2d gravity stress tensor and the problem of the calculation of the multi-loop amplitudes in the string theory

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Abstract

It is shown that the vacuum value of the 2d gravity stress tensor in the free field theory is singular in the fundamental region on the complex plane where the genus-$n > 1$ Riemann surface are mapped. Because of the above singularity, one can not construct modular invariant multi-loop amplitudes. The discussed singularity is due to the singularity in the vacuum value of the 2d gravity field that turns out to be on the genus-$n > 1$ Riemann surfaces.

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1 Introduction

It is well known that the conform anomaly \[1,2\] hampers the construction of the string theory in non-critical dimensions. Because of the above anomaly, the string amplitudes in such a theory should depend on a choice of 2-dim. fixed reference metrics. To remove this shortcoming, it has been proposed \[3\] to write down the string amplitudes in the form of the integrals over both the string fields and the 2-dim. metrics. As far as the integration over the above metrics has been performed, in this case the string amplitudes can not depend on fixed reference metrics. So, the discussed theory is expected to be invariant under both the general two-dim. reparameterization group and the Weyl group, like the critical string one.

In the discussed scheme, it is presumed that the conform anomaly is canceled owing to the 2\textit{d} gravity field \[4,5,6\]. On the classical level the action of this field has the Liouville form. Quantum corrections do not change this form, only the parameters in the above action are renormalized \[5,6,7,8\].

The above scheme is mainly reduced \[4,5,6,7,8\] to the free field theory. This theory contains the string and ghost fields and the \(\Phi\) field of the 2\textit{d} gravity, as well. The stress energy-momentum tensor of the \(\Phi\) field includes the linear in \(\Phi\) term to cancel the conform anomaly in the whole stress energy-momentum tensor. Owing to this term, the \(\Phi\) field receives an additional term under two-dim. coordinate transformations except only linear transformations. So on the closed Riemann surfaces of a genus-\(n > 1\) the discussed \(\Phi\) field is necessarily receives an additional term under twists about B-cycles. One concludes by this that the vacuum expectation value \(< \Phi >\) of the \(\Phi\) field is unequal to zero.

In this paper we calculate \(< \Phi >\) on the closed Riemann surfaces of a genus-\(n \geq 1\) using the Schottky group parameterization \[9,10\]. As in \[5,6\], the conformal gauge is employed. We show that \(< \Phi >\) appears to be singular in the fundamental region \(\Omega\) on the complex plane where the Riemann surfaces are mapped.

Moreover, for \(n \geq 2\) the above singularity in \(< \Phi >\) originates the vacuum expectation value of the stress tensor to be singular in the above \(\Omega\) region. Due to the discussed singularity in the stress tensor it appears impossible to construct the multi-loop partition functions, which would be consistent with the requirement of the modular invariance.

Only the one-loop partition function can be constructed because in this
case the stress tensor appears to be non-singular in the above $\Omega$ region on
the complex plane where the genus-1 Riemann surfaces are mapped. Moreover, one can use
the linear kleinian groups instead of the Schottky ones by mapping the genus-1 Riemann surfaces into parallelograms. In this case the $<\Phi>$ value is equal to zero. The one-loop partition function in the discussed parameterization has been calculated in [8]. We show that the same result is obtained also in the Schottky parameterization, the contribution due to
$<\Phi'>$ being taken into account.

The paper is organized as it follows. In Sec.2 we calculate $<\Phi>$, the Schottky groups being employed. We show that $<\Phi>$ is singular in the fundamental region on the complex plane where the Riemann surfaces are mapped. We show also that the above singularity for the Riemann surfaces of a genus-$n \geq 2$ originates the singularity in the vacuum value of the stress energy-momentum tensor. In Sec.3 we discuss the problem of the construction of the multi-loop partition functions. The one-loop partition function is discussed in the Sec.4.

2 Vacuum value of the 2d gravity stress tensor

As it has been noted in the Sec.1, we employ the Schottky parameterization. In this case the genus-$n$ kleinian group is generated by the base Schottky transformations $z \to g_s(z)$ where

$$g_s(z) = \frac{a_s z + b_s}{c_s z + d_s}, \quad a_s d_s - b_s c_s = 1 \quad \text{and} \quad s = 1, 2, \cdots, n. \quad (1)$$

Every transformation $g_s$ maps the $C^{(+)}_s$ circle into the $C^{(-)}_s$ one where

$$C^{(+)}_s = \{z : |c_s z + a_s| = 1\} \quad \text{and} \quad C^{(-)}_s = \{z : |c_s z + d_s| = 1\}. \quad (2)$$

So the fundamental region can be chosen to be the exterior of all the $C^{(-)}_s$ and $C^{(+)}_s$ circles [9, 10]. Also, it is worth-while to note that (1) can be rewritten down as

$$\frac{g_s(z) - u_s}{g_s(z) - v_s} = k_s \frac{z - u_s}{z - v_s} \quad (3)$$
where \( u_s \) and \( v_s \) are the fixed point coordinates and \( k_s \) is the multiplier. One can think that \(|k_s| < 1\).

As in \([5, 6]\), the conformal gauge is used. The stress energy-momentum tensor \( T \) is given by

\[
T = T_m + T_{gh} + T_\Phi
\]  

(4)

where \( T_m \) denotes the contribution due to the string fields, \( T_{gh} \) is the ghost contribution and \( T_\Phi \) is the contribution of the 2\( d \) gravity field \( \Phi \). The \( T_m \) value is given by

\[
T_m = -\frac{1}{2} \partial X_M \partial X^M
\]  

(5)

where \( X^M \) are the string fields, \( M = 1, 2 \cdots d \). The above fields are normalized as ( the \( M \) index is omitted )

\[
X(z_1, \bar{z}_1)X(z_2, \bar{z}_2) \rightarrow -\ln |z_1 - z_2|^2
\]  

(6)

at \( z_1 \rightarrow z_2 \). The \( T_{gh} \) is given by

\[
T_{gh} = 2b \partial c + (\partial b)c
\]  

(7)

where \( b \) is the tensor ghost field and \( c \) is the vector ghost one, they being normalized as

\[
c(z_1, \bar{z}_1)b(z_2, \bar{z}_2) \rightarrow -\frac{1}{z_1 - z_2}
\]  

(8)

at \( z_1 \rightarrow z_2 \). And \( T_\Phi \) in (1) is equal to \([5, 6]\)

\[
T_\Phi = -\frac{1}{2} \partial \Phi \partial \Phi + f \partial^2 \Phi.
\]  

(9)

The \( \Phi \) field in (3) is normalized by the condition that

\[
\Phi(z_1, \bar{z}_1)\Phi(z_2, \bar{z}_2) \rightarrow -\ln |z_1 - z_2|^2
\]  

(10)

at \( z_1 \rightarrow z_2 \). Being determined from the condition that (4) has not the conform anomaly \([5, 6]\), the \( f \) factor appears to be

\[
f^2 = \frac{25 - d}{12}.
\]  

(11)
It is follows from (9) that under the conformal \( z = z(\tilde{z}) \) transformations the \( \Phi \) field is changed as

\[
\Phi(z) = \tilde{\Phi}(\tilde{z}) - f \ln \left| \frac{\partial z}{\partial \tilde{z}} \right|^2
\]  

(12)

where \( \tilde{\Phi}(\tilde{z}) \) is the 2d gravity field in the new coordinate system. So under the kleinian group transformations (1) the vacuum value \(<\Phi(z, \bar{z})>\) of the \( \Phi \) field is changed as

\[
<\Phi(g_s(z), \bar{z})> = <\Phi(z, \tilde{z})> + 2f \ln |c_sz + d_s|^2.
\]  

(13)

In this case the change of \(<\Phi'>> = \partial_z <\Phi>\) under the transformations (10) is given by

\[
<\Phi'(g_s(z))> = (c_sz + d_s)^2[<\Phi'(z)> + 2f \partial \ln (c_sz + d_s)].
\]  

(14)

To calculate \(<\Phi(z, \tilde{z})>\) we construct the Green function \( K(z, z') \) as

\[
K(z, z') = \sum_{(\Gamma)} \frac{1}{(c_{\Gamma}z + d_{\Gamma})^2[gr(z) - z']}.
\]  

(15)

In (15) the summation is performed over all the group products \( \Gamma = \{ z \rightarrow g_{\Gamma}(z) \} \) of the base transformations (10), \( \Gamma = I \) being included. The \( I \) symbol denotes the identical transformation and \( gr(z) = (a_{\Gamma}z + b_{\Gamma})(c_{\Gamma}z + d_{\Gamma})^{-1} \). The above \( K(z, z') \) Green function satisfies the conditions that

\[
K(g_s(z), z') = Q_s(z)^2 K(z, z'), \quad K(z, g_s(z')) = 2\pi i J'_s(z) + K(z, z')
\]  

(16)

where \( Q_s(z) = (c_sz + d_s) \) and \( J'_s(z) \) are 1-forms, \( J'_s(z) = \partial_z J_s(z) \).

To calculate \(<\Phi(z)>\) we start with the following relation

\[
<\Phi'(z)> = -\int_{C(z)} K(z, z') <\Phi'(z')> \frac{dz'}{2\pi i}
\]  

(17)

where the infinitesimal contour \( C(z) \) surrounds the \( z \) point in the positive direction. In addition to eq.(15), eq.(17) employs that \(<\Phi(z)>\) is a holomorphic function of \( z \). We presume that \(<\Phi(z)>\) has no singularities inside the fundamental region. In this case one can reduce the right side of (17) to the integrals over the boundaries \( C_s^{(+)} \) and \( C_s^{(-)} \) of the fundamental region.
Ω, both \( C^{(+)} \) and \( C^{(-)} \) being given by (2). Furthermore, the \( z \to g_s(z) \) replacement allows to reduce the integral along the \( C^{(+)} \) contour to the integral along \( C^{(-)} \). In this case, employing eqs. (14) and (14), one can derive from (17) that

\[
< \Phi'(z) > = -2f \sum_{s=1}^{n} \int_{C_s^{(-)}} K(z, z') \frac{c_s d z'}{c_s z + d_s} + \sum_{s=1}^{n} h_s J'_s(z)
\]  

(18)

where \( h_s \) are constants. For \( h_s \) to be arbitrary, the change of (18) under every kleinian group transformation (1) differs from (13) by a constant term \( H_s \). The above \( h_s \) constants are just determined from the conditions that \( H_s = 0 \) for \( s = 1, 2, \cdots, n \).

Eq.(18) gives the desired \( < \Phi'(z) > \) value. It follows from (18) and from (16) that

\[
< \Phi'(z) > \to -\frac{2fn}{z} \quad \text{at} \quad z \to \infty
\]  

(19)

At the same time, every non-singular 1-form decreases at \( z \to \infty \) not slower than \( const \cdot z^{-2} \). As example, the \( J'_s(z) \) 1-forms in (18) possess this property. So one concludes that \( < \Phi'(z) > \) is singular at \( z \to \infty \).

Owing to the presence of \( < \Phi'(z) > \), an additional contribution into the vacuum value \( < T_\Phi(z) > \) of the 2d gravity stress tensor \( T_\Phi \) arises. To determine the above contribution, one can substitute in eq.(9) the \( < \Phi'(z) > \) value instead of \( \Phi' \). Because of the discussed contribution, it appears that

\[
< T_\Phi(z) > \to -\frac{2f^2 n(n - 1)}{z^2} \quad \text{at} \quad z \to \infty.
\]  

(20)

Eq.(20) follows from (9) and (19). It is useful to remind that every non-singular 2-form decreases at \( z \to \infty \) not slower than \( const \cdot z^{-4} \). So, for \( n \geq 2 \), \( < T_\Phi(z) > \) is singular at \( z \to \infty \).

One could construct the \( < \Phi(z) > \) value to be non-singular at \( z \to \infty \). In this case \( < \Phi(z) > \) appears to be singular at finite values of \( z \). In any case, for \( n \geq 2 \) the vacuum value \( < T_\Phi(z) > \) of the stress tensor \( T_\Phi \) turns out to be singular in the fundamental region on the complex \( z \)-plane where the Riemann surfaces are mapped. In the next Sec. we show that the above singularity prevents constructing the multi-loop partition functions.
3 Problem of the multi-loop partition functions

Intuition suggests that a singularity of the vacuum value of the stress tensor does the theory to be self-contradictory. To argue this statement in a more rigorous manner, one can note that in the considered theory [5, 6] the string amplitudes are expected to be independent of fixed reference metrics. In this case, the genus-\(n\) partition functions for \(n \geq 2\) satisfy the equations [11] that are non other than Ward identities. In [11] the discussed equations have been obtained in the critical string theory, but the same consideration can be performed for the non-critical string [5, 6].

The partition functions \(Z_n(\{q_N, \overline{q}_N\})\) depend on \((3n - 3)\) complex moduli \(q_N\) together with \(\overline{q}_N\), which are complex conjugated to \(q_N\). The first of the discussed equations for \(Z_n(\{q_N, \overline{q}_N\})\) is given by [11]

\[
\sum_N \chi_N(z) \frac{\partial \ln Z_n(\{q_N, \overline{q}_N\})}{\partial q_N} = \langle T(z) \rangle - \sum_N \frac{\partial}{\partial q_n} \chi_N(z) \quad (21)
\]

and the second equation is obtained from (21) by the complex conjugation. The \(\langle T(z) \rangle\) value denotes the vacuum expectation of the stress tensor (4) and \(\chi_N(z)\) are 2-tensor zero modes. It is worthwhile to note that in eq.(21), the ghost contribution to \(\langle T(z) \rangle\) is calculated in the special ghost scheme [11]. Unlike the well known scheme [10, 12], in the considered scheme [11], the ghost vacuum correlator has no unphysical poles. Among other things, the above correlator takes into account those contributions to the partition functions, which are due to both the moduli volume form and ghost zero modes. In the critical string theory solution of both eq.(21) and its complex conjugated fully determine the partition functions apart only an independent of moduli factor. But in the non-critical string theory [5, 6] eq.(21) has no solutions.

Indeed, being non-singular in the fundamental region, every \(\chi_N(z)\) decreases not slower than \(const \cdot z^{-4}\) at \(z \to \infty\). So the left side of (21) decreases at least as \(z^{-4}\) at \(z \to \infty\). At the same time, the right part of (21) decreases only as \(z^{-2}\), as it follows from eq.(20). So in this case there are no the partition functions, which are consistent with the requirement that the multi-loop amplitudes are independent of a choice of fixed references metrics. It can be proved that the above equation (21) appears to be the modular in-
variant. So the absence of the solutions of eq.(21) is meant that there are no the multi-loop partition functions, which are consistent with the requirement of the modular invariance. Only the genus-1 partition function can be constructed to be modular invariant, as it explained in the Sec.4.

4 One-loop partition function

We write down the one-loop amplitudes \( A_1 \) as

\[
A_1 = \int Z_1(k, \bar{k}) \langle V \rangle d\omega d\bar{\omega}
\]

where \( Z_1(k, \bar{k}) \) is the genus-1 partition function and \( \langle V \rangle \) is the vacuum value of the vertex product integrated over the Riemann surface. The complex period \( \omega \) is given in the terms of the \( k \) multiplier in (3) as

\[
\omega = \frac{1}{2\pi i} \ln k.
\]

As it follows from (20), for \( n = 1 \) the singularity in \( \langle \Phi \rangle \) does not originate the singularity in the vacuum value of the stress tensor. In this case \( \langle \Phi \rangle \) is given by

\[
\langle \Phi \rangle = -f \ln \left| \frac{(z - u)(z - v)}{(u - v)^2} \right|^2 + \text{const}
\]

where \( u \) and \( v \) are the coordinates of the fixed points of the Schottky transformation (1), \( f^2 \) being given by (11). Indeed, one verifies that (24) satisfies (13). It is follows from (24) that

\[
\langle \Phi' \rangle = -f \left[ \frac{1}{z - u} + \frac{1}{z - v} \right].
\]

We write down the vacuum value \( < T_\Phi > \) of the \( T_\Phi \) stress tensor as

\[
<T_\Phi(z)> = <T_\Phi^{(1)}(z)> + T_\Phi^{(2)}(z)
\]

where \( <T_\Phi^{(1)}(z)> \) is calculated in the terms of the vacuum correlator of the \( \Phi \) fields in the accordance with the rules [10] of the conform theory and

\[
T_\Phi^{(2)}(z) = -\frac{1}{2}(<\Phi'(z)>)^2 + f <\Phi''(z)>
\]

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It is follows from (25) and (27) that
\[ T_{\Phi}^{(2)}(z) = \frac{1}{2} f^2 \frac{(u-v)^2}{(z-u)^2(z-v)^2}. \quad (28) \]

One can see from (28) that \( T_{\Phi}^{(2)} \) has no singularities in the fundamental region on the complex \( z \) plane, the \( z = \infty \) point being included.

Moreover, instead of the Schottky parameterization, one can map every genus-1 Riemann surfaces into the parallelogram on the complex \( w \) plane. In this case the kleinian group is given by the linear coordinate transformations
\[ w \rightarrow w + 1 \quad w \rightarrow w + \omega \quad (29) \]
where \( \omega \) is the complex period. The above parameterization can be obtained from that where the Schottky group are employed by the mapping
\[ w = \frac{1}{2\pi i} \ln \frac{z-u}{z-v} \quad (30) \]
with both \( u \) and \( v \) to be the coordinates of the fixed points of the transformation \([1]\). The first of eqs.(29) corresponds to \( 2\pi \)-twist around the \( C^{(+)} \) or \( C^{(-)} \) circle on the complex \( z \) plane, both \( C^{(+)} \) and \( C^{(-)} \) being given by \([1]\). The second of eqs.(29) corresponds to the mapping \([1]\) on the \( z \) plane.

It is follows from (12), (24) and (30) that, the kleinian group (29) being employed, the \( \langle \Phi \rangle \) value is independent of \( w \) and, therefore, \( \langle \Phi' \rangle = 0 \). This result is natural. Indeed, one can see from (12) that \( \langle \Phi \rangle \) is scalar under the kleinian group (29). Besides, \( \langle \Phi' \rangle \) is the holomorphic function of \( z \). Therefore, \( \langle \Phi \rangle \) appears to be the zero mode of the Laplacian. So \( \langle \Phi \rangle = const \) and \( \langle \Phi' \rangle = 0 \). In this case the one-loop partition function can be calculated by the method given in [13], as it has been done in [8].

In the above paper [8] the \( Z_1(k, \bar{k}) \) partition function has been obtained to be
\[ Z_1(k, \bar{k}) = const \frac{|k|^\frac{1}{2} |y|^2 \tilde{Z}_m}{Im\omega} \quad (31) \]
where \( const \) is independent of \( k \) and \( \bar{k} \); \( \omega \) is defined by (23). The \( y = y(k) \) function is given by
\[ y(k) = \prod_{n=1}^{\infty} (1 - k^n). \quad (32) \]
The $\tilde{Z}_m$ multiplier in (31) is due to the string fields. In the considered scheme

$$\tilde{Z}_m = \frac{|k \frac{1}{2\pi} y|^{-2d}}{(Im\omega)^d/2},$$

(33)
y being defined by (32). The factor before $\tilde{Z}_m$ in (31) is due to the ghost and 2d gravity fields. It has been noted in [8] that $\tilde{Z}_m$ being defined by (33), is modular invariant and $Z_1(k, \bar{k})$ satisfies the condition of the modular invariance of the one-loop amplitudes.

The same result is obtained in the Schottky parameterization, the contribution due to $<\Phi'>$ being taken into account. This proves that the discussed contribution is really necessary to obtain the correct result. To prove the above statement, we note that $Z_1(k, \bar{k})$ satisfies eq.(21) just as the multi-loop partition functions do it. In this case, however, the ghost scheme [11] should be modified to take into account the contribution to $Z_1$ due to the zero modes of the vector ghosts. It could be achieved by a suitable choice of the ghost vacuum correlator. But this problem is not considered in the present paper. Instead we note that the ghost contribution to $Z_1(k, \bar{k})$ is the same as in the critical string theory. And we use the critical string calculations [11, 12, 14] to determine the contribution discussed.

We write down the desired $Z_1(k, \bar{k})$ value as

$$Z_1(k, \bar{k}) = Z_\Phi(k, \bar{k})Z_{gh}(k, \bar{k})Z_m(k, \bar{k})$$

(34)

where $Z_\Phi$, $Z_{gh}$ and $Z_m$ are due to the 2d gravity field, the ghost fields and the string fields, respectively. In this case both $Z_\Phi(k, \bar{k})$ and $Z_{gh}(k, \bar{k})$ are determined by the following equations

$$\chi_k(z)\frac{\partial \ln Z_m(k, \bar{k})}{\partial k} = <T_m(z)>,\quad (35)$$

$$\chi_k(z)\frac{\partial \ln Z_\Phi(k, \bar{k})}{\partial k} = <T^{(1)}_\Phi(z) + T^{(2)}_\Phi(z)>,\quad (36)$$

their complex conjugated being added, too. The $T^{(1)}_\Phi(z)$ and $T^{(2)}_\Phi(z)$ values are given by (26) and (28).

It is worth-while to note that the $\chi_k(z)$ tensor zero mode in (35) and in (36) is normalized by [11]

$$\int_{C(-)} \chi_k(z) \frac{(z - u)(z - v)}{k(u - v)} \frac{dz}{2\pi i} = -1$$

(37)
where \( C(-) \) circle is defined by eq.(2). So eq.(28) can be rewritten down as

\[
T_{\Phi}^{(2)}(z) = \frac{f^2}{2k} \chi_k(z).
\]  

(38)

In the Schottky parameterization the \( < T_m > \) value is calculated by the same methods that have been used \[11, 12, 14\] in the critical string theory. As the result, the \( Z_m \) value being calculated from (35) turns out to be

\[
Z_m = \frac{|y|^{-2d}}{(Im\omega)^{d/2}}.
\]  

(39)

where \( y \) is given by (32). Furthermore, using eq.(39) together with the results given in \[13\], we obtain that

\[
Z_{gh} = \frac{|y|^2}{(Im\omega)|k|^2}.
\]  

(40)

Furthermore, \( < T_{\Phi}^{(1)} > \) is calculated by the same method as \( < T_m > \). Moreover, \( T_{\Phi}^{(2)} \) is given by (38), \( f^2 \) being given by (11). In this case one can find from (36) that

\[
Z_{\Phi} = \frac{|k|^{(1-d)/2}|y|^2}{(Im\omega)}.
\]  

(41)

Being compared with (31) and with (33), eqs. (34) and (39)-(41) show that every from the \( Z_m \), \( Z_{gh} \) and \( Z_{\Phi} \) values depends on the parameterization of the Riemann surfaces to be chosen. This is because the conform anomaly presents in every of the \( T_m \), \( T_{gh} \) and \( T_{\Phi} \) values. So every of the above values itself does not appear to be a tensor. But the whole \( Z_1 \) partition function being defined by (31) and (33) coincides with that calculated from (34) and (39)-(41). And the above the partition function satisfies the condition that the one-lop amplitudes are due to be modular invariant. At the same time, the construction of the multi-loop partition functions in the scheme \[4, 6\] appears to be impossible.

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