

## 1 Read *Polchinski* Sections 1.3 and 1.4:

Read, *mostly* understood. □

## 2 Spinning Closed String in AdS Space:

For a classical spinning string, we have Nambu–Goto action:

$$S_{NG} = -T \int d\tau d\sigma \sqrt{-\det \gamma_{ab}}, \quad \gamma_{ab} = G_{\mu\nu} \partial_a X^\mu \partial_b X^\nu \quad (1)$$

Here  $G_{\mu\nu}$  is the spacetime metric.  $\gamma_{ab}$  can be treated as the induced metric on the worldsheet.

In AdS space we have:

$$ds^2 = R^2 (-\cosh^2 \rho dt^2 + d\rho^2 + \sinh^2 \rho d\Omega^2) \quad (2)$$

Where  $d\Omega^2$  is the metric of a unit  $(d-2)$ -sphere  $S^{d-2}$ . For convenience let's define unit  $S^{d-2}$  metric  $G_{ij}^1$ , and raise or lower the  $i, j, \dots$  indices using  $G_{ij}^1$  instead of  $G_{ij}$ , i.e.,

$$G_{ij}^1 = G_{ij} / (R^2 \sinh^2 \rho), \quad i, j = 2, \dots, d-1 \quad (3)$$

Furthermore, we consider the special case that the closed string is *folded*, like a rubber band stretched along a line; in this case we can choose the worldsheet parameter  $(\tau, \sigma) = (t, \rho)$  while  $\Omega = \Omega(t, \rho) = \Omega(\tau, \sigma)$ , which leads to the following decomposition:

$$\partial_a X^\mu = \delta_a^\mu + \delta_i^\mu \partial_a \Omega^i, \quad a = 0, 1, \quad i = 2, \dots, d-1, \quad (4)$$

$$\begin{aligned} \gamma_{ab} &= G_{\mu\nu} \partial_a X^\mu \partial_b X^\nu \\ &= G_{ab} + G_{ij} \partial_a \Omega^i \partial_b \Omega^j \\ &= G_{ab} + R^2 \sinh^2 \rho G_{ij}^1 \partial_a \Omega^i \partial_b \Omega^j \\ &= R^2 \left\{ \begin{pmatrix} -\cosh^2 \rho & \\ & 1 \end{pmatrix} + \sinh^2 \rho \begin{pmatrix} (\partial_a \Omega)^2 & \partial_a \Omega \cdot \partial_b \Omega \\ \partial_b \Omega \cdot \partial_a \Omega & (\partial_b \Omega)^2 \end{pmatrix} \right\} \end{aligned} \quad (5)$$

Here  $\partial_a \Omega \cdot \partial_b \Omega \equiv \partial_a \Omega^i \partial_b \Omega_i \equiv G_{ij}^1 \partial_a \Omega^i \partial_b \Omega^j$ , and we have:

$$\begin{aligned} \det \gamma_{ab} &= (R^2)^2 \left\{ \sinh^4 \rho \det (\partial_a \Omega^i \partial_b \Omega_i) \right. \\ &\quad \left. + \sinh^2 \rho ((\partial_a \Omega)^2 - (\partial_b \Omega)^2 \cosh^2 \rho) \right. \\ &\quad \left. - \cosh^2 \rho \right\}, \quad (6) \\ \sqrt{-\det \gamma_{ab}} &= R^2 \left\{ \cosh^2 \rho - \sinh^2 \rho ((\partial_a \Omega)^2 - (\partial_b \Omega)^2 \cosh^2 \rho) \right. \\ &\quad \left. - \sinh^4 \rho \det (\partial_a \Omega^i \partial_b \Omega_i) \right\}^{1/2} \end{aligned}$$

Mark the end points of the string with  $\rho = r(t)$ , then the total length of such closed folded string is  $\ell = 4r$ . We then have:

$$S = -4TR^2 \int dt \int_0^r d\rho \sqrt{\cosh^2 \rho - \sinh^2 \rho ((\partial_a \Omega)^2 - (\partial_b \Omega)^2 \cosh^2 \rho) - \sinh^4 \rho \det (\partial_a \Omega^i \partial_b \Omega_i)} \quad (7)$$

Further simplification comes from the fact that, due to rotational symmetry, the string's motion can be restricted in a plane where its position is characterized by some angle  $\theta = \Omega^{i_0} \in \{\Omega^i\}_i$ . In this case other angle parameters  $\Omega^i|_{i \neq i_0} = 0$ , and the action is further reduced to:

$$S = -4TR^2 \int dt \int_0^r d\rho \sqrt{\cosh^2 \rho - \sinh^2 \rho ((\partial_a \theta)^2 - (\partial_b \theta)^2 \cosh^2 \rho)} = \int dt \int_0^r d\rho \mathcal{L}, \quad (8)$$

$$\mathcal{L} = -4TR^2 \sqrt{\cosh^2 \rho - \omega^2 \sinh^2 \rho}, \quad \omega = \partial_t \theta, \partial_\rho \theta = 0 \quad (9)$$

We consider the special solution  $\theta = \omega t$ , while in general the endpoint  $r = r(t)$  could be dynamical; variation of the action w.r.t.  $r(t)$  gives<sup>1</sup>:

$$0 = \delta S = -4TR^2 \int dt \int_r^{r+\delta r} d\rho \sqrt{\cosh^2 \rho - \omega^2 \sinh^2 \rho} = -4TR^2 \int dt \sqrt{\cosh^2 r - \omega^2 \sinh^2 r} \delta r, \quad (10)$$

$$\omega^2 = \frac{\cosh^2 r}{\sinh^2 r} = \coth^2 r \quad (11)$$

Note that if  $\omega$  is constant, then  $r$  must be fixed by (11). Taking  $\theta$  as the only dynamical variable, it is then straight-forward to write the energy  $E$  and angular momentum  $J$  for such folded closed string:

$$\omega = \dot{\theta}, \quad \Pi = \frac{\partial \mathcal{L}}{\partial \omega} = 4TR^2 \frac{\omega \sinh^2 \rho}{\sqrt{\cosh^2 \rho - \omega^2 \sinh^2 \rho}}, \quad (12)$$

$$J = \int_0^r d\rho \Pi = 4TR^2 \int_0^r d\rho \frac{\omega \sinh^2 \rho}{\sqrt{\cosh^2 \rho - \omega^2 \sinh^2 \rho}}, \quad (13)$$

$$E = \int_0^r d\rho (\Pi \omega - \mathcal{L}) = 4TR^2 \int_0^r d\rho \frac{\cosh^2 \rho}{\sqrt{\cosh^2 \rho - \omega^2 \sinh^2 \rho}}, \quad (14)$$

In the large string limit,  $r \rightarrow \infty$ ,  $\omega = \coth r \rightarrow 1$ . Expand in terms of  $\epsilon = \omega - 1 > 0$ , we find that  $r = \frac{1}{2} \ln \left(1 + \frac{2}{\epsilon}\right) \sim \frac{1}{2} \ln \frac{2}{\epsilon}$ , or alternatively,  $e^{2r} \cdot \epsilon \sim 2$ . With some help from Mathematica<sup>TM</sup>, we get:

$$\begin{aligned} E - J &= 4TR^2 \int_0^r d\rho \frac{\cosh^2 \rho - \omega \sinh^2 \rho}{\sqrt{\cosh^2 \rho - \omega^2 \sinh^2 \rho}} = 4TR^2 \int_0^r d\rho \left(1 + \frac{\epsilon^2}{8} \sinh^2(2\rho) + \mathcal{O}(\epsilon^3)\right) \\ &= 4TR^2 \left(r \left(1 - \frac{\epsilon^2}{16} + \mathcal{O}(\epsilon^3)\right) + \mathcal{O}(1)\right) = \left(2TR^2 \ln \frac{2}{\epsilon}\right) \left(1 - \frac{\epsilon^2}{16} + \mathcal{O}(\epsilon^3)\right) \\ &\sim 2TR^2 \left(\ln \frac{2}{\epsilon}\right) \end{aligned} \quad (15)$$

Similarly,  $J \sim 4TR^2 \int_0^r d\rho \sinh^2 \rho \sim TR^2 \left(\frac{2}{\epsilon}\right)$ , this gives:

$$E - J \sim 2TR^2 \ln \frac{J}{TR^2} \quad (16)$$

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<sup>1</sup>The above reasoning is confirmed in e.g. [arXiv:hep-th/0204051](https://arxiv.org/abs/hep-th/0204051).

### 3 Special Conformal Transformations:

$$x^\mu \xrightarrow{K(a)} \tilde{x}^\mu = \frac{x^\mu + x^2 a^\mu}{1 + 2a \cdot x + a^2 x^2} \quad (17)$$

(a) Under special conformal transformation  $K(a)$ , metric  $\delta_{\mu\nu} \mapsto g_{\mu\nu}$  while:

$$g_{\alpha\beta} d\tilde{x}^\alpha d\tilde{x}^\beta = \delta_{\mu\nu} dx^\mu dx^\nu, \quad g_{\alpha\beta} = \delta_{\mu\nu} \frac{\partial x^\mu}{\partial \tilde{x}^\alpha} \frac{\partial x^\nu}{\partial \tilde{x}^\beta} \quad (18)$$

To calculate this we have to know the inverse transformation  $x = K^{-1}(a) \tilde{x}$ . First, notice the following decomposition<sup>2</sup> of  $K(a)$ :

$$\tilde{x}^\mu = \frac{\frac{x^\mu}{x^2} + a^\mu}{\frac{1}{x^2} + \frac{2a \cdot x}{x^2} + a^2} = \frac{\frac{x^\mu}{x^2} + a^\mu}{\left| \frac{x^\mu}{x^2} + a^\mu \right|^2}, \quad (19)$$

$$\text{i.e. } K(a): x^\mu \xrightarrow{I} \frac{x^\mu}{x^2} \xrightarrow{T(a)} y^\mu = \frac{x^\mu}{x^2} + a^\mu \xrightarrow{I} \tilde{x}^\mu = \frac{y^\mu}{y^2}, \quad (20)$$

$$\text{i.e. } \frac{\tilde{x}^\mu}{\tilde{x}^2} = \frac{y^\mu}{y^2} \Big/ \frac{1}{y^2} = y^\mu = \frac{x^\mu}{x^2} + a^\mu \quad (21)$$

From (21), we see that the transformation parameter  $a^\mu$  composes linearly:  $K(b)K(a) = K(a+b)$ , therefore  $K^{-1}(a) = K(-a)$ , and we have:

$$x^\mu = K(-a) \tilde{x}^\mu = \frac{\tilde{x}^\mu - \tilde{x}^2 a^\mu}{1 - 2a \cdot \tilde{x} + a^2 \tilde{x}^2} = \frac{\tilde{y}^\mu}{y^2}, \quad (22)$$

$$\begin{aligned} \frac{\partial x^\mu}{\partial \tilde{x}^\alpha} &= \frac{\partial x^\mu}{\partial \tilde{y}^\sigma} \frac{\partial \tilde{y}^\sigma}{\partial \tilde{x}^\alpha} = \left( \frac{\partial}{\partial \tilde{y}^\sigma} \frac{\tilde{y}^\mu}{\tilde{y}^2} \right) \frac{\partial}{\partial \tilde{x}^\alpha} \left( \frac{\tilde{x}^\sigma}{\tilde{x}^2} - a^\sigma \right) = \left( \frac{\partial}{\partial \tilde{y}^\sigma} \frac{\tilde{y}^\mu}{\tilde{y}^2} \right) \left( \frac{\partial}{\partial \tilde{x}^\alpha} \frac{\tilde{x}^\sigma}{\tilde{x}^2} \right) \\ &= (\tilde{y}^2 \delta_\sigma^\mu - 2\tilde{y}^\mu \tilde{y}_\sigma) (\tilde{x}^2 \delta_\alpha^\sigma - 2\tilde{x}^\sigma \tilde{x}_\alpha) / (\tilde{y}^4 \tilde{x}^4), \end{aligned} \quad (23)$$

$$\begin{aligned} g_{\alpha\beta} &\stackrel{(18)}{=} \delta_{\mu\nu} (\tilde{y}^2 \delta_\sigma^\mu - 2\tilde{y}^\mu \tilde{y}_\sigma) (\tilde{x}^2 \delta_\alpha^\sigma - 2\tilde{x}^\sigma \tilde{x}_\alpha) (\tilde{y}^2 \delta_\rho^\nu - 2\tilde{y}^\nu \tilde{y}_\rho) (\tilde{x}^2 \delta_\beta^\rho - 2\tilde{x}^\rho \tilde{x}_\beta) / (\tilde{y}^8 \tilde{x}^8) \\ &\stackrel{\Sigma_{\mu,\nu}}{=} \tilde{y}^{-4} \delta_{\sigma\rho} (\tilde{x}^2 \delta_\alpha^\sigma - 2\tilde{x}^\sigma \tilde{x}_\alpha) (\tilde{x}^2 \delta_\beta^\rho - 2\tilde{x}^\rho \tilde{x}_\beta) / \tilde{x}^8 \\ &\stackrel{\Sigma_{\sigma,\rho}}{=} \tilde{y}^{-4} \tilde{x}^{-4} \delta_{\alpha\beta} \end{aligned} \quad (24)$$

We see that  $g_{\alpha\beta} = f(x) \delta_{\alpha\beta}$ , with coefficient:

$$f(x) = \tilde{y}^{-4} \tilde{x}^{-4} \stackrel{(20)}{=} \frac{x^4}{\tilde{x}^4} \stackrel{(21)}{=} (1 + 2a \cdot x + a^2 x^2)^2 \quad (25)$$

□<sub>(a)</sub>

(b) In 2D with  $z = x^1 + ix^2$ ,  $x^\mu \sim (z, \bar{z})$ , we see from (21) that:

$$\frac{x^\mu}{x^2} \sim \frac{z}{|z|^2} = \frac{1}{\bar{z}} \mapsto \frac{1}{\bar{z}} + a, \quad \text{i.e. } z \mapsto w = \frac{1}{\frac{1}{\bar{z}} + a} = \frac{z}{1 + z\bar{a}} \quad (26)$$

Expand in the  $\bar{a} \rightarrow 0$  limit, we find that  $w = z(1 - z\bar{a} - \dots) \sim z - z^2\bar{a}$ , i.e. it is generated by:

$$K_{\bar{z}} = -z^2 \partial_z = -z^2 \partial, \quad \partial \equiv \partial_z \quad (27)$$

<sup>2</sup>See *Di Francesco et al*, and also [github.com/davidsd/ph229](https://github.com/davidsd/ph229).

Note that when considering non-holomorphic functions, we have to consider  $(z, \bar{z})$  as *two* independent variables; hence the anti-holomorphic transformation  $\bar{z} \mapsto \bar{w} = \frac{\bar{z}}{1+\bar{z}a} \sim \bar{z} - \bar{z}^2 a$  provides another degree of freedom, namely:

$$K_\mu \sim (K_{\bar{z}} = -z^2 \partial, K_z = -\bar{z}^2 \bar{\partial}), \quad (28)$$

$$\partial \equiv \partial_z, \quad \bar{\partial} \equiv \partial_{\bar{z}}$$

Similarly, for translation  $z \mapsto z + a$  and its conjugate, we have  $P_\mu \sim (P_z = \partial, P_{\bar{z}} = \bar{\partial})$ . However, dilation and rotation are both encoded in a complex rescaling  $z \mapsto \lambda z$ ,  $\lambda = re^{i\theta} \in \mathbb{C}$ ; we have:

$$z \mapsto \lambda z, \quad \lambda = re^{i\theta} \in \mathbb{C}, \quad \begin{aligned} \delta r &\longleftrightarrow D = z \partial + \bar{z} \bar{\partial}, \\ \delta \theta &\longleftrightarrow M = i(z \partial - \bar{z} \bar{\partial}), \end{aligned} \quad (29)$$

In summary, we have  $\text{span}_{\mathbb{R}} \{P_\mu, K_\mu, D, M\} = \mathfrak{so}(3, 1)$  generating the “global” transformation subgroup of the 2D conformal group; here, the  $\mathfrak{so}(3, 1)$  boost is a linear combination<sup>3</sup> of  $P_\mu$  and  $K_\mu$ . More specifically, in 2D any holomorphic or anti-holomorphic function gives a conformal transformation, hence the (classical) 2D conformal group is generated by:

$$\ell_m = z^{m+1} \partial, \quad \bar{\ell}_m = \bar{z}^{m+1} \bar{\partial}, \quad m \in \mathbb{Z} \quad (30)$$

i.e. the *Witt algebra* (or Virasoro algebra  $\mathbf{Vir}_c$  with  $c = 0$ ). It is clear that a (complexified)  $\mathfrak{so}(3, 1)$  lives inside  $\mathbf{Vir}_c$ , i.e.,

$$\begin{aligned} \mathfrak{so}(3, 1)^\mathbb{C} &= \text{span}_{\mathbb{C}} \{P_\mu, K_\mu, D, M\} \\ &= \text{span}_{\mathbb{C}} \{\ell_m, \bar{\ell}_m \mid m = 0, \pm 1\} = \mathfrak{sl}(2, \mathbb{R})^\mathbb{C} \oplus_{\mathbb{C}} \mathfrak{sl}(2, \mathbb{R})^\mathbb{C} \subset \mathbf{Vir}_c \end{aligned} \quad (31)$$

■

**4** *bc CFT*:

$$S = \frac{1}{2\pi} \int d^2 z \, b \bar{\partial} c \quad (32)$$

Stress tensor of a theory can be obtained via variation over the metric, or equivalently, over the fields  $\phi^i$  with  $\delta\phi$  induced by some *local* spacetime translation  $x^\mu \mapsto x^\mu + \delta x^\mu$ ,  $\delta x^\mu = \epsilon(x) a^\mu$ . Here  $\epsilon(x)$  is any compactly supported bump function, centered around some point  $x_0$ .

In 2D, we have  $\mu = z, \bar{z}$ ; for  $\phi(z, \bar{z})$  with conformal weight  $(h, \bar{h})$ , consider  $z \mapsto z'$ ,  $\bar{z} \mapsto \bar{z}'$ . For convenience, let's first consider a generic variation  $\delta z = \epsilon(z, \bar{z})$  before restricting to spacetime translation; we have:

$$\phi'(z', \bar{z}') = \left(\frac{dz'}{dz}\right)^{-h} \left(\frac{d\bar{z}'}{d\bar{z}}\right)^{-\bar{h}} \phi(z, \bar{z}), \quad (33)$$

$$\tilde{\delta}\phi = (-h \partial\epsilon - \bar{h} \bar{\partial}\bar{\epsilon}) \phi, \quad (34)$$

$$\delta\phi = \tilde{\delta}\phi - \frac{\partial\phi}{\partial x^\mu} \delta x^\mu = (-h \partial\epsilon - \bar{h} \bar{\partial}\bar{\epsilon}) \phi - \epsilon \partial\phi - \bar{\epsilon} \bar{\partial}\phi, \quad (35)$$

Here we use  $\tilde{\delta}\phi$  to denote the “internal” variation related to the conformal weights.

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<sup>3</sup>See e.g. [github.com/davidsd/ph229](https://github.com/davidsd/ph229).

Note that  $\phi = b, c$  are anti-commuting Grassmann numbers, variation of the action gives:

$$\begin{aligned}\delta S[b, c, \delta b, \delta c] &= \frac{1}{2\pi} \int d^2 z (\delta b \bar{\partial} c + b \bar{\partial} \delta c) \\ &= \frac{1}{2\pi} \int d^2 z (-\bar{\partial} c \delta b - \bar{\partial} b \delta c) + \frac{1}{2\pi} \int d^2 z \bar{\partial}(b \delta c)\end{aligned}\quad (36)$$

For *unknown*  $b, c$  and arbitrary  $\delta b, \delta c$ , the second term is reduced to a boundary term at infinity and can be dropped; imposing  $\delta S = 0$  gives the equation of motion (EOM):  $\bar{\partial} b = \bar{\partial} c = 0$ .

On the other hand, for *on-shell*  $b, c$  and compactly supported  $\phi = \delta b, \delta c$  given in (35), the first term in (36) vanishes while  $\delta S_0 = 0$  still holds; this gives:

$$\begin{aligned}0 = \delta S_0 &= \frac{1}{2\pi} \int d^2 z \bar{\partial}(b \delta c) = \frac{1}{2\pi} \int d^2 z \bar{\partial}(-(1-\lambda)bc\partial\epsilon - b\partial c\epsilon) \\ &= \frac{1}{2\pi} \int d^2 z (-(1-\lambda)bc\bar{\partial}\partial\epsilon - b\partial c\bar{\partial}\epsilon)\end{aligned}\quad (37)$$

Here we've distributed the  $\bar{\partial}$  operator and dropped all terms that vanish automatically by EOM. Next we shall collect the  $\partial\epsilon, \bar{\partial}\epsilon$  terms; integrating by parts on the first integrand gives:

$$\begin{aligned}0 = \delta S_0 &= \frac{1}{2\pi} \int d^2 z ((1-\lambda)\partial(bc) - b\partial c)\bar{\partial}\epsilon \\ &= \frac{1}{2\pi} \int d^2 z ((\partial b)c - \lambda\partial(bc))\bar{\partial}\epsilon \\ &= -\frac{1}{2\pi} \int d^2 z \epsilon(z, \bar{z}) \partial_{\bar{z}}((\partial b)c - \lambda\partial(bc))\end{aligned}\quad (38)$$

Notice that we have obtained a conserved current using a generic  $\delta z = \epsilon(z, \bar{z}), \delta \bar{z} = \bar{\epsilon}(z, \bar{z})$ ; by setting  $\epsilon = \epsilon(z)$ , we get a energy momentum tensor<sup>4</sup>:

$$T(z) = :(\partial b)c: - \lambda\partial(:bc:)\quad (39)$$

Normal ordering is added manually to remove singular terms.

To compute  $TT$  OPE, we need the OPE of  $b(z)c(0)$ ; this is obtained by examining the following path integral, which is zero since the integrand is a total functional derivative:

$$0 = \int \mathcal{D}b \mathcal{D}c \frac{\delta}{\delta \phi} (e^{-S} \psi)\quad (40)$$

Taking  $\phi, \psi = b, c$ , this generates operator equations such as  $\bar{\partial} b(z)c(0) = 2\pi\delta^2(z, \bar{z})$ . Note that  $\bar{\partial}(\frac{1}{z}) = 2\pi\delta^2(z, \bar{z})$ , which gives:

$$b(z)c(0) \sim c(z)b(0) \sim \frac{1}{z}, \quad b(z)b(0) \sim 0 \sim c(z)c(0)\quad (41)$$

With the  $bc$  OPE in hand, the  $TT$  OPE is computed directly with brute force, by repeatedly applying Wick's theorem. This gives:

$$T(z)T(0) \sim \frac{-6\lambda^2 + 6\lambda - 1}{z^4} + \dots\quad (42)$$

In general we have  $-6\lambda^2 + 6\lambda - 1 = \frac{c}{2}$ ; for  $\lambda = 2$  this gives  $c = -26$ . ■

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<sup>4</sup>Note that the energy momentum tensor obtained in this way is generally *not* unique: it can be off by a boundary term or any term that vanish on-shell; see Luboš' comment at [physics.stackexchange.com/a/96100](https://physics.stackexchange.com/a/96100), also [arXiv:1601.03616](https://arxiv.org/abs/1601.03616). However, it is possible to fix this redundancy by considering  $Tb$  OPE and match its conformal dimension. I would like to thank 林般 for pointing this out.

**5 Free Fermion CFT:**

$$S = \int d^2z \psi_i \bar{\partial} \psi^i, \quad \psi^i = \psi_i^*, \quad \psi_i = \psi_i(z) \quad (43)$$

(a) Mode expansion of such chiral fermion is given by:

$$\psi_i = \sum_{k \in \mathbb{Z} + \frac{1}{2}} \frac{b_{ik}}{z^{k+\frac{1}{2}}}, \quad b_{ik} = \frac{1}{2\pi i} \oint dz z^{k-\frac{1}{2}} \psi_i \quad (44)$$

Canonical quantization is achieved by simply imposing anti-commutation relations; this is justified by mapping the system onto a cylinder, then  $b_{ik}$ 's indeed map to modes on the spatial circle<sup>5</sup>. The only non-zero commutators are:

$$\{b_{ik}, b_q^{j\dagger}\} = \delta_{k+q,0} \delta_i^j \quad (45)$$

This gives the only non-zero 2-point functions:

$$\begin{aligned} \langle \psi_i(z) \psi^j(w) \rangle &= \sum_{k,q \in \mathbb{Z} + \frac{1}{2}} \frac{1}{z^{k+\frac{1}{2}}} \frac{1}{w^{q+\frac{1}{2}}} \langle b_{ik} b_q^{j\dagger} \rangle \\ &= \sum_{k,q \in \mathbb{Z} + \frac{1}{2}} \frac{1}{z^{k+\frac{1}{2}}} \frac{1}{w^{q+\frac{1}{2}}} \langle 0 | \{b_{ik}, b_q^{j\dagger}\} | 0 \rangle = \frac{\delta_i^j}{z-w} \end{aligned} \quad (46)$$

Note that  $b_k^i |0\rangle = 0, \forall k \geq \frac{1}{2}$ .

(b)(c) Combining two  $\psi$  expansions gives the mode expansion of  $J_i^j = :\psi_i(z) \psi^j(z):$ , namely:

$$J_i^j(z) = \sum_{k \in \mathbb{Z}} \frac{(J_i^j)_k}{z^{k+1}}, \quad (J_i^j)_k = \sum_{q \in \mathbb{Z} + \frac{1}{2}} :b_{iq} b_{k-q}^{j\dagger}: \quad (47)$$

It is in fact more convenient to obtain the  $JJ$  OPE first, and then use it to find the  $[J_0, J_0]$  mode commutator<sup>6</sup>; note that  $\psi_i(z) \psi^j(w)$  contraction gives  $\frac{\delta_i^j}{z-w}$ , we have:

$$J_i^j(z) J_k^l(w) \sim \frac{\delta_i^l \delta_k^j}{z^2} + \frac{\delta_k^j J_i^l(w) - \delta_i^l J_k^j(w)}{z}, \quad (48)$$

$$[(J_i^j)_0, (J_k^l)_0] = \frac{1}{(2\pi i)^2} \oint_0 dw \oint_w dz J_i^j(z) J_k^l(w) = \delta_i^l (J_k^j)_0 - \delta_k^j (J_i^l)_0 \quad (49)$$

(d) Similar to  $bc$  CFT, we have:

$$T(z) = \frac{1}{2} (: \psi_i \partial \psi^i : - : \partial \psi_i \psi^i :), \quad T(z) T(w) \sim \frac{n/2}{(z-w)^4} + \frac{2T(w)}{(z-w)^2} + \frac{\partial T(w)}{z-w} \quad (50)$$

With each (complex) field contributing  $\frac{1}{2} \times 2$  central charge<sup>7</sup>.

<sup>5</sup>This can be proven rigorously by considering operator equations like in the  $bc$  CFT problem.

<sup>6</sup>I would like to thank 谷夏 for providing this hint.

<sup>7</sup>In fact a complex (Dirac) fermion can be “treated like” (*dual to*) a boson; this is *bosonization*.

(e) For real fermions, there is an additional reality condition:

$$\psi^i = \psi_i^\star = \psi_i \quad (51)$$

The canonical quantization still holds without the extra adjoint, same as the 2-point function:

$$\langle \psi_i(z) \psi_j(w) \rangle = \frac{\delta_{ij}}{z - w} \quad (52)$$

Similar holds for  $J_{ij} = :\psi_i \psi_j:$  and its OPE, but we no longer need to distinguish upper/lower indices; we have:

$$J_{ij}(z) J_{kl}(0) \sim \frac{-\delta_{ik}\delta_{jl} + \delta_{il}\delta_{jk}}{z^2} + \frac{-\delta_{ik}J_{jl}(0) + \delta_{il}J_{jk}(0) + \delta_{jk}J_{il}(0) - \delta_{jl}J_{ik}(0)}{z} \quad (53)$$

$$[(J_{ij})_0, (J_{kl})_0] = -\delta_{ik}(J_{jl})_0 + \delta_{il}(J_{jk})_0 + \delta_{jk}(J_{il})_0 - \delta_{jl}(J_{ik})_0 \quad (54)$$

This is precisely the  $\mathfrak{o}(n)$  algebra. ■