)] z_l = z_l = \partial_2 l(2, 2), \tag{2.50}
\]
Let us fix the unknown function A by (2.26). Then Eqs. (2.51) and (2.52) vanish, and Eq. (2.53) becomes
\[
\frac{\partial_{12}}{Z_{12}} \{ \partial_1 [l((2, 2) + \Gamma_2 + \gamma^2 \Gamma_2 + \Lambda] \partial_2 C_2 \\
+ \frac{1}{2} l((2, 2) + \Gamma_2 + \gamma^2 \Gamma_2) \partial_2 C_2 \\
+ \frac{1}{2} D_1 [l((2, 2) + \Gamma_2 + \gamma^2 \Gamma_2) DC_2] \} = \frac{\partial_{12}}{Z_{12}} \{ \partial_1 [l((2, 2) + \Gamma_2 + \gamma^2 \Gamma_2) DC_2] \} + \cdots. \tag{2.54}
\]
Collecting terms proportional to \( \theta_{12} / Z_{12} \) one finally finds Eq. (2.43) to be
\[
- \frac{\partial_{12}}{Z_{12}} Q_2 + \cdots. \tag{2.55}
\]
This completes the proof of (2.19f).

Other remaining equations in (2.19) can be derived in a similar way by carrying out straightforward but long calculations.
Theorem 4.1. This shows that $\zeta^2$ is not the super Schouten determinant. Here it should be remarked that the super Schouten determinant is given by

$$\zeta = \frac{\theta_{\gamma}(\phi_{\gamma})}{\theta_{\gamma}(\phi_{\gamma})},$$

where $\zeta$. From this we take a limit $Z \to Z'$. On the other hand, the last term of (2.69) does not vanish.

$\zeta$ is a transformation rule for $\delta$ and $\gamma$.

This theorem is known to give the super Schouten determinant. Here it should be remarked that the super Schouten determinant is given by

$$\zeta = \frac{\theta_{\gamma}(\phi_{\gamma})}{\theta_{\gamma}(\phi_{\gamma})},$$

where $\zeta$. From this we take a limit $Z \to Z'$. On the other hand, the last term of (2.69) does not vanish.

We would like to prove the following:

$$\zeta^2 = \frac{\theta_{\gamma}(\phi_{\gamma})}{\theta_{\gamma}(\phi_{\gamma})} + \frac{\theta_{\gamma}(\phi_{\gamma})}{\theta_{\gamma}(\phi_{\gamma})} \zeta^2.$$
3. The \( N = 2 \) Superstring

The chiral \( N = 2 \) superstring field \( X^\mu(Z) \) is a function of the superspace variable \( Z = (z, \theta, \bar{\theta})^{10,11} \). In terms of component fields we get the expansion

\[
X^\mu(Z) = x^\mu(z) + \partial \bar{\psi}^\mu(z) + \partial \psi^\mu(z) + \frac{1}{2} \partial \theta \partial \bar{\theta}^\mu(z), \quad (\mu = 0, 1, 2, \ldots, d - 1). \quad (3.1)
\]

Let us introduce an infinitesimal superdistance \( dZ \) by

\[
dZ = dz + \frac{1}{2} (d\theta d\bar{\theta} + \theta d\bar{\theta}). \quad (3.2)
\]

A super analytic transformation \( Z \rightarrow \tilde{Z} \) transforms \( dZ \) according to

\[
d\tilde{Z} = (D\bar{\theta})(D\theta) dZ, \quad (3.3)
\]

where \( D \) and \( \bar{D} \) are covariant derivatives defined by

\[
D = \frac{\partial}{\partial \theta} + \frac{1}{2} \theta \frac{\partial}{\partial z}, \quad \bar{D} = \frac{\partial}{\partial \bar{\theta}} + \frac{1}{2} \bar{\theta} \frac{\partial}{\partial z} \quad (3.4)
\]

with properties

\[
D^2 = D^2 = 0, \quad (3.5)
\]

\[
\{D, \bar{D}\} = \frac{\partial}{\partial z}, \quad (3.6)
\]

\[
D = (D\theta)\bar{D}, \quad \bar{D} = (D\bar{\theta})\bar{D}. \quad (3.7)
\]

We also introduce a super line element squared \( ds^2 = |edZ|^2 \), where \( e \) is a super-einbein. We require that the one-form \( edZ \) is invariant under the transformation \( Z \rightarrow \tilde{Z} \). This yields a transformation rule of \( e \)

\[
e = (D\bar{\theta})(D\theta)e. \quad (3.8)
\]

From this we find that superaffine connections \( \Gamma = De/e \) and \( \bar{\Gamma} = \bar{D}e/e \) are subject to transformation rules

\[
\Gamma = (D\theta)\bar{\Gamma} + \frac{\partial \bar{\theta}}{D\theta}, \quad \bar{\Gamma} = (D\bar{\theta})\Gamma + \frac{\partial \theta}{D\bar{\theta}}. \quad (3.9)
\]

The primary superconformal field \( \Phi_{pf} \) is characterized by a pair of weights \( (p, q) \) and by its parity \( \sigma \) with respect to interchange \( \theta \leftrightarrow \bar{\theta} \). We denote this characteristic by \( (p, q)^\sigma \). Here we give a list of the assignment of weights and parity of primary superconformal fields with which we will concern in what follows.

<table>
<thead>
<tr>
<th>Fields</th>
<th>((p, q)^\sigma)</th>
</tr>
</thead>
<tbody>
<tr>
<td>Superstring</td>
<td>(X^\mu)</td>
</tr>
<tr>
<td>Superdilaton</td>
<td>(\Phi)</td>
</tr>
<tr>
<td>Superantighost</td>
<td>(B)</td>
</tr>
</tbody>
</table>
| Superghost   | \(C\)             | \((-1,-1)^+\)     | (3.10)

As for the \( \sigma \)-parity for \( B \) and \( C \), we consider the super BRST charge

\[
Q_B = \frac{1}{2\pi i} \oint dz d\theta d\bar{\theta} Q(Z). \quad (3.11)
\]

Since the volume element \( d\theta d\bar{\theta} \) changes its sign under an interchange \( \theta \leftrightarrow \bar{\theta} \), the BRST current \( Q(Z) \), which will be explicitly defined by (3.24) below, should have a negative \( \sigma \)-parity. Based on this observation, it is possible to find proper assignments of the \( \sigma \)-parity as in table (3.10). Namely, we have assigned the \( \sigma \)-parity of \( B \) to be negative and that of \( C \) to be positive.

The transformation rule of \( \Phi_{pf} \) is given by

\[
\Phi_{pf}(Z) = (D\bar{\theta})^p(D\theta)^q \Phi_{pf}(\tilde{Z}). \quad (3.12)
\]

Two kinds of supercovariant derivatives \( D - p\Gamma \) and \( \bar{D} - q\bar{\Gamma} \) are necessary to give the following transformation rules:

\[
(D - p\Gamma)\Phi_{pf} = (D\bar{\theta})^p(D\theta)^{q+1}(\bar{D} - q\bar{\Gamma})\Phi_{pf}, \quad (3.13)
\]

\[
(\bar{D} - p\bar{\Gamma})\Phi_{pf} = (D\bar{\theta})^{p+1}(D\theta)^q(\bar{D} - q\bar{\Gamma})\Phi_{pf}. \quad (3.14)
\]

Fundamental contractions on an arbitrary higher-genus super Riemann surface \( M_2 \) are given by

\[
<X^\mu(Z_1)X^\nu(Z_2)> = \eta^\mu\nu[-2\ln Z_{12} + f(Z_1, Z_2)]. \quad (3.15)
\]
\[
\frac{\partial q}{\partial e} + \frac{\partial q}{\partial e} = (z, z)_S
\]

with the super Schrödinger derivative \( (z, z)_S \)

\[
(z, z)_S + (z, z)_S(q \partial)(q \partial) = (z, z)_R
\]

which is subject to a transformation rule.

\[
(z, z)_R = \zeta + \zeta - \zeta \partial + \partial \zeta = \zeta
\]

is an \( q = N \) superfield configuration.

where \( q = \xi \) is the dimension of matter fields for \( p = 0 \) and \( q = \eta \) and \( \xi \) (the index of matter fields for \( p = 1 \) and \( \eta \) for \( p = 2 \)).

\[
(z, z)_R = \zeta (\zeta - \xi_\xi + \xi_\eta) - \xi_\xi (\zeta - \xi_\xi + \xi_\eta) = \zeta_r
\]

is a number tensor of the type \((1, 1)\) which is explicitly given by

\[
(z, z)_R = \zeta \left( 1 - \xi_\xi + \xi_\eta \right)
\]

For the case \( \xi = \eta = 0 \), we have the fundamental countermeasure

\[
\xi_\xi + \xi_\eta = \xi_\xi
\]

Note that the above four generators have all negative \( \xi \) -parity. A similar algebra is valid for the superfield of matter field is the dimension of matter fields for \( p = 1 \) and \( \eta \) for \( p = 2 \).

We then consider the following four generators: which are arbitrary basis for \( \zeta \) and \( q \) and are holomorphic except for \( \zeta \) and \( q \).

\[
\xi_\xi + \xi_\eta = \xi_\xi
\]

where \( \xi_\xi + \xi_\eta = \xi_\xi \) and \( \xi_\xi + \xi_\eta = \xi_\xi \) for \( \zeta \).
with a constraint

\[ \hat{c} = 2(1 - \alpha^2). \]  

(3.34)

Here we have set, for brevity,

\[ \delta = \tilde{\theta}_{12}/Z_{12} \]  

(3.35)

\[ \Delta_0 = \tilde{\theta}_{12}/Z_{12} \]  

(3.36)

\[ \Delta_1 = (\tilde{\theta}_{12}/Z_{12})D_2 - (\theta_{12}/Z_{12})D_2 \]  

(3.37)

\[ \Delta_2 = (\tilde{\theta}_{12}/Z_{12})\partial_2 \]  

(3.38)

and by \( A \sim B \) we mean that \( A \) is equal to \( B \) except for finite terms.

Contrary to the \( N = 1 \) algebra, there appear no superdilaton terms in \( J \) and \( Q^{(1)} \), and, moreover, the parameter \( \alpha \) can be chosen to be zero. This means that the \( N = 2 \) algebra is valid without the superdilaton field. In this case the dimensions of \( X^\alpha \) is just \( \hat{c} = d = 2 \). This coincides with the result obtained by Gomis and Suzuki in genus-zero case. When we have the superdilaton in \( T \), the constraint (3.34) implies \( \hat{c} = d + 1 \leq 2 \), namely, \( \hat{c} = 1 \) or \( \hat{c} = 2 \). The \( \hat{c} = 1 \) case is realized by the superdilaton only without the superstring \( X^\alpha \), and the \( \hat{c} = 2 \) case is realized by the superstring in dimension one plus the superdilaton.

From (3.33e) we obtain the formula

\[ T(Z) = [Q_B, G(Z)], \]  

(3.39)

where \( Q_B \) is the super BRST charge, namely, the stress-energy tensor \( T \) is BRST-trivial. We also see from (3.33a) and (3.33i) that the total central charge is zero. These features are characteristic of topological theories.

4. Concluding remarks

On arbitrary higher-genus super Riemann surfaces we have shown that the \( N = 1 \) and \( N = 2 \) superstrings realize supertopological algebras which are superfield extensions of the twisted \( N = 2 \) superconformal algebra. In the \( N = 1 \) case the superdilaton field is necessary in order to validate the algebra, whereas in the \( N = 2 \) case the superdilaton is not always necessary for the algebra. The \( N = 1 \) supertopological algebra is realized by the superstring in dimension \( d = 1 \) plus the free superdilaton field or by the superdilaton only without the superstring. This conclusion is different from that of other similar works.\(^7\) Fujikawa and Suzuki\(^7\) obtained the result of \( d = \pm \infty \), where there is no \( D(D - \Gamma)(D - 2\Gamma)C \) term in \( Q(Z) \). Bershadsky et al.\(^9\) discussed that any superstring theory in any dimension has the \( N = 3 \) twisted supersymmetry. However, their approach is not the full superspace formulation. Their conclusion about dimensionality has been derived by introducing the Coulomb gas model of the matter fields. On the other hand, we have not introduced such a model.

The \( N = 2 \) supertopological algebra is also realized by three models: (i) the superstring in two dimensions, (ii) the superstring in one dimension plus the superdilaton, (iii) the superdilaton only without the superstring.

The conclusion drawn from (i) is the same as the one obtained by Gomis and Suzuki\(^9\) which corresponds to ours in genus-zero case. On the other hand, the other two models yield different results from ours, since the \( \sigma \) parity of \( T_m \) has not been considered seriously.

From the mathematical point of view we have found some new results: (i) the \( N = 2 \) superprojective connection (3.30) expressed by the superaffine connection \( \Gamma \), (ii) the transformation rules (3.13) and (3.14) for supercovariant derivatives \( D - p\Gamma \) and \( \bar{D} - q\Gamma \) in the \( N = 2 \) case.
References