Dilatonic Supergravity in Two Dimensions and the Disappearance of Quantum Black Hole

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ABSTRACT

We analyze a supergravity theory coupled to a dilaton and superconformal matters in two dimensions. This theory is classically soluble and we find all the solutions appeared in Callan, Giddings, Harvey and Strominger’s dilatonic gravity also satisfy the constraints and the equations of motion in this supersymmetric theory. We quantize this theory by following the procedure of Distler, Hlousek and Kawai. In the quantum action, the cosmological term is renormalized to vanish. As a result, any solution corresponding to classical black hole does not appear in the quantum theory, which should be compared with the non-supersymmetric case.
1. Introduction

The dilaton gravity theory proposed by Callan, Giddings, Harvey and Strominger\cite{1}(CGHS) is very instructive for the understanding of black hole physics. Especially, the problems associated with Hawking radiation\cite{2} have been discussed by using this toy model.\textsuperscript{3–17} In the original paper by CGHS, the quantum effects, such as the Hawking radiation and its back reaction of the metric, were expected to be described by adding correction term, which only comes from the conformal anomaly,\textsuperscript{18} to the classical action. Several authors,\textsuperscript{7–9} however, have claimed that the procedure of David\textsuperscript{19} and of Distler and Kawai\textsuperscript{20} is necessary when we quantize this theory consistently. By using this procedure, it has been found\textsuperscript{12,13} that the quantum theory has no lower bound in energy and it has been conjectured that this problem will be resolved by supersymmetrizing the theory. In this paper, we propose a supergravity theory coupled to a dilaton in two dimensions. This theory is also classically soluble and we find all the solutions found in Callan, Giddings, Harvey and Strominger’s dilatonic gravity also satisfy the constraints and the equations of motion in this supersymmetric theory. We quantize this theory by following the procedure of Distler, Hlousek and Kawai.\textsuperscript{21} In the quantum action, the cosmological term is renormalized to vanish. As a result, any solution corresponding to classical black hole does not appear in the quantum theory. This might tell that supersymmetry would forbid the existence of black hole in quantum theory even in higher dimensions. In the next section, we propose the classical action of dilatonic supergravity by using the tensor calculus by Higashijima, Umetsu, Yu,\textsuperscript{22,23} which is based on conformal supergravity.\textsuperscript{23–25} We show that all the solutions, including black hole solutions, in CGHS theory are also solution of this supersymmetric theory. Here it is interesting to observe that the cosmological constant is always positive semi-definite as in CGHS model. This situation is analogous to that of four dimensional supergravity, which can be constructed only in the anti-de Sitter space. In Section 3, we quantize this action following the procedure of Distler, Hlousek and Kawai. The cosmological term cannot appear if we require that the quantum action has superconformal symmetry. This tells that,
in quantum theory, there is not any solution corresponding to black hole solutions in the classical theory. By bosonizing the fermion fields, we find that the equations of motion which are obtained from the effective action of the quantum theory are Liouville equations. The last section is devoted to summary and discussion.

2. Classical Black Hole

We start from the following action. This action will describe the effective action of superstring in two dimensional black hole background. We use the notations and the tensor calculus in the papers by Higashijima, Uematsu, Yu.

\[
S = \frac{1}{2\pi} \int d^2 x \left( -2[\tilde{\Phi}^2 \otimes W]_{\text{inv}} - 4[\tilde{\Phi} \otimes T(\tilde{\Phi})]_{\text{inv}} + 4\lambda [\tilde{\Phi}^2]_{\text{inv}} + \sum_i \frac{1}{2} [\tilde{\Sigma}_i \otimes T(\tilde{\Sigma}_i)]_{\text{inv}} \right)
\]

\[
= \frac{1}{2\pi} \int d^2 x e \left\{ -R\phi^2 + 2S(\phi f' + \tilde{\zeta}\zeta) - 4\phi\tilde{\zeta}\sigma^{\mu\nu}\psi_{\mu\nu} + 4g^{\mu\nu}\partial_\mu\phi\partial_\nu\phi + 4\tilde{\zeta}\gamma^\mu\partial_\mu\zeta - 4f'^2 - 4\tilde{\psi}_\nu\gamma^\mu\gamma^\nu\zeta\partial_\mu A - \frac{1}{2}\tilde{\zeta}\tilde{\psi}_\nu\gamma^\mu\gamma^\nu\psi_\mu \right. \\
+ \left. 4\lambda \left( 2\phi f' - \tilde{\zeta}\zeta + \phi\tilde{\psi}_\mu\gamma^\mu\zeta + \frac{1}{2}\tilde{\psi}_\mu\sigma^{\mu\nu}\psi_\nu + S\phi^2 \right) \right.
\]

\[
+ \sum_i \left( -\frac{1}{2} g^{\mu\nu}\partial_\mu a_i\partial_\nu a_i - \frac{1}{2}\tilde{\xi}_i\gamma^\mu\partial_\mu\xi_i + \frac{1}{2} G^2_i \\
+ \frac{1}{2}\tilde{\psi}_\nu\gamma^\mu\gamma^\nu\xi_i\partial_\mu a_i - \frac{1}{16}\xi_i\xi_i\tilde{\psi}_\nu\gamma^\mu\gamma^\nu\psi_\mu \right) \right\}
\]

Here \(\Sigma_i\)'s are matter scalar multiplets \(\Sigma_i = (a_i, \xi_i, G'_i), i = 1, \cdots, N\). A scalar multiplet \(\Phi = (\phi, \zeta, f')\) is given in terms of a dilaton multiplet \(\Phi = (\phi, \chi, F')\) by

\[
\tilde{\Phi} \equiv e^{-\Phi} = (e^{-\phi}, -e^{-\phi}\chi, -e^{-\phi}(F' + \frac{1}{2}\tilde{\chi}\chi)).
\]
$W$ is a curvature multiplet,

$$W = (S, \eta, -S^2 + \frac{1}{2} R - \frac{1}{2} \bar{\psi}^\mu \gamma^\nu \psi_{\mu \nu} + \frac{1}{4} S \bar{\psi}^\mu \psi_\mu) .$$

(2.3)

Here $\eta$ and $\psi_{\mu \nu}$ are defined by,

$$\eta \equiv - \frac{1}{2} S \gamma^\mu \psi_\mu + \frac{1}{2} e^{-\frac{1}{2} \epsilon^{\mu \nu} \gamma_5} \psi_{\mu \nu} ,$$
$$\psi_{\mu \nu} \equiv D_\mu \psi_\nu - D_\nu \psi_\mu .$$

(2.4)

$T(\Sigma)$ is a kinetic multiplet which is defined for a scalar multiplet $\Sigma = (a, \xi, G')$ and $[\Sigma]_{\text{inv}}$ expresses the invariant Lagrangian density which is given by

$$[\Sigma]_{\text{inv}} \equiv e [G' + \frac{1}{2} \bar{\psi}^\mu \gamma^\mu \xi + \frac{1}{2} \alpha \bar{\psi}_\mu \sigma^{\mu \nu} \psi_\nu + Sa] .$$

(2.5)

The action (2.1) is, by construction, invariant under the following local supersymmetry transformation,

$$\delta e_\mu^a = \bar{\epsilon} \gamma^a \psi_\mu$$
$$\delta \psi_\mu = 2(\partial_\mu + \frac{1}{2} \omega_\mu \gamma_5 + \frac{1}{2} \gamma_\mu S) \epsilon$$
$$\delta S = - \frac{1}{2} S \bar{\epsilon} \gamma^\mu \psi_\mu + \frac{1}{2} i e^{-\frac{1}{2} \epsilon^{\mu \nu} \gamma_5} \psi_{\mu \nu}$$
$$\delta \phi = \bar{\epsilon} \chi$$
$$\delta \chi = \{ F' + \gamma^\mu (\partial_\mu \phi - \frac{1}{2} \bar{\psi}_\mu \chi) \} \epsilon$$
$$\delta F' = \bar{\epsilon} \gamma^\mu \left\{ \left( \partial_\mu + \frac{1}{2} \omega_\mu \gamma_5 \right) \chi ight. \right.$$
$$\left. - \frac{1}{2} \gamma^\nu \left( \partial_\nu A - \frac{1}{2} \bar{\psi}_\nu \chi \right) \psi_\mu - \frac{1}{2} F' \psi_\mu \right\}$$
$$\delta a_i = \bar{\epsilon} \xi_i$$
$$\delta \xi_i = \{ G' i + \gamma^\mu (\partial_\mu a_i - \frac{1}{2} \bar{\psi}_\mu \xi_i) \} \epsilon$$
$$\delta G'_i = \bar{\epsilon} \gamma^\mu \left\{ \left( \partial_\mu + \frac{1}{2} \omega_\mu \gamma_5 \right) \xi_i ight. \right.$$
$$\left. - \frac{1}{2} \gamma^\nu \left( \partial_\nu \phi - \frac{1}{2} \bar{\psi}_\nu \xi_i \right) \psi_\mu - \frac{1}{2} G' i \psi_\mu \right\} .$$

Here $\epsilon$ is an anti-commuting spinor parameter of local supersymmetry transforma-
tion and $\omega_\mu$ is the spin connection and given by
\[
\omega_\mu = -ie^{-1}e_{a\mu}e^{\lambda\nu}\partial_\lambda e^a_\nu - \frac{1}{2}\bar{\psi}_\mu\gamma^\lambda\gamma^\lambda\psi_\lambda. \tag{2.7}
\]

We now show that all the classical solutions found in Ref.1 satisfy the constraints and the equations of motion which are derived from the action (2.1). In order to do this, we set all the fermionic fields to vanish, which are solutions of all the constraints and the equations of motion which are given by the variation of the fermionic fields $\psi_\mu$, $\chi$ (or $\zeta$) and $\xi_i$. Then by integrating the the auxiliary fields $S$, $F'$ (or $f'$) and $G'_i$, we obtain the following classical action which has appeared in the paper by Callan, Giddings, Harvey and Strominger, \(^1\)
\[
S = \frac{1}{2\pi} \int d^2x \sqrt{-g} \left[ e^{-2\phi} \left( -R + 4(\nabla\phi)^2 + 4\lambda^2 \right) - \frac{1}{2} \sum_i g^{\mu\nu} \partial_\mu a_i \partial_\nu a_i \right]. \tag{2.8}
\]

This tells that all the classical solutions found in Ref.1, including the solutions describing the formation of a black hole by collapsing matter, are also solutions of this supersymmetric theory.

We note that we cannot construct a supersymmetric model of the dilaton gravity when the cosmological constant $\lambda^2$ is negative.

### 3. Quantum Effects

In the original paper by CGHS, the quantum effects were expected to be described by adding correction term, which only comes from the conformal anomaly, \(^1\)\(^8\) to the classical action. In the supersymmetric model proposed here, this corresponds to add the following term, \(^2\)
\[
S_{\text{anomaly}} = \frac{\kappa}{2\pi} \int d^2x \left[ -\frac{1}{2} \partial_\mu \rho \partial^\mu \rho - \frac{1}{2} \bar{\psi}\gamma^\mu \partial_\mu \psi + \frac{1}{2} eS^2 \right], \tag{3.1}
\]
\* The local supersymmetric form of this action is given in Ref.28.
when we choose the following superconformal gauge fixing condition,

\[ g_{\mp} = -\frac{1}{2}e^{2\rho}, \quad g_{\pm} = 0, \quad \psi_{\mu} = \gamma_{\mu}\psi. \tag{3.2} \]

Here \( \kappa = \frac{8-N}{4} \) is a constant, which should be determined by the conformal anomaly. We need more counterterms since the quantum action should have superconformal symmetry when we choose the superconformal gauge (3.2). By following de Alwis’ paper, we assume the kinetic term is given by

\[ S_{\text{kin}} = \frac{1}{2\pi} \int d^2x d^2\theta \left[ -4 e^{-2\Phi} (1 + h(\hat{\Phi})) \bar{D}\hat{\rho}D\hat{\Phi} \\
+ 2 e^{-2\Phi} (1 + \bar{h}(\hat{\Phi})) (\bar{D}\hat{\Phi}D\hat{\rho} + \bar{D}\hat{\rho}D\hat{\Phi}) + \kappa \bar{D}\hat{\rho}D\hat{\rho} \right]. \tag{3.3} \]

Here we have used superfield notations and \( \hat{\Phi} \) and \( \hat{\rho} \) are superfields defined by

\[ \hat{\Phi} \equiv \phi + \bar{\theta}\chi + \frac{1}{2} \bar{\theta}\theta F', \]
\[ \hat{\rho} \equiv \rho + \bar{\theta}\psi + \frac{1}{2} \bar{\theta}\theta S. \tag{3.4} \]

\( \theta \) and \( \bar{\theta} \) are anti-commuting coordinates and \( D \) and \( \bar{D} \) are covariant derivatives. If we define new fields \( \hat{X} \) and \( \hat{Y} \) by

\[ \hat{X} = 2 \sqrt{\frac{2}{|\kappa|}} \int dte^{-2t} \sqrt{(1 + \bar{h}(t))^2 + \kappa e^{2t}(1 + h(t))}, \]
\[ \hat{Y} = \sqrt{2|\kappa|} \left( \hat{\rho} - \frac{1}{\kappa} + \frac{2}{\kappa} \int dte^{-2t} \bar{h}(t) \right). \tag{3.5} \]

the kinetic term is rewritten by

\[ S_{\text{kin}} = \frac{1}{2\pi} \int d^2x d^2\theta [\mp \bar{D}\hat{X}D\hat{X} \pm \bar{D}\hat{Y}D\hat{Y}] \tag{3.6} \]

Here upper/lower signs correspond to \( \kappa > 0 / \kappa < 0 \), respectively. From now on, we will consider the case with lower signs in Eq.(3.6) since another case can be treated in a similar way.
If we assume that there is any interaction term with respect to \( \hat{X} \) and \( \hat{Y} \), the energy momentum tensor \( T \) has the following form,

\[
T = T_X + T_Y + T_\Sigma + T_{\text{ghost}},
\]

\[
T_X = -\frac{1}{2}(\partial X \partial X - \chi_X \partial \chi_X),
\]

\[
T_Y = \frac{1}{2}(\partial Y \partial Y - \chi_Y \partial \chi_Y) + \sqrt{\frac{|\kappa|}{2}} \partial^2 Y.
\]

Here we have written \( \hat{X} \) and \( \hat{Y} \) fields in the components \( \hat{X} = X + \bar{\theta}\chi_X + \frac{1}{2} \bar{\theta}\theta F_X \) and \( \hat{Y} = Y + \bar{\theta}\chi_Y + \frac{1}{2} \bar{\theta}\theta F_Y \). \( T_\Sigma \) and \( T_{\text{ghost}} \) are energy momentum tensors of matter fields and ghost fields and they contribute to the central charge by \( \frac{3}{2} N \) and \( -15 \), respectively. The contribution to the central charge by \( T_X \) is \( \frac{3}{2} \) and that by \( T_Y \) is given by

\[
c_Y = \frac{3}{2}(1 + 4\kappa) = \frac{3}{2}(-N + 9).
\]

Therefore the total central charge \( c \) vanishes: \( c = \frac{3}{2} N - 15 + \frac{3}{2} + c_Y = 0 \). We now introduce interaction term \( V \) so that the term does not violate the super-conformal symmetry. This requires that \( V \) should be given by a vertex operator \( V = e^{\alpha \hat{X} + \beta \hat{Y}} \) whose conformal dimension is \( (\frac{1}{2}, \frac{1}{2}) \), i.e.,

\[
\frac{1}{2} \alpha^2 - \frac{1}{2} \beta(\beta + \sqrt{2|\kappa|}) = \frac{1}{2}.
\]

If we impose the condition that \( T \) is proportional to \( e^{\rho-\phi} \) in the weak coupling limit, we find \( \alpha = \beta \), i.e.,

\[
\alpha = \beta = \sqrt{\frac{1}{2|\kappa|}} = \sqrt{\frac{2}{|N-8|}}.
\]
Therefore we find the quantum theory is described by the following effective action:

\[
S_q = \frac{1}{2\pi} \int d^2x \ d^2\theta \left[ \bar{D} \hat{X} D \hat{X} - \bar{D} \hat{Y} D \hat{Y} + 2\tilde{\lambda}e\sqrt{\frac{2}{|N-8|}(X+\bar{Y})} + \sum_{i=1}^{N} \bar{D} \hat{\Sigma} D \hat{\Sigma} \right]
\]

\[
= \frac{1}{2\pi} \int d^2x \left[ -\left( \partial_{\mu}X \partial^{\mu}X - i\bar{\chi}X \gamma^\mu \partial_{\mu} \chi_X - F_X^2 \right) \\
+ \left( \partial_{\mu}Y \partial^{\mu}Y - i\bar{\chi}Y \gamma^\mu \partial_{\mu} \chi_Y - F_Y^2 \right) \\
+ \tilde{\lambda} \sqrt{\frac{2}{|N-8|}} e\sqrt{\frac{2}{|N-8|}(X+\bar{Y})} \left\{ F_X + F_Y \\
- \frac{1}{2} \sqrt{\frac{2}{|N-8|}} (\bar{\chi}X + \bar{\chi}Y)(\chi_X + \chi_Y) \right\} \\
+ \sum_{i} \frac{1}{2} \left\{ -\partial_{\mu}a_i \partial^{\mu}a_i + i\xi_i \gamma^\mu \partial_{\mu} \xi_i + G_i^2 \right\} \right]
\]

(3.11)

Here \( \hat{\Sigma}_i \)'s are matter superfields: \( \hat{\Sigma}_i = a_i + \bar{\theta} \xi_i + \frac{1}{2} \bar{\theta} \theta G_i \). By integrating auxiliary fields \( F_X, F_Y \) and \( G_i \), we obtain,

\[
S'_{q} = \frac{1}{2\pi} \int d^2x \left[ -\left( \partial_{\mu}X \partial^{\mu}X - i\bar{\chi}X \gamma^\mu \partial_{\mu} \chi_X \right) \\
+ \left( \partial_{\mu}Y \partial^{\mu}Y - i\bar{\chi}Y \gamma^\mu \partial_{\mu} \chi_Y \right) \\
- \frac{\tilde{\lambda}}{|N-8|} e\sqrt{\frac{2}{|N-8|}(X+\bar{Y})} \sqrt{\frac{2}{|N-8|}} (\bar{\chi}X + \bar{\chi}Y)(\chi_X + \chi_Y) \\
+ \sum_{i} \left\{ -\partial_{\mu}a_i \partial^{\mu}a_i + i\xi_i \gamma^\mu \partial_{\mu} \xi_i \right\} \right]
\]

(3.12)

Note that terms like \( e^{2 \sqrt{\frac{2}{|N-8|}(X+\bar{Y})}} \) do not appear. When we consider more general interaction term \( V =: e^{\alpha \bar{X} + \beta \bar{Y}} \), which satisfy Equation (3.9) and \( \alpha \neq \beta \), the term like \( V =: e^{2(\alpha \bar{X} + \beta \bar{Y})} \) does not appear as long as \( V \) is exactly marginal, i.e., \( V(x)V(y) \sim O((x-y)^\delta), \delta > 0 \), which tells \( V(x)^2 = 0 \).
The equations of motion for $X$ and $Y$ are given by

$$0 = \partial_\mu \partial^\mu X - \frac{\tilde{\lambda}}{2} \left( \frac{2}{|N-8|} \right)^{\frac{3}{2}} e^{\sqrt{\frac{2}{|N-8|}(X+Y)}} (\bar{\chi}X + \bar{\chi}Y)(\chi X + \chi Y)$$

$$0 = - \partial_\mu \partial^\mu Y - \frac{\tilde{\lambda}}{2} \left( \frac{2}{|N-8|} \right)^{\frac{3}{2}} e^{\sqrt{\frac{2}{|N-8|}(X+Y)}} (\bar{\chi}X + \bar{\chi}Y)(\chi X + \chi Y)$$

(3.13)

If we consider solutions where all the fermion fields vanish, $X$ and $Y$ are given by the sums of holomorphic and anti-holomorphic functions. This tells that there is not any solution corresponding to black hole solution in the classical theory. Supersymmetry forbids the existence of black hole in the quantum theory.

If we bosonize the fermion fields $\chi X$ and $\chi Y$: $\chi X \pm \chi Y \sim \pm e^{\mp \vartheta}$, we obtain the following equations of motion,

$$0 = \partial_\mu \partial^\mu X - \frac{\tilde{\lambda}}{2} \left( \frac{2}{|N-8|} \right)^{\frac{3}{2}} e^{\sqrt{\frac{2}{|N-8|}(X+Y)+\vartheta}}$$

$$0 = - \partial_\mu \partial^\mu Y - \frac{\tilde{\lambda}}{2} \left( \frac{2}{|N-8|} \right)^{\frac{3}{2}} e^{\sqrt{\frac{2}{|N-8|}(X+Y)+\vartheta}}$$

$$0 = \partial_\mu \partial^\mu \vartheta - \frac{\tilde{\lambda}}{2} \left( \frac{2}{|N-8|} \right)^{\frac{3}{2}} e^{\sqrt{\frac{2}{|N-8|}(X+Y)+\vartheta}}$$

(3.14)

We can set $X = Y$ by using the residual symmetry of the reparametrization symmetry or by a coordinate choice. Then we find that $\vartheta$ satisfies the Liouville equation,

$$0 = \partial_\mu \partial^\mu \vartheta - \frac{\tilde{\lambda}}{2} \left( \frac{2}{|N-8|} \right)^{\frac{3}{2}} e^{\vartheta}.$$ 

(3.15)

The equations (3.14) tell that $X$ and $Y$ are given in terms of $\vartheta$,

$$X = -Y = \vartheta + f^+(x^+) + f^-(x^-).$$

(3.16)

* Since the fermion fields $\chi X$ and $\chi Y$ have opposite signatures, these fermion fields can be bosonized by a negative norm boson.
Here $f^\pm$ are arbitrary functions. Note that there are static solutions:

$$f^\pm = 0, \quad e^{-\vartheta} = -\frac{16}{A} \frac{C n^2 (x^+ x^-)^{n-1}}{\{1 - C(x^+ x^-)^n\}^2}.$$ (3.17)

Here $A = \tilde{\lambda} \left( \frac{2}{|N-8|} \right)^{\frac{3}{2}}, C$ is an arbitrary constant and $n$ is an integer.

4. Summary and Discussion

We have analyzed a supergravity theory coupled to a dilaton and superconformal matters in two dimensions. This theory is classically soluble and we have found all the solutions appeared in Callan, Giddings, Harvey and Strominger’s dilatonic gravity also satisfy the constraints and the equations of motion in this supersymmetric theory. When we quantize this theory following the procedure of Distler, Hlousek and Kawai, the cosmological term is renormalized to vanish in the quantum action. As a result, any solution corresponding to classical black hole does not appear in the quantum theory, which should be compared with the non-supersymmetric case.

It should be amazing that supersymmetry forbids the existence of quantum black hole although classical black hole is allowed to exist. One of the motivations of the present work was to build a theory of dilaton gravity where the Bondi mass of black hole is bounded from below and the theory has a ground state. The above motivation becomes, however, irrelevant since there is no quantum black hole in our model.

However, from the alternative viewpoint, our model gives us an interesting conjecture. Let us assume that the quantum black hole disappears even in four dimensions although it is difficult to take account of quantum effects owing to the non-renormalizability of gravity. In four dimensions, the Schwarzschild radius of black holes and the Compton wave length of elementary particles become comparable at the Planck length scale. This might suggest that we should include
the black hole–like states in the Hilbert space. From the observation done in the present article, however, we might conjecture that, if supersymmetry is realized at the Planck scale, such black hole–like states need not to be included as quantum states in quantum gravity where only quantum states expressing smooth space-time structure are admitted.

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