String Theory №6 [WIP]

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1 Type 0 Superstrings

A closed superstring theory consists of sectors labeled by the boundary conditions of $(\psi, \tilde{\psi})$ along with suitable GSO projections $(-1)^F = \pm 1$. Here we follow the discussions of *Polchinski*, with R: $\alpha = 1$ and NS: $\alpha = 0$.

There are also some consistency conditions: by modular invariance, there must be at least one left-moving R sector and at least one right-moving R sector; on the other hand, the OPE must close, and since $\mathrm{R}\times\mathrm{R} = \mathrm{NS}$ there must be corresponding NS sector for each R sector.

If we includes only the (NS, NS) and the (R, R) sectors, then both must exist due to the above conditions. In fact, closure of OPE implies that the (NS+, NS+) sector must exist. In addition, NS− sector must be paired with another NS− sector due to the level matching condition of the closed string, i.e. it is possible to have a (NS−, NS−) sector.

The full possiblities can be generated by enumerating all possible (R, R) sectors (there is $2\times 2 = 4$ of them), while applying an extra consistency check that all pairs of vertex operators O_1, O_2 are mutually local, i.e.

$$\exp i\pi\Big(F_1\alpha_2 - F_2\alpha_1 - \tilde{F}_1\tilde{\alpha}_2 + \tilde{F}_2\tilde{\alpha}_1\Big) = 1 \tag{1}$$

If $O_1 \in (\mathrm{NS+}, \mathrm{NS+})$, then we have $\alpha_1 = \tilde{\alpha}_1 = 0 = F_1 = \tilde{F}_1$, hence the above factor is always trivial; for $O_1 \in (\mathrm{R}, \mathrm{R})$, however, $\alpha_1 = \tilde{\alpha}_1 = 1$, which yields a non-trivial constraint for the second operator: $F_2 - \tilde{F}_2 = F_1\alpha_2 - \tilde{F}_1\tilde{\alpha}_2 = \alpha_2\big(F_1 - \tilde{F}_1\big) \pmod 2$, assuming $\alpha_2 = \tilde{\alpha}_2$. With $\alpha_2 = 0$ this gives $F_2 = \tilde{F}_2$, and with $\alpha_2 = 1$ this gives $F_2 - \tilde{F}_2 = F_1 - \tilde{F}_1$, which means that all (R, R) sectors have the same sign difference between F and $\tilde{F}$. The possible solutions can then be narrowed down to:

$$0\mathrm{A}:\ (\mathrm{NS+}, \mathrm{NS+}),\ (\mathrm{NS-}, \mathrm{NS-}),\ (\mathrm{R+}, \mathrm{R-}),\ (\mathrm{R-}, \mathrm{R+}), \tag{2}$$

$$0\mathrm{B}:\ (\mathrm{NS+}, \mathrm{NS+}),\ (\mathrm{NS-}, \mathrm{NS-}),\ (\mathrm{R+}, \mathrm{R+}),\ (\mathrm{R-}, \mathrm{R-}), \tag{3}$$

And additionally, (NS+, NS+) with any *single one* of the 4 possible (R, R) sectors. (4)

If there are two (R, R) sectors, then there must be an acompanying (NS−, NS−) sector due to the closure of OPE. It is straight-foward to check that these possiblilities are all valid under the above constaints: (0) level matching of closed strings, (1) mutual locality, (2) closure of OPE, and (3) (apparent) modular invariance (not sufficient yet, to be checked below).

(a) The torus partition function of the theory breaks up into a product of independent sums over the bosonic X and fermionic $(\psi, \tilde{\psi})$ oscillators. The bosonic part is identical to the bosonic string situation, therefore modular invariant; to check the total partition function for modular invariance, we will look at the fermionic contributions $Z = Z_{\psi,\tilde{\psi}}$ explicitly.

Similar to the Type II case, the building block of Z is given by:

$$Z^{\alpha}{}_{\beta} = \mathrm{Tr}_{\alpha}\left[(-1)^{\beta F} q^{H}\right] \tag{5}$$

Where α, β labels the periodicity in the spatial and temporal directions (σ^1, σ^2); note that for fermionic fields, anti-periodicity in the time direction gives the simple trace, while the periodic path integral gives the trace weighted by $(-1)^F$, as is explained in *Polchinski*, Appendix A.

✒ PAST WORK, AS TEMPLATE ✒

(a) Mode expansion of X CFT is[1]:

$$\partial X(z) = -i\sqrt{\frac{\alpha'}{2}}\sum_{m=-\infty}^{\infty}\frac{\alpha_m}{z^{m+1}}, \quad \bar\partial X(\bar z) = -i\sqrt{\frac{\alpha'}{2}}\sum_{m=-\infty}^{\infty}\frac{\tilde\alpha_m}{\bar z^{m+1}}, \tag{6}$$

$$X = x - i\sqrt{\frac{\alpha'}{2}}\left(\alpha_0 \ln z + \tilde\alpha_0 \ln \bar z\right) + i\sqrt{\frac{\alpha'}{2}}\sum_{m\neq 0}\frac{1}{m}\left(\frac{\alpha_m}{z^m} + \frac{\tilde\alpha_m}{\bar z^m}\right), \tag{7}$$

Momentum p is the charge for *spacetime* translation; we have:

$$X \mapsto X + \text{const}, \quad j_a = \frac{i}{\alpha'}\,\partial_a X, \tag{8}$$

$$p = \frac{1}{2\pi i}\oint_C \left(\mathrm{d}z\, j - \mathrm{d}\bar z\, \tilde j\,\right) = \frac{1}{\alpha'}\sqrt{\frac{\alpha'}{2}}\left(\alpha_0 + \tilde\alpha_0\right) = \sqrt{\frac{1}{2\alpha'}}\left(\alpha_0 + \tilde\alpha_0\right) \tag{9}$$

Additionally, for compact free boson, X is only defined modulo $2\pi R$; therefore, states after $X + 2\pi R$ translation should be identical to the original states, i.e.

$$e^{ip\,(2\pi R)} = \mathbb{1}, \quad p = \frac{n}{R}, \quad n \in \mathbb{Z} \tag{10}$$

This, in fact, holds for any field theory[2] defined for $X \in S^1$, including the ordinary quantum mechanics (a classical field theory) on S^1.

On the other hand, there are additional constraints in string theory: for the state of a *single* closed string, there is a discrete translational symmetry on the *worldsheet*:

$$X(\sigma^1 + 2\pi) \cong X(\sigma^1), \quad X(\sigma^1 + 2\pi) = X(\sigma^1) + 2\pi R w, \quad w \in \mathbb{Z} \tag{11}$$

With some definite winding number w. In $(z, \bar z)$ coordinates, we have:

$$2\pi R w = X\!\left(z\, e^{2\pi i}, \bar z\, e^{-2\pi i}\,\right) - X(z, \bar z) = -i\sqrt{\frac{\alpha'}{2}}\,2\pi i\left(\alpha_0 - \tilde\alpha_0\right) = 2\pi\sqrt{\frac{\alpha'}{2}}\left(\alpha_0 - \tilde\alpha_0\right), \tag{12}$$

$$p = \frac{p_L + p_R}{2}, \quad p_L = \sqrt{\frac{2}{\alpha'}}\,\alpha_0, \quad p_R = \sqrt{\frac{2}{\alpha'}}\,\tilde\alpha_0, \tag{13}$$

$$p_{L,R} = \frac{n}{R} \pm \frac{wR}{\alpha'}\,, \tag{14}$$

$$X = x - i\,\frac{\alpha'}{2}\left(p_L \ln z + p_R \ln \bar z\right) + i\sqrt{\frac{\alpha'}{2}}\sum_{m\neq 0}\frac{1}{m}\left(\frac{\alpha_m}{z^m} + \frac{\tilde\alpha_m}{\bar z^m}\right), \tag{15}$$

For the oscillator expressions for L_0, recall that:

$$T(z) = -\frac{1}{\alpha'} :\!\partial X\, \partial X\!: = \sum_m \frac{L_m}{z^{m+2}}, \tag{16}$$

$$L_{m\neq 0} = \frac{1}{2}\sum_l \alpha_{m-l}\alpha_l, \quad L_0 = \frac{1}{2} :\sum_l \alpha_{-l}\alpha_l : \sim \frac{\alpha' p_L^2}{4} + \sum_{l>0}\alpha_{-l}\alpha_l, \tag{17}$$

[1] Again we follow the convention of *Polchinski*.

[2] Reference: discussions in *Polchinski*, Chapter 8.

The L_0 expression may be off by some normal ordering constant; this ambiguity can be resolved by considering:

$$2L_0\,|0,0;n=w=0\rangle = (L_1L_{-1}-L_{-1}L_1)\,|0,0;p_L=p_R=0\rangle = 0-0=0 \tag{18}$$

Therefore the normal ordering constant is, in fact, trivial, and we have:

$$L_0 = \frac{\alpha' p_L^2}{4} + \sum_{l>0}\alpha_{-l}\alpha_l, \quad \tilde{L}_0 = \frac{\alpha' p_R^2}{4} + \sum_{l>0}\tilde{\alpha}_{-l}\tilde{\alpha}_l, \tag{19}$$

(b) The torus partition function is given by:

$$\langle \mathbb{1}\rangle_{T^2} \equiv Z(\tau=\tau_1+i\tau_2) = \int \mathcal{D}X e^{-S} = \mathrm{Tr}\, e^{-(2\pi\tau_2)H} e^{i\,(2\pi\tau_1)P} \tag{20}$$

Here P generates *worldsheet* translation along σ^1, not to be confused with p which generates *spacetime* translation; with $z=e^{-iw},\ w=\sigma^1+i\sigma^2$,

$$\begin{aligned} T^0{}_1 = \eta^{00}\,(\partial_0\sigma^2)\,T_{21} = -iT_{12} &= -i\left(T_{ww}\,(\partial_1 w)(\partial_2 w) + T_{\bar{w}\bar{w}}\,(\partial_1\bar{w})(\partial_2\bar{w})\right)\\ &= T_{ww}-T_{\bar{w}\bar{w}}\\ &= \left(T_{zz}\,(\partial_w z)^2 + \tfrac{c}{24}\right) - \left(T_{\bar{z}\bar{z}}\,(\partial_{\bar{w}}\bar{z})^2 + \tfrac{\tilde{c}}{24}\right)\\ &= T(z)\,(-iz)^2 - \tilde{T}(\bar{z})\,(+i\bar{z})^2 + \frac{c-\tilde{c}}{24}, \end{aligned} \tag{21}$$

$$\begin{aligned} P = \int \frac{\mathrm{d}\sigma_1}{2\pi}\left(-T^0{}_1\right) &= -\int\frac{\mathrm{d}\sigma_1}{2\pi}\,T(z)\,(-iz)^2 + \int\frac{\mathrm{d}\sigma_1}{2\pi}\,\tilde{T}(\bar{z})\,(+i\bar{z})^2 - \frac{c-\tilde{c}}{24}\\ &= +\oint \frac{\mathrm{d}z}{2\pi\,(-iz)}\,T(z)\,(-iz)^2 + \oint\frac{\mathrm{d}\bar{z}}{2\pi\,(+i\bar{z})}\,\tilde{T}(\bar{z})\,(+i\bar{z})^2 - \frac{c-\tilde{c}}{24}\\ &= \oint\frac{\mathrm{d}z}{2\pi i}\,zT(z) - \oint\frac{\mathrm{d}\bar{z}}{2\pi i}\,\bar{z}\tilde{T}(\bar{z}) - \frac{c-\tilde{c}}{24}\\ &= L_0 - \tilde{L}_0 - \frac{c-\tilde{c}}{24}\\ &= \left(L_0 - \tfrac{c}{24}\right) - \left(\tilde{L}_0 - \tfrac{\tilde{c}}{24}\right),\\ H = \int\frac{\mathrm{d}\sigma_1}{2\pi}\,T^0{}_0 &= \int\frac{\mathrm{d}\sigma_1}{2\pi}\,T_{22}\\ &= L_0 + \tilde{L}_0 - \frac{c+\tilde{c}}{24}\\ &= \left(L_0 - \tfrac{c}{24}\right) + \left(\tilde{L}_0 - \tfrac{\tilde{c}}{24}\right), \end{aligned} \tag{22}$$

Here we've used the fact that $\oint \frac{\mathrm{d}\bar{z}}{\bar{z}} = \oint \frac{\mathrm{d}\bar{z}}{\bar{z}} = 2\pi i$. Therefore,

$$Z(\tau) = \mathrm{Tr}\, e^{-(2\pi\tau_2)H} e^{i\,(2\pi\tau_1)P} = \mathrm{Tr}\, q^{L_0-\frac{c}{24}}\,\bar{q}^{\tilde{L}_0-\frac{\tilde{c}}{24}}, \quad q = e^{2\pi i\tau} \tag{23}$$

Using the expressions in (a), we find that L_0 action on a state $|\psi\rangle$ created by $\alpha_{-l},\tilde{\alpha}_{-l}$ yields the sum of occupation numbers N_l weighted by l:

$$L_0\,|\psi\rangle = \left(\frac{\alpha' k_L^2}{4} + \sum_{l>0} l\cdot N_l\right)|\psi\rangle \tag{24}$$

With $c = \tilde{c} = 1$, we obtain:

$$\begin{aligned} Z(\tau) &= (q\bar{q})^{-\frac{1}{24}} \sum_{n,w} e^{-2\pi\tau_2\alpha' \frac{k_L^2+k_R^2}{4}} e^{2\pi i\tau_1\alpha' \frac{k_L^2-k_R^2}{4}} \sum_{(N_l),(\tilde{N}_l)} q^{\sum_{l>0} l\cdot N_l} \bar{q}^{\sum_{l>0} l\cdot \tilde{N}_l} \\ &= (q\bar{q})^{-\frac{1}{24}} \sum_{n,w} e^{-\pi\tau_2\left(\frac{\alpha' n^2}{R^2}+\frac{w^2R^2}{\alpha'}\right)+2\pi i\tau_1 nw} \sum_{(N_l),(\tilde{N}_l)} \prod_{l>0} q^{l\cdot N_l}\bar{q}^{l\cdot\tilde{N}_l} \\ &= |\eta(\tau)|^{-2} \sum_{n,w} e^{-\pi\tau_2\left(\frac{\alpha' n^2}{R^2}+\frac{w^2R^2}{\alpha'}\right)+2\pi i\tau_1 nw} \end{aligned} \tag{25}$$

We've simplified the contributions from the oscillator modes using $\eta(\tau)$, since they are identical to the oscillator contributions of the non-compact $X \in \mathbb{R}^1$:

$$\begin{aligned} (q\bar{q})^{-\frac{1}{24}} \sum_{(N_l),(\tilde{N}_l)} \prod_{l>0} q^{l\cdot N_l}\bar{q}^{l\cdot\tilde{N}_l} &= (q\bar{q})^{-\frac{1}{24}} \prod_{l>0} \sum_{N_l,\tilde{N}_l=0}^{\infty} q^{l\cdot N_l}\bar{q}^{l\cdot\tilde{N}_l} \\ &= (q\bar{q})^{-\frac{1}{24}} \prod_{l>0} \frac{1}{1-q^l}\frac{1}{1-\bar{q}^l} = |\eta(\tau)|^{-2} \end{aligned} \tag{26}$$

In the $R \to \infty$ limit, only the $w = 0$ modes survive; all other modes are exponentially suppressed by the $e^{-\pi\tau_2 w^2R^2/\alpha'}$ factor; i.e.

$$\begin{aligned} Z(\tau) &= |\eta(\tau)|^{-2} \sum_{n,w} \exp\left\{-\pi\tau_2\left(\frac{\alpha' n^2}{R^2}+\frac{w^2R^2}{\alpha'}\right)+2\pi i\tau_1 nw\right\} \\ &\to |\eta(\tau)|^{-2} \sum_{n} \exp\left\{-\pi\tau_2\,\frac{\alpha' n^2}{R^2}\right\}, \quad k = \frac{n}{R} \\ &\to |\eta(\tau)|^{-2}\, V\!\int \frac{\mathrm{d}k}{2\pi} \exp\left\{-\pi\tau_2\,\alpha' k^2\right\} \\ &= V\,|\eta(\tau)|^{-2}\left(4\pi^2\alpha'\tau_2\right)^{-\frac{1}{2}} \\ &\equiv V \cdot Z_X(\tau) = 2\pi R\, Z_X(\tau) \end{aligned} \tag{27}$$

We recover the partition function $V \cdot Z_X(\tau)$ for non-compact X, as expected.

(c) Using the Poisson resummation formula, we find that:

$$Z(\tau) = 2\pi R\, Z_X(\tau) \sum_{m,w} \exp\left(-\frac{\pi R^2 |m - w\tau|^2}{\alpha'\tau_2}\right) \tag{28}$$

$Z_X(\tau)$ is modular invariant by the properties of the Dedekind $\eta(\tau)$ function, as is demonstrated for the non-compact X in *Polchinski*.

The sum, on the other hand, is naturally invariant under $T\colon \tau \mapsto \tau + 1$, by making a change of variables $m \mapsto m + w$. It is also invariant under $S\colon \tau \mapsto -1/\tau$ with $m \mapsto -w, w \mapsto m$ [3]. Therefore, $Z(\tau)$ is modular invariant.

[3] Reference: *Polchinski*.

$\boxed{2}$ $\mathbb{Z}_2$ Orbifold

The $\mathbb{Z}_2$ orbifold is constructed by imposing an additional identification on $X \in S^1$:

$$X \cong -X \tag{29}$$

The target space is then reduced to $S^1/\mathbb{Z}_2 \cong [0, \pi R]$.

(a) The first contributions to the orbifold partition function comes from the states that are invariant reflection r; we have:

$$\mathrm{Tr}_{S^1/\mathbb{Z}_2} = \mathrm{Tr}_{S^1} \frac{1+r}{2} = \frac{1}{2}\,\mathrm{Tr}_{S^1} + \frac{1}{2}\,\mathrm{Tr}_{S^1} \circ r \tag{30}$$

Acting on $q^{L_0 - \frac{c}{24}}\,\bar{q}^{\tilde{L}_0 - \frac{\tilde{c}}{24}}$, the first term gives $\frac{1}{2}\,Z_{S^1}(\tau)$ where Z_{S^1} is the S^1 partition function we've obtained in $\boxed{1}$.

For the second term, note that:

$$r: \quad \Big|(N_l), (\tilde{N}_l); n, w\Big\rangle \longmapsto (-1)^{\sum_l (N_l + \tilde{N}_l)}\Big|(N_l), (\tilde{N}_l); -n, -w\Big\rangle \tag{31}$$

In particular, it reverses n, w, hence r insertion gives vanishing amplitude unless $n = w = 0$; the summation is very much similar to the Z_{S^1} case, i.e. we have:

$$\begin{aligned} \frac{1}{2}\,\mathrm{Tr}_{S^1}\left(r\, q^{L_0 - \frac{c}{24}}\,\bar{q}^{\tilde{L}_0 - \frac{\tilde{c}}{24}}\right) &= \frac{1}{2}\,(q\bar{q})^{-\frac{1}{24}} \prod_{l>0} \sum_{N_l, \tilde{N}_l = 0}^{\infty} (-1)^{N_l + \tilde{N}_l} q^{l \cdot N_l} \bar{q}^{l \cdot \tilde{N}_l} \\ &= \frac{1}{2}\,(q\bar{q})^{-\frac{1}{24}} \prod_{l>0} \frac{1}{1-(-q^l)}\,\frac{1}{1-(-\bar{q}^l)} = \left|\frac{\eta(\tau)}{\theta_2(\tau)}\right| \end{aligned} \tag{32}$$

Where we've used the fact that[4]: $q^{-\frac{1}{24}} \prod_{l>0} \frac{1}{1-(-q^l)} = \sqrt{2}\sqrt{\frac{\eta(\tau)}{\theta_2(\tau)}}$. Therefore, the total contributions from r–invariant states are:

$$\frac{1}{2}\,Z_{S^1}(\tau) + \left|\frac{\eta(\tau)}{\theta_2(\tau)}\right| \tag{33}$$

(b) With $X \cong -X$, new possibilities emerge as the boundary condition along σ^1:

$$X(\sigma^1 + 2\pi) \cong X(\sigma^1), \quad X(\sigma^1 + 2\pi) = \pm X(\sigma^1) + 2\pi R w, \quad w \in \mathbb{Z} \tag{34}$$

The "$-$" sign corresponds to the *twisted states*. Due to the anti-periodicity, ∂X has a half-integer mode expansion:

$$\partial X\big(z\, e^{2\pi i}\big) = -\partial X(z), \tag{35}$$

$$\partial X(z) = -i\sqrt{\frac{\alpha'}{2}} \sum_{m=-\infty}^{\infty} \frac{\alpha_{m-\frac{1}{2}}}{z^{m+\frac{1}{2}}}, \quad \bar{\partial} X(\bar{z}) = -i\sqrt{\frac{\alpha'}{2}} \sum_{m=-\infty}^{\infty} \frac{\tilde{\alpha}_{m-\frac{1}{2}}}{\bar{z}^{m+\frac{1}{2}}}, \tag{36}$$

$$X = x + i\sqrt{\frac{\alpha'}{2}} \sum_{m=-\infty}^{\infty} \frac{1}{m+\frac{1}{2}}\left(\frac{\alpha_{m+\frac{1}{2}}}{z^{m+\frac{1}{2}}} + \frac{\tilde{\alpha}_{m+\frac{1}{2}}}{\bar{z}^{m+\frac{1}{2}}}\right), \tag{37}$$

Apply the boundary condition on X, and we find that $x = \pi R w'$; however, due to the identification $X + 2\pi R \cong X \cong -X$, there are only two inequivalent choices: $x = 0$ and $x = \pi R$, which correspond to the string localized around either of the two fixed points of the $\mathbb{Z}_2$ action.

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[4] Reference: Blumenhagen & Plauschinn, *Introduction to CFT*, and also *Polchinski*.

Much similar to the case in $\boxed{1}$, we have:

$$\left[\alpha_{\frac{1}{2}+l}, \alpha_{-\frac{1}{2}-l}\right] = \frac{1}{2} + l, \tag{38}$$

$$L_{m\neq 0} = \frac{1}{2}\sum_{l} \alpha_{m-\frac{1}{2}-l}\alpha_{\frac{1}{2}+l}, \quad L_0 = \frac{1}{2} : \sum_{l} \alpha_{-\frac{1}{2}-l}\alpha_{\frac{1}{2}+l} : \sim \sum_{l\geq 0} \alpha_{-\frac{1}{2}-l}\alpha_{\frac{1}{2}+l} \tag{39}$$

We can use the same trick to fix the normal ordering constant in L_0; this time it is non-trivial:

$$L_{-1} = \frac{1}{2}\,\alpha_{-\frac{1}{2}}^2 + \sum_{l\geq 0} \alpha_{-\frac{1}{2}-l}\alpha_{\frac{1}{2}+l}, \quad L_1 = \frac{1}{2}\,\alpha_{\frac{1}{2}}^2 + \sum_{l>0} \alpha_{\frac{1}{2}-l}\alpha_{\frac{1}{2}+l}, \tag{40}$$

$$\begin{aligned} L_0 \left|0, 0; x\right\rangle &= \frac{1}{2}\left(L_1 L_{-1} - L_{-1} L_1\right)\left|0, 0; x\right\rangle \\ &= \frac{1}{2} \times \frac{1}{4}\,\alpha_{\frac{1}{2}}^2 \alpha_{-\frac{1}{2}}^2 \left|0, 0; x\right\rangle - 0 \\ &= \frac{1}{16}\left|0, 0; x\right\rangle, \end{aligned} \tag{41}$$

$$L_0 = \frac{1}{16} + \sum_{l\geq 0} \alpha_{-\frac{1}{2}-l}\alpha_{\frac{1}{2}+l} = \frac{1}{16} + \sum_{l\geq 0}\left(l + \frac{1}{2}\right) N_{l+\frac{1}{2}} = \frac{1}{16} + \sum_{l>0}\left(l - \frac{1}{2}\right) N_{l-\frac{1}{2}}, \tag{42}$$

The trace can then be computed, following the same recipe as before:

$$\begin{aligned} \mathrm{Tr}_{S^1}\left(\frac{1+r}{2}\, q^{L_0 - \frac{c}{24}}\, \bar{q}^{\tilde{L}_0 - \frac{\tilde{c}}{24}}\right) &= (q\bar{q})^{-\frac{1}{24}+\frac{1}{16}} \prod_{l+\frac{1}{2}\in\mathbb{Z}^+} \sum_{N_l, \tilde{N}_l = 0}^{\infty} \frac{1 + (-1)^{N_l + \tilde{N}_l}}{2}\, q^{l\cdot N_l}\, \bar{q}^{l\cdot\tilde{N}_l} \times 2 \\ &= \frac{1}{2}\,(q\bar{q})^{+\frac{1}{48}} \left\{ \prod_{l>0} \left|\frac{1}{1 - q^{l-\frac{1}{2}}}\right|^2 + \prod_{l>0}\left|\frac{1}{1 + q^{l-\frac{1}{2}}}\right|^2 \right\} \times 2 \\ &= \left|\frac{\eta(\tau)}{\theta_4(\tau)}\right| + \left|\frac{\eta(\tau)}{\theta_3(\tau)}\right| \end{aligned} \tag{43}$$

There is an extra factor of 2 from the number of twisted sectors: $x = 0$ and $x = \pi R$.

(c) The full partition function is therefore:

$$Z(\tau) = \frac{1}{2}\, Z_{S^1}(\tau) + \left|\frac{\eta(\tau)}{\theta_2(\tau)}\right| + \left|\frac{\eta(\tau)}{\theta_4(\tau)}\right| + \left|\frac{\eta(\tau)}{\theta_3(\tau)}\right| \tag{44}$$

The first term is modular invariant, as is proved in $\boxed{1}$.

The remaining terms are also modular invariant, due to the transformational properties of η and θ functions[5]:

$$T \circlearrowleft \left|\frac{\eta(\tau)}{\theta_2(\tau)}\right| \overset{S}{\longleftrightarrow} \left|\frac{\eta(\tau)}{\theta_4(\tau)}\right| \overset{T}{\longleftrightarrow} \left|\frac{\eta(\tau)}{\theta_3(\tau)}\right| \circlearrowleft S \tag{45}$$

Therefore, the full partition function is modular invariant.

[5]Reference: *Blumenhagen & Plauschinn.*

$\boxed{3}$ Torus 4-point function in bc CFT

$$\left\langle c(w_1)\, b(w_2)\, \tilde{c}(\bar{w}_3)\, \tilde{b}(\bar{w}_4) \right\rangle = \int \mathscr{D}b\, \mathscr{D}\tilde{b}\, \mathscr{D}c\, \mathscr{D}\tilde{c}\ c(w_1)\, b(w_2)\, \tilde{c}(\bar{w}_3)\, \tilde{b}(\bar{w}_4)\, e^{-S'} \equiv Z' \tag{46}$$

First we argue that only the zero modes of the insertions survive the path integral[6]. In fact, as anti-commuting replacements of the gauge degrees of freedom, ghost modes are *defined* to be the eigenvalues of $P^\dagger P$, where P is the conformal Killing differential[7]. More specifically, given a conformal Killing vector (CKV) $\delta\sigma^a$, the conformal Killing equation can be written as:

$$P\, \delta\sigma = 0 \tag{47}$$

While $P^\dagger \delta' g = 0$ gives moduli variation $\delta' g_{ab}$ of the metric. Roughly speaking, P captures the variation of gauge fixing under an arbitrary gauge transformation; naturally, CKV's are given by $(\ker P)$, while $(\det P) \sim \Delta_{FP}$ is the Faddeev–Popov functional measure near the gauge slice. $(\det P)$ can then be calculated with:

$$\delta\sigma^a \mapsto c^a, \quad \delta' g_{ab} \mapsto b_{ab}, \quad \Delta_{FP} \sim \det P \sim \int \mathscr{D}b\, \mathscr{D}\tilde{b}\, \mathscr{D}c\, \mathscr{D}\tilde{c}\ e^{-S'}, \tag{48}$$

$$S' = \frac{1}{2\pi} \int \mathrm{d}^2\sigma\, g^{1/2}\, b_{ab}\, (P\cdot c)^{ab} = \frac{1}{2\pi} \int \mathrm{d}^2 w\, \left(b\, \bar{\partial}_w c + \tilde{b}\, \partial_w \tilde{c} \right) \tag{49}$$

In the end we have chosen conformal gauge, such that[8] $P \sim (\bar{\partial}_w, \partial_w)$, $P^\dagger P \sim -\bar{\partial}_w \partial_w = -\nabla^2$. In the $w = \sigma^1 + i\sigma^2$ coordinates, CKV's are simple translations: $c^a = \text{const}$; with $z = e^{-iw}$, it gets mapped to $c^z = c^w\, \partial_w z = c^w\, (-iz)$, which agrees with the zero mode c_0 in the $c(z)$ expansion:

$$c(z) = \sum_{m=-\infty}^{\infty} \frac{c_m}{z^{m+1-\lambda}} = c_0\, z + \sum_{m\neq 0} \frac{c_m}{z^{m-1}}, \quad \lambda = 2 \tag{50}$$

Now we are finally ready to prove our argument: for anti-commuting variables like $c(z)$,

$$\int \mathscr{D}c \sim \prod_m \int \mathrm{d}c_m \sim \prod_m \frac{\partial}{\partial c_m} \tag{51}$$

Since c_0 corresponds to a CKV, $P\cdot c_0 = 0$, therefore it vanishes in $S' = \int \mathrm{d}^2\sigma\, (b\cdot P\cdot c)$; for the path integral to be non-zero, there has to be some additional c_0 insertions, i.e.

$$Z' \sim \int \mathscr{D}b\, \mathscr{D}\tilde{b}\, \mathscr{D}c\, \mathscr{D}\tilde{c}\ c_0 b_0 \tilde{c}_0 \tilde{b}_0\, e^{-S'} \sim \left(\frac{1}{\sqrt{\tau_2}}\right)^4 \int \mathscr{D}'b\, \mathscr{D}'\tilde{b}\, \mathscr{D}'c\, \mathscr{D}'\tilde{c}\ e^{-S'}, \quad \int \mathscr{D}'c \sim \prod_{m\neq 0} \int \mathrm{d}c_m \tag{52}$$

Note the additional $\left(\frac{1}{\sqrt{\tau_2}}\right)^4$ factor coming from the zero modes[9]; this has to do with the normalization of the zero modes, each contributing a factor of $\frac{1}{\sqrt{A}}$, where $A \sim \tau_2$ is the volume (surface

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[6] I would like to thank 谷夏 for some very helpful discussions about this problem.

[7] Reference: *Polchinski*, Chapter 3 & 5.

[8] References:

- Nakahara, *Geometry, Topology and Physics*;
- Blumenhagen et al, *Basic Concepts of String Theory*.

[9] Reference: *Di Francesco et al.*

area) of the torus. On a different note, since it is very difficult, if not impossible, to keep track of various (often divergent) constant factors in the path integral, we have been and will be calculating Z' up to an overall constant coefficient.

Now we have to deal with the path integral over non-zero modes. Note that the holomorphic mode expansion (50) is incomplete for our purpose: it gives the *on-shell* mode expansion, while our path integral should go over all possible configurations, including the off-shell modes, which is *not* holomorphic. However, on $T^2 = S^1 \times S^1$, the full modes are simple[10]:

$$-\nabla^2 \psi_{n_1,n_2} = \lambda_{n_1,n_2} \psi_{n_1,n_2}, \tag{53}$$

$$\begin{aligned}\psi_{n_1,n_2} &= \exp\left(i\left(n_1\tilde{\sigma}^1 + n_2\tilde{\sigma}^2\right)\right), \quad \tilde{\sigma}^2 = \frac{\sigma^2}{\tau_2}, \quad \tilde{\sigma}^1 = \sigma^1 - \sigma^2\frac{\tau_1}{\tau_2}, \\ &= \exp\left\{i\left(n_1\sigma^1 + \frac{n_2 - n_1\tau_1}{\tau_2}\,\sigma^2\right)\right\},\end{aligned} \tag{54}$$

Here we first use the "rectangular" coordinates $(\tilde{\sigma}^1, \tilde{\sigma}^2) \in [0, 2\pi]^2$ to write down the obvious eigenfunctions ψ_{n_1,n_2}, and then relate them back to the (σ^1, σ^2) coordinates. Therefore, we have:

$$\begin{aligned}\lambda_{n_1,n_2} &= \left\{n_1^2 + \left(\frac{n_2 - n_1\tau_1}{\tau_2}\right)^2\right\} \\ &= \frac{1}{\tau_2^2}\left\{(n_1\tau_2)^2 + \left(n_1\tau_1 - n_2\right)^2\right\} \\ &= \frac{1}{\tau_2^2}\left|n_1\tau - n_2\right|^2,\end{aligned} \tag{55}$$

$$\det{}' P \sim \left(\prod_{n_1,n_2}{}' \sqrt{\lambda_{n_1,n_2}}\right)^2 \sim \prod_{n_1,n_2}{}' \lambda_{n_1,n_2} \tag{56}$$

The determinant can be computed with ζ-function regularization, as is performed in detail in *Di Francesco*; the result can be nicely summarized using the Eisenstein series, as shown in *Nakahara*:

$$E(\tau, s) = \sum_{n_1,n_2}{}' \frac{\tau_2^s}{\left|n_1\tau - n_2\right|^{2s}}, \tag{57}$$

$$\det{}' P \sim \prod_{n_1,n_2}{}' \frac{1}{\tau_2^2}\left|n_1\tau - n_2\right|^2 \sim \tau_2 \exp\left\{-\partial_s E'(\tau, s)_{s=0}\right\} = \tau_2^2\left|\eta(\tau)\right|^4 \tag{58}$$

Finally, we have:

$$Z' \sim \tau_2^{-2} \det{}' P \sim \tau_2^{-2}\tau_2^2\left|\eta(\tau)\right|^4 \sim \left|\eta(\tau)\right|^4 \tag{59}$$

4 Torus Propagator as a Trace

$$w' \to 0, \quad \left\langle \partial_w X(w)\, \partial_{w'} X(w')\right\rangle = \mathrm{Tr}\left(\partial_w X(w)\, \partial_{w'} X(w')\, q^{L_0 - \frac{c}{24}}\, \bar{q}^{\tilde{L}_0 - \frac{\tilde{c}}{24}}\right) \tag{60}$$

[10] References: (1) *Nakahara*, (2) *Di Francesco et al.*, and (3) http://theory.uchicago.edu/~sethi/Teaching/P483-W2018/p483-sol3.pdf.

Here we've dropped the time ordering in the $w' \to 0$ limit. Recall the mode expansion of ∂X in $\boxed{1}$; we see that only the "diagonal" components of $\partial X(w)\,\partial X(w')$ survive in the trace, i.e.

$$
\begin{aligned}
\partial_w X(w)\,\partial_{w'} X(w') &= (\partial_w z)(\partial_{w'} z')\,\partial_z X(z)\,\partial_{z'} X(z'), \quad z = e^{-iw}, \quad 1 \le |z| \le e^{2\pi\tau_2} \\
&\sim -\frac{\alpha'}{2} \sum_{n=-\infty}^{\infty} \frac{\alpha_{-n}\alpha_n}{z^{-n+1} z'^{n+1}} (-iz)(-iz') \\
&= \frac{\alpha'}{2}\left(\alpha_0^2 + \sum_{n>0} \left(\left(\frac{z}{z'}\right)^n + \left(\frac{z'}{z}\right)^n \right) \alpha_{-n}\alpha_n + \sum_{n>0} n \left(\frac{z'}{z}\right)^n \right) \\
&= \frac{\alpha'}{2}\left(\alpha_0^2 + \sum_{n>0} \left(\left(\frac{z}{z'}\right)^n + \left(\frac{z'}{z}\right)^n \right) \alpha_{-n}\alpha_n + \frac{zz'}{(z-z')^2} \right)
\end{aligned} \tag{61}
$$

The last term is a normal ordering constant; here it is naturally regularized by $\left(\frac{z'}{z}\right)^n$.

The α_0^2 term can be substituted with spacetime momentum p; we have:

$$
p = \sqrt{\frac{1}{2\alpha'}}\,(\alpha_0 + \tilde\alpha_0) = \sqrt{\frac{1}{2\alpha'}}\,2\alpha_0 = \sqrt{\frac{2}{\alpha'}}\,\alpha_0, \tag{62}
$$

$$
\partial_w X(w)\,\partial_{w'} X(w') \sim \frac{\alpha'}{2}\left(\frac{\alpha' p^2}{2} + \sum_{n>0} \left(\left(\frac{z}{z'}\right)^n + \left(\frac{z'}{z}\right)^n \right) nN_n \right) \tag{63}
$$

On the other hand, the partition function is:

$$
\begin{aligned}
Z(\tau) = \langle \mathbb{1} \rangle &= (q\bar q)^{-\frac{1}{24}}\, V \int \frac{\mathrm{d}k}{2\pi}\, e^{-\pi\tau_2\,\alpha' k^2} \sum_{(N_l),(\tilde N_l)} q^{\sum_{l>0} l\cdot N_l}\, \bar q^{\sum_{l>0} l\cdot \tilde N_l} \\
&= (q\bar q)^{-\frac{1}{24}}\, V \int \frac{\mathrm{d}k}{2\pi}\, e^{-\pi\tau_2\,\alpha' k^2} \sum_{(N_l),(\tilde N_l)} \prod_{l>0} q^{l\cdot N_l} \bar q^{l\cdot \tilde N_l} \\
&= |\eta(\tau)|^{-2}\, V \int \frac{\mathrm{d}k}{2\pi}\, e^{-\pi\tau_2\,\alpha' k^2}
\end{aligned} \tag{64}
$$

We can work out $Z^{-1}\langle \partial X \partial X \rangle$ by considering term by term insertion of the $\partial X \partial X$ mode expansion into the above expression. For the $\frac{\alpha' p^2}{2}$ term, we have a contribution of:

$$
\frac{\displaystyle\int \frac{\mathrm{d}k}{2\pi}\, \frac{\alpha' k^2}{2}\, e^{-\pi\tau_2\,\alpha' k^2}}{\displaystyle\int \frac{\mathrm{d}k}{2\pi}\, e^{-\pi\tau_2\,\alpha' k^2}} = \frac{\alpha'}{2}\,\frac{1}{2\cdot \pi\alpha'\tau_2} = \frac{1}{4\pi\tau_2} \tag{65}
$$

For the nN_n insertion, we have a contribution of:

$$
\begin{aligned}
\frac{\displaystyle\sum_{(N_l)} nN_n\, q^{\sum_{l>0} l\cdot N_l}}{\displaystyle\sum_{(N_l)} q^{\sum_{l>0} l\cdot N_l}} &= \frac{\displaystyle\sum_{(N_l)} nN_n \prod_{l>0} q^{l\cdot N_l}}{\displaystyle\sum_{(N_l)} \prod_{l>0} q^{l\cdot N_l}} = \frac{\displaystyle\sum_{N_n=0}^{\infty} nN_n\, q^{n\cdot N_n}}{\displaystyle\sum_{N_n=0}^{\infty} q^{n\cdot N_n}} = \frac{nq^n \dfrac{\partial}{\partial(q^n)} \displaystyle\sum_{N_n=0}^{\infty} q^{n\cdot N_n}}{\displaystyle\sum_{N_n=0}^{\infty} q^{n\cdot N_n}} \\
&= \frac{nq^n \frac{\partial}{\partial(q^n)} \frac{1}{1-q^n}}{\frac{1}{1-q^n}} = \frac{nq^n}{1-q^n}
\end{aligned} \tag{66}
$$

Therefore, the complete result is given by:

$$
\begin{aligned}
\frac{1}{Z(\tau)}\langle \partial_w X(w)\,\partial_{w'}X(w')\rangle &= \frac{\alpha'}{2}\left(\frac{1}{4\pi\tau_2} + \sum_{n>0}\left(\left(\frac{z}{z'}\right)^n + \left(\frac{z'}{z}\right)^n\right)\frac{nq^n}{1-q^n} + \frac{zz'}{(z-z')^2}\right) \\
&\xrightarrow[z'\to 1]{w'\to 0} \frac{\alpha'}{2}\left(\frac{1}{4\pi\tau_2} + \sum_{n>0}\left(z^n + z^{-n}\right)\frac{nq^n}{1-q^n} + \frac{z}{(z-1)^2}\right)
\end{aligned} \tag{67}
$$

On the other hand, the torus propagator is given by:

$$
G'(w,\bar{w};w',\bar{w}') = -\frac{\alpha'}{2}\,\ln|f(w-w',\tau)|^2 + \frac{\alpha'}{4\pi\tau_2}\left(\mathrm{Im}\,(w-w')\right)^2, \tag{68}
$$

$$
f(w,\tau) \equiv \theta_1\left(\frac{w}{2\pi}\,\middle|\,\tau\right) = 2\,e^{\frac{i\pi\tau}{4}}\sin\frac{w}{2}\prod_{m>0}^{\infty}(1-q^m)(1-z^{-1}q^m)(1-zq^m), \quad z = e^{-iw} \tag{69}
$$

We find that $\partial_w\partial_{w'}G'$ contains the same zero mode contribution $\frac{\alpha'}{8\pi\tau_2}$ and normal ordering contribution $\frac{\alpha'}{2}\frac{z}{(z-1)^2}$ as in (67):

$$
\partial_w\partial_{w'}G'(w,\bar{w};w',\bar{w}')_{w'=0} = \frac{\alpha'}{8\pi\tau_2} + \frac{\alpha'}{2}\partial_w^2\ln f(w,\tau), \tag{70}
$$

$$
\partial_w^2\ln f(w,\tau) = \partial_w^2\ln\sin\frac{w}{2} + \partial_w^2\sum_{m>0}\Big(\ln(1-zq^m) + \ln(1-z^{-1}q^m)\Big), \tag{71}
$$

$$
\partial_w^2\ln\sin\frac{w}{2} = \partial_w^2\ln\sin\frac{w}{2} = -\frac{1}{4\sin^2\frac{w}{2}} = \frac{1}{2\,(\cos w - 1)} = \frac{1}{z+z^{-1}-2} = \frac{z}{(z-1)^2}, \tag{72}
$$

The remaining parts come from oscillator modes; they also match with (67), but the equivalence is less obvious: we have[11]:

$$
\begin{aligned}
\partial_w^2\sum_{m>0}\ln(1-zq^m) &= \partial_w^2\sum_{m>0}\sum_{n>0}-\frac{1}{n}\,(zq^m)^n \\
&= \sum_{n>0}\partial_w^2\left(-\frac{1}{n}\,z^n\right)\sum_{m>0}q^{mn}, \quad \partial_w = -iz\,\partial_z \\
&= \sum_{n>0}-\frac{(-in)^2}{n}\,z^n\cdot\frac{q^n}{1-q^n} \\
&= \sum_{n>0}z^n\frac{nq^n}{1-q^n},
\end{aligned} \tag{73}
$$

$$
\partial_w^2\sum_{m>0}\ln(1-z^{-1}q^m) = \sum_{n>0}z^{-n}\frac{nq^n}{1-q^n}, \tag{74}
$$

This is precisely the contribution from oscillator modes in (67). Therefore, we have:

$$
\frac{1}{Z(\tau)}\langle \partial_w X(w)\,\partial_{w'}X(w')\rangle_{w'=0} = \partial_w\partial_{w'}G'(w,\bar{w};w',\bar{w}')_{w'=0} \tag{75}
$$

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[11] Reference: http://theory.uchicago.edu/~sethi/Teaching/P483-W2018/p483-sol3.pdf. I would like to thank Lucy Smith for providing this hint.