Interactions of Discrete States with Nonzero Ghost Number
in $c = 1$ 2D Gravity

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Abstract

We study the interactions of the discrete states with nonzero ghost number in $c = 1$ two-dimensional (2D) quantum gravity. By using the vertex operator representations, it is shown that their interactions are given by the structure constants of the group of the area preserving diffeomorphism similar to those of vanishing ghost number. The effective action for these states is also worked out. The result suggests the whole system has a BRST-like symmetry.
Much attention has recently been paid to nonperturbative treatment of two-dimensional (2D) quantum gravity in terms of the matrix model [1, 2]. Most of the remarkable results there are now understood in the continuum approach using the techniques in the conformal field theory [3, 4]. There still remain many problems to be clarified, however, in order to understand the full theory. The problem with which we are mainly concerned in this paper is the interaction of the so-called “discrete states” in the $c = 1$ conformal field theory coupled to 2D gravity [5-9].

In the conformal gauge, the $c = 1$ quantum gravity may be regarded effectively as a string theory in two dimensions with suitable background charge. Thus it is expected that there is only the degree of freedom corresponding to the “center of mass” or “tachyon”, since there are no transverse directions. However, it has been found that there exist other discrete degrees of freedom in the $c \leq 1$ theory coupled to the 2D quantum gravity, both in the matrix models and in the Liouville approach [5-9].

The nature of these “discrete states” is most effectively studied in the Liouville theory. By using the BRST formulation of the $c = 1$ quantum gravity, all the physical states characterized by the BRST cohomology have been enumerated and it has been found that there are indeed an infinite number of physical states with ghost number $N_{FP} = 0, \pm 1$ at the discrete values of momenta [6-11]. These states may be interpreted as higher string states [6], but they exist only for fixed values of momenta, allowing for no usual particle interpretation.

Recently it has been pointed out that the dynamics of these discrete states are governed by the symmetry group of the area preserving diffeomorphism [10-15]. For $c = 1$, these states with $N_{FP} = 0$ are known to form representations of the $SU(2)$ Kac-Moody algebra [10-13,15,16]. By using the vertex operator representation of these states by means of the $SU(2)$ algebra, Klebanov and Polyakov have computed the three point interactions and have proposed an effective action for these discrete states [11]. This action is further made complete by including the scalar degrees of freedom [15]. However, the interactions involving those extra states with $N_{FP} \neq 0$ have not been studied and their role in the theory remains elusive. It is thus interesting to examine their interactions and
try to clarify their role in order to understand the whole structure of the theory.

The purpose of this paper is to fill the above gap and, in particular, determine the three point interactions of the extra states with \( N_{FP} \neq 0 \). We will do this using the vertex operator representation of these states given for the \( c = 1 \) theory [13].

Denoting the \( c = 1 \) matter and Liouville fields by \( X(z) \) and \( \phi(z) \), respectively, one has the energy-momentum tensor

\[
T(z) = -\frac{1}{2} (\partial X)^2 - \frac{1}{2} (\partial \phi)^2 + \sqrt{2} \partial^2 \phi. \tag{1}
\]

The discrete physical state are characterized by the \( SU(2) \) current algebra defined by

\[
J^\pm(z) = e^{\pm i\sqrt{2}X(z)}, \quad J^0(z) = \frac{1}{\sqrt{2}} i \partial X(z). \tag{2}
\]

The states with \( N_{FP} = 0 \) are obtained by acting the following operators on the physical vacuum \( \equiv |0 > X, \phi \otimes c_1 |0 >_{bc} \)

\[
\tilde{\Psi}^{(\pm)}_{Jm}(z) = \left( \frac{(J + m)!}{(2J)!(J - m)!} \right) (J^-)^{-m} e^{i\sqrt{2}JX(z)+\sqrt{2}(1\pm J)\phi(z)}, \tag{3}
\]

where \( J^- \) is the contour integral over the current \( J^- \) around \( z \) and \( J = 1, 2, \cdots; m = -J, -J + 1, \cdots, J - 1, J \). These generate spin \( J \) multiplet with \( N_{FP} = 0 \). Their interactions are now known to be described by the symmetry of the area preserving diffeomorphism [11, 15].

The extra states with \( N_{FP} \neq 0 \) also fall into representations of the \( SU(2) \) current algebra. It has been shown that these states are generated by the following operators [13]

\[
\tilde{\Psi}^{(-)}_{J-1,m}(z) \sim (J^-)^{-m-1} \int_z \frac{dw \, c(w)J^-(w)}{2\pi i (w - z)} e^{i\sqrt{2}JX(z)+\sqrt{2}(1+J)\phi(z)}, \tag{4}
\]

\[
\tilde{\Psi}^{(+)}_{J-1,m}(z) \sim (J^-)^{-m-1} \int_z \frac{dw \, b(w)e^{-iX(w)/\sqrt{2}}-\phi(w)/\sqrt{2}e^{i\sqrt{2}JX(z)+\sqrt{2}(3/2-J)\phi(z)},} {2\pi i (w - z)} \tag{5}
\]

where the caret on the ghost field \( c(w) \) in eq. (4) means that the zero mode \( c_0 \) is removed. As has been noted in ref. [13], this is necessary to make the created states in the "relative cohomology" and these states form spin \( (J - 1) \) representation.

*For simplicity, we consider integer spin \( J \) in this paper.

1In other words, we subtract the states proportional to \( \partial c(z) \) in the "absolute cohomology" which are in spin \( J \) representation.
Applying the operator (4) on the physical vacuum, we easily see that these create states with \( N_{FP} = 1 \)
\[
\tilde{\Psi}_{J-1,m}(0)|0 > \otimes c_1|0 >_{bc}
\sim (J_0^-)^{J-m-1} \int_0^{2\pi} \sum_{n \neq 0} c_n w^{-n} e^{-i\sqrt{2}X(w)}|p_X = \sqrt{2}(J-1), p_\phi = -\sqrt{2}iJ >, \\
= (J_0^-)^{J-m-1} \sum_{n \geq 1} c_n S_{2J-1-n} \left( -\frac{\sqrt{2}X(w)}{n} \right) |p_X = \sqrt{2}(J-1), p_\phi = -\sqrt{2}iJ >, 
\]
in agreement with ref. [8]. Here we have used the Schur polynomial defined by
\[
\exp \left( \sum_{k \geq 0} x_k z^k \right) = \sum_{k \geq 0} S_k(x) z^k. 
\]
We can show that other states created by (3)-(5) also agree with those given in ref. [8] and that these states are BRST invariant.

In order to examine the interactions of these states, we first follow ref. [11] and compute the operator product expansions (OPEs). From the ghost number conservation and the dependence on the zero modes of \( X \) and \( \phi \), we should have
\[
\tilde{\Psi}_{J_1-1,m_1}^{(+)}(z) \tilde{\Psi}_{J_2-1,m_2}^{(-)}(0) = \cdots + \frac{1}{z} \sum F_{J_1-1,m_1,J_1+J_2-2,-m_1-m_2}^{J_2,-m_2} \Psi_{J_2-2,m_2}^{(-)}(0) + \cdots, \\
F_{J_1-1,m_1,J_1+J_2-2,-m_1-m_2}^{J_2,-m_2} = C_{J_1-1,m_1,J_1+J_2-2,-m_1-m_2}^{J_2,-m_2} g(J_1,J_2), 
\]
where \( C \) are the Clebsch-Gordan coefficients and \( g(J_1,J_2) \) is an unknown function to be determined. For \( J_3 = J_1 + J_2 - 1, m_3 = -m_1 - m_2 \), we have
\[
C_{J_1-1,m_1,J_3-1,m_3}^{J_2,-m_2} = \frac{(-1)^{J_1-1-m_1} N(J_3-1,m_3)}{N(J_1-1,m_1)N(J_2,m_2)} [m_1 J_2 - m_2 (J_1 - 1)], 
\]
\[
N(J,m) = \sqrt{\frac{(J-m)!(J+m)!}{(2J-1)!}}. 
\]
Notice that \( \tilde{\Psi}_{J-1,m}^{(\pm)} \) have spin \((J-1)\). We thus see that the last factor in eq. (9a) is the structure constant of the area preserving diffeomorphism [17].

Our next task is to determine \( g(J_1,J_2) \). For this purpose, we consider the special case \( m_1 = -J_1 + 2 \) and \( m_2 = -J_2 \). One finds
\[
\sqrt{2(J_1 - 1)\Psi^{(+)}_{J_1-1,-J_1+2}(z)\tilde{\Psi}^{(-)}_{J_1+J_2-2, J_1+J_2-2}(0)}
= \frac{1}{z} \int \frac{dw}{2\pi i} \int \frac{dv}{2\pi i} \int \frac{du}{2\pi i} (w-v)(w-1)^{-2J_1+1}(w-u)^{-2VJ_1+2J_2-2}(v-u)^{-2} \\
\times (v-1)^{-2}J_1V^{2J_1+2J_2-3}(1-u)^{-2J_1-1}u^{-2J_1-2}J_2+3 \tilde{\Psi}^{(-)}_{J_2,J_2}(0),
\]

where we have used the fact that the contraction \(<b(z)c(w)> \sim w^2/z^2(z-w)\) because the zero mode \(c_0\) is removed from \(c(w)\). The \(u\)-integration is then deformed to be winding around \(u = v\) and \(w\). We thus get

\[
-\frac{1}{z} \int \frac{dw}{2\pi i} \int \frac{dv}{2\pi i} \int \frac{du}{2\pi i} [(1-w)^{2J_1-1}w^{2J_1-2J_2+3} - (1-v)^{2J_1-1}v^{2J_1-2J_2+3}]
\times \frac{\partial^2}{\partial u \partial v} [(w-v)(w-1)^{-2d_1+1}w^{2J_1+2J_2-2}(v-1)^{-2J_1+1}v^{2J_1+2J_2-3} \tilde{\Psi}^{(-)}_{J_2,J_2}(0)]
\]

\[
= -\frac{(2J_1+2J_2-3)!}{(2J_1-3)!(2J_2-1)!} \tilde{\Psi}^{(-)}_{J_2,J_2}(0).
\]

Combining eq. (11) with (8) and (9), we find

\[
F = \frac{(-1)^{J_1-1-m_1}N(J_3-1, m_2)}{N(J_1-1, m_1)N(J_2, m_2)} \frac{(2J_3-1)!}{(2J_1-3)!(2J_2-1)! \sqrt{2J_2(J_1-1)(J_3-1)}} [m_1 J_2 - m_2 (J_1-1)].
\]

In principle, it should be possible to compute other OPEs, but we have not been able to complete it. Here, instead of pursuing this line, let us directly compute the three point functions involving these states. This is equivalent to computing all possible OPEs at once. This also serves as a check of our above result.

It is easy to see that the only nonvanishing function is

\[
<0|\Psi^{(+)}_{J_2,m_2}(z_1)c(z_1)\tilde{\Psi}^{(+)}_{J_1-1,m_1}(z_2)c(z_2)\tilde{\Psi}^{(-)}_{J_1+J_2-2,-m_1-m_2}(0)c(0)|0>,
\]

which will be a constant that has the same structure as the coefficient \(F\) in eq. (8). There is no other coupling of those states with “tachyonic” states.
To compute this, we again specialize to the case $m_1 = -J_1 + 1$ and $m_2 = -J_2 + 1$. First we compute (13) with zero mode $c_0$ included and then subtract the contribution from $c_0$. The first contribution (times $\sqrt{2J_2}$) is given by

$$\int_{z_1} \int_{z_2} \int_{2\pi i} \int_{2\pi i} \int_{0} \int_{2\pi i} (z_1 - z_2)^{2J_1 + 2J_2} z_1^{-2J_1 + 1} z_2^{-2J_2} (w - z_1)^{-2J_2} (z_1 - v)^{-2J_2}$$

$$\times (z_1 - u)^{2J_1 + 1} (w - v)(w - z_2)^{-2J_1 + 1} (w - u)^{-2J_1 + 2} (v - z_2)^{-2J_1 + 2} \times (v - u)^{-2J_1 + 2} (z_2 - u)^{2J_1} u^{2J_1 - 2J_2 + 2}.$$ 

(14)

At first sight, this appears to be $SL_2$ non-invariant and hence depends on $z_1$ and $z_2$. We will see that actually this gives a constant. 

The contour integration over $u$ is again deformed and one finds only the contribution form $u = \infty$. The result turns out to be

$$- \oint_{J_1} \oint_{J_2} \oint_{2\pi i} \oint_{2\pi i} (z_1 - z_2)^{2J_1 + 2J_2} z_1^{-2J_1 + 1} z_2^{-2J_2} (w - z_1)^{-2J_2} (w - z_2)^{-2J_1 + 1}$$

$$\times u^{2J_1 + 2J_2 - 2} (z_1 - v)^{-2J_2} (v - z_2)^{-2J_1 + 1} u^{2J_1 + 2J_2 - 2}.$$ 

(15)

We then perform the $w$-integration, which is deformed to be around $w = z_2$ and $\infty$. The contribution from $w = z_2$ drops out because of symmetry and we are left with

$$- \oint_{J_2} \oint_{2\pi i} (2J_2 z + (2J_1 - 1) z_2 - v)(z_1 - v)^{-2J_2} (v - z_2)^{-2J_1 + 1} u^{2J_1 + 2J_2 - 2}.$$ 

(16)

Performing the $v$ integration, we finally get

$$- \frac{(2J_1 + 2J_2 - 2)!}{(2J_1 - 2)!(2J_2 - 1)!},$$

(17)

which is independent of $z_1$ and $z_2$.

The contribution from the zero-mode term, on the other hand, can be computed similarly. Using the correlator of ghosts

$$< 0|c(z_1)b(v)c(z_2)c_0uc(0)|0 >= \frac{u z_1^2 z_2^2 (z_1 - z_2)}{v^2 (z_1 - v)(v - z_2)},$$

(18)

\footnote{This can be understood from the fact that (14) is made $SL_2$ invariant by multiplying with $\lim_{Z \to \infty} \frac{Z - u}{Z} = 1$.}
we find it is given by
\[
\oint \frac{dw}{2\pi i} \oint \frac{dv}{2\pi i} \oint \frac{du}{2\pi i} (z_1 - z_2)^{2J_1 + J_2 - 2} \frac{z_1^{-2} J_1 + 2 z_2^{-2} J_2 + 1}{z_2^{-2}} \frac{(w - z_1)^{-2} J_2 (z_1 - v)^{-2} J_2}{(w - z_2)^{-2} J_2 (w - u)^{-2} J_2 + 2 (v - z_2)^{-2} J_1 + 1 (v - u)^{-1} \times v^{2J_1 + 2J_2 - 3} (z_2 - w)^{2J_1 - 1} u^{-2J_1 - 2J_2 + 2}, \tag{19}
\]

which is \( SL_2 \) invariant. By a similar procedure, one finds
\[
-2J_2 \frac{(2J_1 + 2J_2 - 3)!}{(2J_1 - 2)!(2J_2 - 1)!} \tag{20}
\]

Hence the three point correlation (12) for \( m_1 = -J_1 + 1 \) and \( m_2 = -J_2 + 1 \), which is given by the difference of (17) and (20), is
\[
-\frac{(2J_1 + 2J_2 - 3)!}{(2J_1 - 3)!(2J_2 - 1)!} \tag{21}
\]

Comparing this with the general structure of the correlation, we find it is given precisely by (12).

If we redefine the fields by
\[
\Psi^{(\pm)}_{J,m}(z) = -\left[ N(J, m)(2J - 1)! \sqrt{\frac{J}{2}} \right]^{\pm 1} \Psi^{(\pm)}_{J,m}(z),
\]
\[
\tilde{\Psi}^{(\pm)}_{J-1,m}(z) = (-1)^{J-m} N(J - 1, m)(2J - 3)! \sqrt{\frac{J - 1}{2}} \tilde{\Psi}^{(\pm)}_{J-1,m}(z),
\]
\[
\tilde{\Psi}^{(-)}_{J-1,m}(z) = \left[ N(J - 1, m)(2J - 3)! (2J - 1) \sqrt{\frac{J - 1}{2}} \right]^{-1} \tilde{\Psi}^{(-)}_{J-1,m}(z), \tag{22}
\]

we get
\[
< 0 | \Psi^{(s)}_{J_2,m_2}(z_1) c(z_1) \tilde{\Psi}^{(+)}_{J_1-1,m_1}(z_2) c(z_2) \tilde{\Psi}^{(-)}_{J_1+J_2-2,-m_1-m_2}(0) c(0) | 0 > = m_1 J_2 - m_2 (J_2 - 1), \tag{23}
\]

Therefore, introducing variables \( g^{(s),A}_{J,m} \) and \( \tilde{g}^{(s),A}_{J-1,m} (s = \pm) \) for these states, the effective action for the three-point interactions is given by
\[
S_{3,gh} = -g_0 \sum_{J_1,m_1,J_2,m_2,A,B,C} [J_2 m_1 - (J_1 - 1) m_2] f^{ABC}_{J-1, J_2, J_2, -m_1 - m_2} h^{(-),A}_{J_1, J_2} \tilde{g}^{(+),B}_{J_1-1, m_1} g^{(+),C}_{J_2, m_2} \int d\phi, \tag{24}
\]
where we have introduced the open string coupling constant $g_0$ and the Chan-Paton index $A$ in the adjoint representation of a Lie group. This can be rewritten in terms of a field defined as

$$\Phi(\phi, \theta, \varphi) = \sum_{s,A,J,m} T^A g_{s,A} M^s(J, m) D^J_{m,0} (\varphi, \theta, 0) e^{i(sJ - 1)\phi},$$

$$\bar{\Phi}^{(\pm)}(\phi, \theta, \varphi) = \sum_{A,J,m} T^A \bar{g}^{(\pm)A} M^{\pm}(J, m) D^J_{m,0} (\varphi, \theta, 0) e^{i(\pm J - 1)\phi}, \quad (25)$$

where $M^s$ are normalization constants defined by

$$M^+(J, m) = \frac{(J - 1)!}{\sqrt{(2J - 1)!}} N(J, m), \quad M^-(J, m) = \frac{(-1)^m (2J + 1) \sqrt{(2J - 1)!}}{4\pi (J - 1)! N(J, m)}. \quad (26)$$

Note that the fields with $N_{NF} \neq 0$ have opposite statistics to those with $N_{FP} = 0$. Using the Poisson brackets for the rotation matrix

$$\{D_{m_1,0}^{J_1}, D_{m_2,0}^{J_2}\} = \sqrt{\frac{N(J_3, m_3)}{N(J_1, m_1) N(J_2, m_2)}} \left[ \frac{(2J_1 - 1)!(2J_2 - 1)!}{(2J_3 - 1)!} \frac{(J_3 - 1)!}{(J_1 - 1)!(J_2 - 1)!} (J_2 m_1 - J_1 m_2) D_{m_3,0}^{J_3} \right], \quad (27)$$

we finally obtain

$$S_3^{(1)} = -2ig_0 \int d\phi e^{2\phi} \int_{S^2} d^2x \varepsilon^{ij} \text{Tr} \left( \bar{\Phi}^{(-)} \frac{\partial \bar{\Phi}^{(+)}}{\partial x^i} \frac{\partial \Phi}{\partial x^j} \right), \quad (28)$$

where $x^i = (\theta, \varphi)$. This action is essentially identical to that for the states with $N_{FP} = 0$.

What is the physical meaning of this result? The form of the effective action reminds us the similar structure in the BRST formulation of the nonabelian gauge theory. Here similarly we suspect that the whole theory may have a BRST-like invariance in the target space, just as in the string field theory. In fact, we can see the symmetry in the present cubic action. For this purpose, it is convenient to write the action for the states without ghost number in terms of the fields

$$\Phi^{(\pm)}(\phi, \theta, \varphi) = \sum_{A,J,m} T^A g_{s,A} M^{\pm}(J, m) D^J_{m,0} (\varphi, \theta, 0) e^{i(\pm J - 1)\phi}. \quad (29)$$
The total action for the cubic terms is then
\[ S_3 = ig_0 \int d\phi e^{2\phi} \int_{S^2} d^2 x \varepsilon^{ij} \text{Tr} \left( \Phi^(-) \frac{\partial \Phi^(+)}{\partial x^i} \frac{\partial \Phi^(+)}{\partial x^j} - 2\tilde{\Phi}^(-) \frac{\partial \Phi^(+)}{\partial x^i} \frac{\partial \Phi^(+)}{\partial x^j} \right), \] (30)

which is invariant under
\[ \delta \Phi^(+)=\lambda \tilde{\Phi}^(+), \quad \delta \tilde{\Phi}^(-)=\lambda \Phi^(-), \]
\[ \delta \tilde{\Phi}^(+)=0, \quad \delta \Phi^(-)=0. \] (31)

Note that this is a nilpotent transformation. This symmetry is similar to the transformation generated by the charge
\[ Q \sim \oint \frac{dz}{2\pi i} b(z)e^{-i(\xi X(z)+\bar{\xi} \bar{X})}/\sqrt{2}, \] (32)

which is the operator to create the states with ghost number from those without it (see eq. (5)), except that this changes the spins of the states (for example, \( Q\Psi^+(J,m) = \tilde{\Psi}^+(J-1/2,m-1/2) \)). The action can be written as
\[ S_3 = ig_0 \int d\phi e^{2\phi} \int_{S^2} d^2 x \varepsilon^{ij} \delta \left[ \text{Tr} \left( \Phi^(-) \frac{\partial \Phi^(+)}{\partial x^i} \frac{\partial \Phi^(+)}{\partial x^j} \right) \right], \] (33)

It is then natural to conjecture that these “ghost degrees of freedom” play the role of canceling part of the contribution from the \( N_{FP}=0 \) states, just as the Faddeev-Popov ghosts. Indeed, Bershadsky and Klebanov [19] have computed the one-loop partition function in \( c=1 \) gravity and found that it contains the contribution only from the primary fields of the form \( e^{i(pX+\bar{p}\bar{X})} \) (those for \( |m| = J \) and “tachyons”) because the contribution of each special primary field with \( |m| \leq J-1 \) cancels with that of the descendant of the previous one. To really check this possibility in our approach, we have to examine the one-loop contribution by using the effective action.

We hope that our finding that the “ghost states” have nonvanishing correlations and that they have the same structure as that of the \( N_{FP}=0 \) states will help to get further insight into the theory.

**Acknowledgement**

We would like to thank N. Sakai and Y. Tanii for useful discussions.
References

[1] E. Brézin and V. Kazakov, Phys. Lett. **B236** (1990) 144; M. R. Douglas and S. Shenker, Nucl. Phys. **B335** (1990) 635; D. J. Gross and A. A. Migdal, Phys. Rev. Lett. **64** (1990) 127; Nucl. Phys. **B340** (1990) 333.

[2] G. Parisi, Phys. Lett. **B238** (1990) 209; D. J. Gross and N. Miljković, Phys. Lett. **B238** (1990) 217; E. Brézin, V. Kazakov and Al. B. Zamolodchikov, Nucl. Phys. **B338** (1990) 637; P. Ginsparg and J. Zinn-Justin, Phys. Lett. **B240** (1990) 333; J. Ambjørn, J. Jurkiewicz and A. Krzywicki, Phys. Lett. **B243** (1990) 209, 213.

[3] J. Distler and H. Kawai, Nucl. Phys. **B321** (1989) 509; F. David, Mod. Phys. Lett. **A3** (1989) 1651.

[4] N. Seiberg, Prog. Theor. Phys. Suppl. **102** (1991) 319; J. Polchinski, in *Strings ’90*, ed. R. Arnowitt et al. (World Scientific, Singapore, 1991) p.62; Nucl. Phys. **B357** (1991) 241.

[5] D. J. Gross and I. Klebanov, Nucl. Phys. **B344** (1990) 475;
D. J. Gross, I. Klebanov and M. Newmann, Nucl. Phys. **B350** (1990)333.

[6] A. M. Polyakov, Mod. Phys. Lett. **A6** (1991) 635.

[7] B. H. Lian and G. J. Zuckerman, Phys. Lett. **B254** (1991) 417; Phys. Lett. **B266** (1991) 21.

[8] P. Bouwknegt, J. M. McCarthy and K. Pilch, CERN preprints, CERN-TH.6162/91 (1991); TH.6279/91 (1991).

[9] K. Itoh and N. Ohta, Fermilab preprint, FERMILAB-PUB-91/228-T (1991), to appear in Nucl. Phys. **B**.

[10] E. Witten, Nucl. Phys. **B373** (1992) 187.

[11] I. Klebanov and A. M. Polyakov, Mod. Phys. Lett. **A6** (1991) 3273.
[12] D. Kutasov, E. Martinec and N. Seiberg, Rutgers preprint, RU-91-49 (1991).

[13] K. Itoh and N. Ohta, Osaka preprint, OS-GE 22-91 (1991), to appear in Prog. Theor. Phys. Suppl. (1992).

[14] P. Bouwknegt, J. M. McCarthy and K. Pilch, CERN preprint, CERN-TH.6346/91 (1991).

[15] Y. Matsumura, N. Sakai and Y. Tani, Tokyo Inst. of Tech. preprint, TIT/HEP-186 (1992).

[16] K. Itoh, Texas A & M preprint, CTP-TAMU-42/91 (1991).

[17] I. Bakas, Phys. Lett. B228 (1989) 57;
    C. Pope, L. Romans and X. Sen, Nucl. Phys. B339 (1990) 191;
    E. Bergshoeff, M. P. Blencowe and K. S. Stelle, Comm. Math. Phys. 128 (1990) 213.

[18] T. Kugo and I. Ojima, Prog. Theor. Phys. Suppl. 66 (1979) 1.

[19] M. Bershadsky and I. R. Klebanov, Nucl. Phys. B360 (1991) 559.