

Homework II

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Problem 1

1.A)

Let M be our $D > 2$ dimensional C^∞ manifold and $\phi : \mathbb{R} \times M \rightarrow M$ be a one parameter family of diffeomorphisms, that is, $\forall t \in \mathbb{R} \mid \phi_t : M \rightarrow M$ is a diffeomorphism such that,

- $\forall p \in M \mid \phi_0(p) = p$
- $\forall p \in M; \forall t, s \in \mathbb{R} \mid \phi_{t+s}(p) = (\phi_t \circ \phi_s)(p)$
- $\forall p \in M \mid \phi(p) : \mathbb{R} \rightarrow M$ is at least C^1

Then this family of diffeomorphisms define in a natural manner a vector which generate these transformations, we define at each point $p \in M$ this vector by it's action in a function $f : M \rightarrow \mathbb{R}$,

$$\xi_p(f) = \left. \frac{d}{dt}((f \circ \phi_t)(p)) \right|_{t=0}$$

And from this we define ξ as a vector field in M , this vector field has as integral curves exactly ϕ . Let's open a little bit more in some chart $x : M \rightarrow \mathbb{R}^D$,

$$\begin{aligned} \xi_p(f) &= \left. \frac{d}{dt}((f \circ x^{-1} \circ x \circ \phi_t)(p)) \right|_{t=0} \\ \xi_p(f) &= \left. \frac{d}{dt}((f \circ x^{-1}) \circ (x \circ \phi_t)(p)) \right|_{t=0} \\ \xi_p(f) &= \partial_\mu(f \circ x^{-1}) \left|_{x \circ \phi_0(p)} \frac{d}{dt}((x \circ \phi_t)^\mu(p)) \right|_{t=0} \\ \xi_p^\mu \partial_\mu(f \circ x^{-1}) \left|_{x(p)} \right. &= \partial_\mu(f \circ x^{-1}) \left|_{x(p)} \frac{d}{dt}((x \circ \phi_t)^\mu(p)) \right|_{t=0} \end{aligned}$$

Where of course ∂_μ is to be interpreted as the derivative of the μ -th component in the chart x . Here we have a clear definition of the values of the ξ vector field in a chart x ,

$$\xi_p^\mu = \left. \frac{d}{dt}((x \circ \phi_t)^\mu(p)) \right|_{t=0}$$

The term inside the derivative is just the pullback of the chart x — the chart can be seen as a \mathbb{R}^D -valued function —, which in it's own can be seen as a new chart x'_t defined by the transformations of the diffeomorphism family ϕ , that is,

$$x'_t = \phi_t^* x = x \circ \phi_t : M \rightarrow \mathbb{R}^D$$

All of this is consistent with our interpretation of the diffeomorphisms being a ‘*coordinate change*’, in principle, with enough derivability of ϕ we can actually write,

$$\begin{aligned} x'_t &= x'_0 + t \left. \frac{d}{dt}(x'_t) \right|_{t=0} + \mathcal{O}(t^2) \\ x_1'^\mu &=: x'^\mu = x^\mu + \xi^\mu + \dots \end{aligned}$$

We just restored the index to not confuse the components of the vector field ξ in the basis x with the vector field itself. That is, we showed that the transformation done by ϕ_1 is equivalent to a ‘*infinitesimal coordinate change*’ by ξ^μ . Actually, all this we did is the special case of a more general type of derivative, the Lie Derivative, given a vector field ξ and it’s family of integral curves ϕ , it’s defined in terms of the pushforward of the object under analysis,

$$\mathcal{L}_\xi T = \frac{d}{dt}(\phi_{-t*}T) \Big|_{t=0}$$

For a $(0, 2)$ tensor, that is, for the metric,

$$\mathcal{L}_\xi g_{\mu\nu} = \nabla_\xi g_{\mu\nu} + g_{\mu\alpha} \nabla_\nu \xi^\alpha + g_{\alpha\nu} \nabla_\mu \xi^\alpha$$

And as the connection is metric compatible,

$$\mathcal{L}_\xi g_{\mu\nu} = \nabla_\nu \xi_\mu + \nabla_\mu \xi_\nu = 2\nabla_{(\mu} \xi_{\nu)}$$

This amounts for the first term in an expansion in t of the diffeomorphism transformed metric $\phi_{-t*}g = g'_{t\mu\nu}$, that is,

$$\begin{aligned} \phi_{-t*}g_{\mu\nu} &=: g'_{t\mu\nu} = g_{\mu\nu} + t\mathcal{L}_\xi g_{\mu\nu} + \mathcal{O}(t^2) \\ g'_{t\mu\nu} &= g_{\mu\nu} + 2t\nabla_{(\mu} \xi_{\nu)} + \mathcal{O}(t^2) \end{aligned}$$

Which also can be interpreted as the ‘*infinitesimal transformation*’ of the metric. Imposing that the initial and transformed metric are conformally flat,

$$\begin{aligned} \exp(2\omega'_t)\eta_{\mu\nu} &= \exp(2\omega)\eta_{\mu\nu} + 2t\nabla_{(\mu} \xi_{\nu)} + \mathcal{O}(t^2) \\ \exp\left(2\omega'_0 + 2t\frac{d}{dt}(\omega'_t) \Big|_{t=0} + \mathcal{O}(t^2)\right)\eta_{\mu\nu} &= \exp(2\omega)\eta_{\mu\nu} + 2t\nabla_{(\mu} \xi_{\nu)} + \mathcal{O}(t^2) \\ \exp(2\omega)\exp\left(2t\frac{d}{dt}(\omega'_t) \Big|_{t=0} + \mathcal{O}(t^2)\right)\eta_{\mu\nu} &= \exp(2\omega)\eta_{\mu\nu} + 2t\nabla_{(\mu} \xi_{\nu)} + \mathcal{O}(t^2) \\ \exp(2\omega)\left(1 + 2t\frac{d}{dt}(\omega'_t) \Big|_{t=0} + \mathcal{O}(t^2)\right)\eta_{\mu\nu} &= \exp(2\omega)\eta_{\mu\nu} + 2t\nabla_{(\mu} \xi_{\nu)} + \mathcal{O}(t^2) \\ \exp(2\omega)2t\frac{d}{dt}(\omega'_t) \Big|_{t=0} \eta_{\mu\nu} &= 2t\nabla_{(\mu} \xi_{\nu)} \\ \frac{d}{dt}(\omega'_t) \Big|_{t=0} g_{\mu\nu} &= \nabla_{(\mu} \xi_{\nu)} \end{aligned}$$

The term $\frac{d}{dt}(\omega'_t) \Big|_{t=0}$ is fully determined by ξ , to see this just contract both sides with the metric,

$$\begin{aligned} \frac{d}{dt}(\omega'_t) \Big|_{t=0} g^{\mu\nu} g_{\mu\nu} &= g^{\mu\nu} \nabla_{(\mu} \xi_{\nu)} \\ \frac{d}{dt}(\omega'_t) \Big|_{t=0} D &= \nabla_\mu \xi^\mu \end{aligned}$$

Substituting this back in the original equation,

$$\nabla_\alpha \xi^\alpha g_{\mu\nu} = D \nabla_{(\mu} \xi_{\nu)} \quad (1.1)$$

This is the condition upon ξ that ensures the diffeomorphism maintain the conformally flatness of the metric. To solve it in it's full generality is hard, so, we'll make some use of symmetry, first, let's write it again in a coordinate free form,

$$\text{Div}_g(\xi)g = \frac{D}{2} \mathcal{L}_\xi g \quad (1.2)$$

Where $\text{Div}_g(\xi)$ is just $\nabla_\alpha \xi^\alpha$, but with the connection ∇ defined with respect with the metric g , now, suppose we have a function $f : M \rightarrow \mathbb{R}$, define $\tilde{g} = \exp(2f)g$, thus,

$$\begin{aligned} \mathcal{L}_\xi \tilde{g} &= \mathcal{L}_\xi(\exp(2f)g) \\ \mathcal{L}_\xi \tilde{g} &= \exp(2f) \mathcal{L}_\xi g + g \mathcal{L}_\xi \exp(2f) \\ \mathcal{L}_\xi \tilde{g} &= \exp(2f) \mathcal{L}_\xi g + g \xi(\exp(2f)) \\ \mathcal{L}_\xi \tilde{g} &= \exp(2f) \mathcal{L}_\xi g + g \xi(f) \exp(2f) \\ \mathcal{L}_\xi \tilde{g} &= \exp(2f) \mathcal{L}_\xi g + \tilde{g} \xi(2f) \end{aligned} \quad (1.3)$$

Also,

$$\begin{aligned} \text{Div}_g(\xi) &= \partial_\alpha \xi^\alpha + \Gamma_{\alpha\lambda}^\alpha \xi^\lambda \\ \text{Div}_g(\xi) &= \partial_\alpha \xi^\alpha + \frac{1}{2} g^{\alpha\beta} (\partial_\alpha g_{\beta\lambda} + \partial_\lambda g_{\beta\alpha} - \partial_\beta g_{\alpha\lambda}) \xi^\lambda \\ \text{Div}_g(\xi) &= \partial_\alpha \xi^\alpha + g^{\alpha\beta} (\partial_\alpha \omega g_{\beta\lambda} + \partial_\lambda \omega g_{\beta\alpha} - \partial_\beta \omega g_{\alpha\lambda}) \xi^\lambda \\ \text{Div}_g(\xi) &= \partial_\alpha \xi^\alpha + (\partial_\lambda \omega + D \partial_\lambda \omega - \partial_\lambda \omega) \xi^\lambda \\ \text{Div}_g(\xi) &= \partial_\alpha \xi^\alpha + D \partial_\lambda \omega \xi^\lambda \end{aligned}$$

And as $\tilde{g} = \exp(2f)g = \exp(2(f + w))\eta$,

$$\begin{aligned} \text{Div}_{\tilde{g}}(\xi) &= \partial_\alpha \xi^\alpha + D \partial_\lambda (f + \omega) \xi^\lambda \\ \text{Div}_{\tilde{g}}(\xi) &= \partial_\alpha \xi^\alpha + D \partial_\lambda \omega \xi^\lambda + D \xi^\lambda \partial_\lambda f \\ \text{Div}_{\tilde{g}}(\xi) &= \text{Div}_g(\xi) + D \xi(f) \end{aligned} \quad (1.4)$$

Multiplying 1.4 by g and subtracting 1.3,

$$\begin{aligned} \text{Div}_{\tilde{g}}(\xi) \tilde{g} - \frac{D}{2} \mathcal{L}_\xi \tilde{g} &= \text{Div}_g(\xi) \tilde{g} + D \tilde{g} \xi(f) - \frac{D}{2} \exp(2f) \mathcal{L}_\xi g - \frac{D}{2} \tilde{g} \xi(2f) \\ \text{Div}_{\tilde{g}}(\xi) \tilde{g} - \frac{D}{2} \mathcal{L}_\xi \tilde{g} &= \exp(2f) \text{Div}_g(\xi) g - \exp(2f) \frac{D}{2} \mathcal{L}_\xi g + D \tilde{g} \xi(f) - D \tilde{g} \xi(f) \\ \text{Div}_{\tilde{g}}(\xi) \tilde{g} - \frac{D}{2} \mathcal{L}_\xi \tilde{g} &= \exp(2f) \left[\text{Div}_g(\xi) g - \frac{D}{2} \mathcal{L}_\xi g \right] \end{aligned}$$

That is, as long as f is sufficiently well behaved, we have,

$$\text{Div}_g(\xi)g = \frac{D}{2}\mathcal{L}_\xi g \Leftrightarrow \text{Div}_{\tilde{g}}(\xi)\tilde{g} = \frac{D}{2}\mathcal{L}_\xi \tilde{g} \quad (1.5)$$

In other words, the vector field ξ which generates diffeomorphisms that preserves the conformally flat condition, does not depend on ω of our conformally flat metric $g = \exp(2\omega)\eta$, so, we can choose $f = -\omega$, such that $\tilde{g} = \eta$, and by 1.5, we just have to solve for,

$$\partial_\alpha \xi^\alpha \eta_{\mu\nu} = D\partial_{(\mu} \xi_{\nu)} \quad (1.6)$$

Which is a lot easier then 1.1, first, we apply ∂^ν to the both sides, relabel the index, apply ∂_ν , symmetrize $\mu \leftrightarrow \nu$ and use 1.6,

$$\begin{aligned} \partial^\nu \partial_\alpha \xi^\alpha \eta_{\mu\nu} &= D\partial^\nu \partial_{(\mu} \xi_{\nu)} \\ \frac{2}{D}\partial_\mu \partial_\alpha \xi^\alpha &= \partial^\nu \partial_\mu \xi_\nu + \partial^\nu \partial_\nu \xi_\mu \\ \frac{2}{D}\partial_\mu \partial_\alpha \xi^\alpha &= \partial^\alpha \partial_\mu \xi_\alpha + \partial^\alpha \partial_\alpha \xi_\mu \\ \frac{2}{D}\partial_\nu \partial_\mu \partial_\alpha \xi^\alpha &= \partial_\nu \partial^\alpha \partial_\mu \xi_\alpha + \partial_\nu \partial^\alpha \partial_\alpha \xi_\mu \\ \frac{2}{D}\partial_{(\nu} \partial_{\mu)} \partial_\alpha \xi^\alpha &= \partial_{(\nu} \partial_{\mu)} \partial^\alpha \xi_\alpha + \partial^\alpha \partial_\alpha \partial_{(\nu} \xi_{\mu)} \\ \frac{2}{D}\partial_\nu \partial_\mu \partial_\alpha \xi^\alpha &= \partial_\nu \partial_\mu \partial_\alpha \xi^\alpha + \frac{\eta_{\mu\nu}}{D}\partial^\alpha \partial_\alpha \partial_\beta \xi^\beta \\ (2-D)\partial_\nu \partial_\mu \partial_\alpha \xi^\alpha &= \eta_{\mu\nu} \partial^\alpha \partial_\alpha \partial_\beta \xi^\beta \end{aligned} \quad (1.7)$$

Contracting with $\eta^{\mu\nu}$ gives,

$$\begin{aligned} (2-D)\partial^\mu \partial_\mu \partial_\alpha \xi^\alpha &= D\partial^\alpha \partial_\alpha \partial_\beta \xi^\beta \\ 2(1-D)\partial^\mu \partial_\mu \partial_\alpha \xi^\alpha &= 0 \\ \partial^\mu \partial_\mu \partial_\alpha \xi^\alpha &= 0 \end{aligned} \quad (1.8)$$

The last step is justified by merely $D \geq 2$, so that up to now, we haven't fully used the hypothesis of being in $D > 2$. So, we'll invoke it now, by using 1.8 in equation 1.7, this is only justified if $D \neq 2$, because, for $D = 2$ the left-hand side of 1.7 is identically zero, but, for $D > 2$ the use of 1.8 in 1.7 results in another constraint,

$$\begin{aligned} (2-D)\partial_\mu \partial_\nu \partial_\alpha \xi^\alpha &= 0 \\ \partial_\mu \partial_\nu \partial_\alpha \xi^\alpha &= 0 \end{aligned} \quad (1.9)$$

It's clear that equation 1.7 allows for much more solutions than 1.9, for example, $\partial \cdot \xi = \cos(k_\mu x^\mu)$ with $k_\mu k^\mu = 0$ is a solution for 1.8, but, isn't for 1.9. To integrate 1.9 is also a lot easier than 1.8,

$$\begin{aligned} \partial_\mu \partial_\nu \partial_\alpha \xi^\alpha &= 0 \\ \partial_\nu \partial_\alpha \xi^\alpha &= a_\nu \end{aligned}$$

$$\partial_\alpha \xi^\alpha = a_\nu x^\nu + b \quad (1.10)$$

With a_ν, b arbitrary constants, this is also automatically a solution of 1.8. To integrate 1.10 is not trivial, as there might be some divergenceless current contributing to ξ^α , that is, the integration gives,

$$\begin{aligned} \partial_\alpha \xi^\alpha &= a_\nu x^\nu + b \\ \xi^\alpha &= \frac{1}{2} a^\alpha_{\mu\nu} x^\mu x^\nu + b^\alpha_\mu x^\mu + c^\alpha + f^\alpha \end{aligned} \quad (1.11)$$

With of course $a^\alpha_{\mu\nu} = a^\alpha_{(\mu\nu)}$, $a^\alpha_{\alpha\nu} = a_\nu$, $b^\alpha_\alpha = b$, c^α being constants, with exception of f^α , which is some non constant divergenceless vector field $\partial_\alpha f^\alpha = 0$. To know what choice of f^α is the correct one we have to go back to equation 1.6, and apply to it ∂_λ ,

$$\frac{2}{D} \partial_\lambda \partial_\alpha \xi^\alpha \eta_{\mu\nu} = \partial_\lambda \partial_\mu \xi_\nu + \partial_\lambda \partial_\nu \xi_\mu$$

Now we permute the index,

$$\begin{aligned} \frac{2}{D} \partial_\lambda \partial_\alpha \xi^\alpha \eta_{\mu\nu} &= \partial_\lambda \partial_\mu \xi_\nu + \partial_\nu \partial_\lambda \xi_\mu \\ \frac{2}{D} \partial_\mu \partial_\alpha \xi^\alpha \eta_{\nu\lambda} &= \partial_\mu \partial_\nu \xi_\lambda + \partial_\lambda \partial_\mu \xi_\nu \\ \frac{2}{D} \partial_\nu \partial_\alpha \xi^\alpha \eta_{\lambda\mu} &= \partial_\nu \partial_\lambda \xi_\mu + \partial_\mu \partial_\nu \xi_\lambda \end{aligned}$$

Sum the two first equations and subtract the third one,

$$\frac{2}{D} (\partial_\lambda \partial_\alpha \xi^\alpha \eta_{\mu\nu} + \partial_\mu \partial_\alpha \xi^\alpha \eta_{\nu\lambda} - \partial_\nu \partial_\alpha \xi^\alpha \eta_{\lambda\mu}) = \partial_\lambda \partial_\mu \xi_\nu + \partial_\nu \partial_\lambda \xi_\mu + \partial_\mu \partial_\nu \xi_\lambda + \partial_\lambda \partial_\mu \xi_\nu - \partial_\nu \partial_\lambda \xi_\mu - \partial_\mu \partial_\nu \xi_\lambda$$

By use of 1.10,

$$\frac{2}{D} (a_\lambda \eta_{\mu\nu} + a_\mu \eta_{\nu\lambda} - a_\nu \eta_{\lambda\mu}) = 2 \partial_\lambda \partial_\mu \xi_\nu$$

And substituting 1.11 in the right-hand side,

$$\frac{2}{D} (a_\lambda \eta_{\mu\nu} + a_\mu \eta_{\nu\lambda} - a_\nu \eta_{\lambda\mu}) = 2 a_{\nu\mu\lambda} + 2 \partial_\lambda \partial_\mu f_\nu$$

We can thus use this to fix f_λ , it's trivial to carry out the integration,

$$\begin{aligned} \partial_\lambda \partial_\mu f_\nu &= \frac{1}{D} (a_\lambda \eta_{\mu\nu} + a_\mu \eta_{\nu\lambda} - a_\nu \eta_{\lambda\mu}) - a_{\nu\mu\lambda} \\ \partial_\mu f_\nu &= \frac{1}{D} (a \cdot x \eta_{\mu\nu} + a_\mu x_\nu - a_\nu x_\mu) - a_{\nu\mu\lambda} x^\lambda + A_{\nu\mu} \\ f_\nu &= \frac{1}{D} \left(a \cdot x x_\nu - \frac{1}{2} a_\nu x \cdot x \right) - \frac{1}{2} a_{\nu\mu\lambda} x^\lambda x^\mu + A_{\nu\mu} x^\mu + B_\nu \\ \partial_\mu f^\mu &= \frac{1}{D} (a \cdot x D + a \cdot x - a \cdot x) - a \cdot x + A^\mu_\mu = A^\mu_\mu \end{aligned} \quad (1.12)$$

Where we just computed also the divergence, which implies the constraint $A^\mu{}_\mu = 0$. Substituting this back in 1.11 get us,

$$\begin{aligned}\xi^\alpha &= \frac{1}{2}a^\alpha{}_{\mu\nu}x^\mu x^\nu + b^\alpha{}_\mu x^\mu + c^\alpha + \frac{1}{D}\left(a \cdot xx^\alpha - \frac{1}{2}a^\alpha x \cdot x\right) - \frac{1}{2}a^\alpha{}_{\mu\lambda}x^\lambda x^\mu + A^\alpha{}_\mu x^\mu + B^\alpha \\ \xi^\alpha &= \frac{1}{D}\left(a \cdot xx^\alpha - \frac{1}{2}a^\alpha x \cdot x\right) + (A^\alpha{}_\mu + b^\alpha{}_\mu)x^\mu + (B^\alpha + c^\alpha) \\ \xi^\alpha &= \frac{1}{D}\left(a \cdot xx^\alpha - \frac{1}{2}a^\alpha x \cdot x\right) + b^\alpha{}_\mu x^\mu + c^\alpha\end{aligned}$$

In the last line we just redefined the tensors. All of this with $a^\alpha, b^\alpha{}_\mu, c^\alpha$ arbitrary constants, but, we didn't really confirmed this is solution of 1.9 is a fully compatible with 1.6, we just checked it satisfy some derived equations from 1.6, now, let's put it to the real test,

$$\begin{aligned}\partial_\alpha \xi^\alpha \eta_{\mu\nu} &= D\partial_{(\mu} \xi_{\nu)} \\ \left[\frac{1}{D}\left(a \cdot xD + a \cdot x - 2\frac{1}{2}a \cdot x\right) + b^\alpha{}_\alpha\right] \eta_{\mu\nu} &= D\left[\frac{1}{D}(a \cdot x\eta_{\mu\nu} + a_{(\mu}x_{\nu)} - a_{(\nu}x_{\mu)}) + b_{(\nu\mu)}\right] \\ a \cdot x\eta_{\mu\nu} + b^\alpha{}_\alpha \eta_{\mu\nu} &= a \cdot x\eta_{\mu\nu} + Db_{(\nu\mu)} \\ b^\alpha{}_\alpha \eta_{\mu\nu} &= Db_{(\nu\mu)} \\ b_{(\nu\mu)} &= \frac{1}{D}b^\alpha{}_\alpha \eta_{\mu\nu}\end{aligned}$$

This constrains the symmetric part of $b_{\mu\nu}$ being a pure trace, so that the degrees of freedom reduces from $b_{\mu\nu}$ to $b_{[\mu\nu]}, b^\alpha{}_\alpha$. This is not the end! We haven't show all the solutions we found are in fact all the solutions from 1.6, as we obtained them from a different approach than direct integration of the equation 1.6, so let's show this, first, 1.6 is a set of first order partial differential equations, there are D^2 such equation, but, they're symmetric in exchange $\mu \leftrightarrow \nu$, so we have to account only for $\frac{D^2+D}{2}$ of them, still, we have to account for the possible constants which can be added to ξ without changing the equation, D of them, and also for another boundary condition on the divergence, 1, so that the full number of constants is,

$$\frac{D^2 + D}{2} + D + 1 = \frac{1}{2}(D + 1)(D + 2)$$

Now let's count the number of constants in the solution we found, namely,

$$\xi_\alpha = \frac{1}{D}\left(a \cdot xx_\alpha - \frac{1}{2}a_\alpha x \cdot x\right) + b_{[\alpha\mu]}x^\mu + \frac{1}{D}b^\mu{}_\mu x_\alpha + c_\alpha$$

Both c^α, a^α contributes with D , $b^\mu{}_\mu$ contributes with 1, and lastly $b_{[\alpha\mu]}$ contributes with $\frac{D^2-D}{2}$, hence, the total number of constants is,

$$\frac{D^2 - D}{2} + 2D + 1 = \frac{1}{2}(D + 1)(D + 2)$$

Matching exactly the number from the original equation, that is, we already got all the solutions, hence, the most general solution of 1.6, for $D > 2$, is given by, just for better reading we make a change of variables that does not affect at all the number of free parameters,

$$\xi_\alpha = 2a \cdot xx_\alpha - a_\alpha x \cdot x + b_{[\alpha\mu]}x^\mu + x_\alpha d + c_\alpha$$

1.B)

We have already counted for the number of independent parameters in the last item as an way of making sure we got all the solutions to the equation, so now it remains to get what is the geometric picture of the transformations ϕ generated by this ξ , first, notice that we have four different kinds of terms in our expression for ξ ,

$$\xi_\alpha = 2a \cdot xx_\alpha - a_\alpha x \cdot x + b_{[\alpha\mu]}x^\mu + x_\alpha d + c_\alpha$$

A quadratic, a two linear and a constant in x . We'll treat one by one, first, let us discuss what is to get ϕ from ξ . As we pointed before,

$$\xi_p^\mu = \left. \frac{d}{dt}((x \circ \phi_t)^\mu(p)) \right|_{t=0}$$

We can deform this definition a little by,

$$\begin{aligned} \xi_{\phi_s(p)}^\mu \Big|_{s=0} &= \left. \frac{d}{dt}((x \circ \phi_t)^\mu(\phi_s(p))) \right|_{t=0} \Big|_{s=0} \\ \xi_{\phi_s(p)}^\mu &= \left. \frac{d}{dt}((x \circ \phi_{t+s})^\mu(p)) \right|_{t=0} \\ \xi_{\phi_s(p)}^\mu &= \left. \frac{d}{d(t+s)}((x \circ \phi_{t+s})^\mu(p)) \right|_{t+s=s} \frac{d(t+s)}{dt} \Big|_{t=0} \\ \xi_{\phi_s(p)}^\mu &= \frac{d}{ds}((x \circ \phi_s)^\mu(p)) \end{aligned}$$

Seeing $\xi_{\phi_s(p)}$ as a function of the coordinates of the point $\phi_s(p)$,

$$\begin{aligned} \xi^\mu((x \circ \phi_s)(p)) &= \frac{d}{ds}((x \circ \phi_s)^\mu(p)) \\ \xi^\mu(x \circ \phi_s) &= \frac{d}{ds}(x \circ \phi_s)^\mu \\ \xi^\mu(x'_s) &= \frac{d}{ds}(x'_s)^\mu \end{aligned}$$

That is, to solve for ϕ , we just have to solve this first order differential equation for the coordinates of $\phi_s(p)$, and as we had already defined, $x \circ \phi_s(p) := x'_s$, so that, we have to solve for,

$$\frac{d}{ds}x'_{s\alpha} = \dot{x}'_{s\alpha} = 2a \cdot x'_s x'_{s\alpha} - a_\alpha x'_s \cdot x'_s + b_{[\alpha\mu]}x'^\mu_s + x'_{s\alpha} d + c_\alpha$$

In this form it's obvious that it is too hard to solve, so we rather solve for each parameter individually, starting by,

- c^α

$$\begin{aligned}\frac{d}{ds}x'_{s\alpha} &= c_\alpha \\ x'_{s\alpha} &= sc_\alpha + x'_{0\alpha} \\ x'_{s\alpha} &= sc_\alpha + x_\alpha\end{aligned}$$

In other words, the transformation parametrized by c^α describe a translation of the coordinate x^α ,

- d

$$\begin{aligned}\frac{d}{ds}x'_{s\alpha} &= x'_{s\alpha}d \\ x'_{s\alpha} &= A_\alpha \exp(sd)\end{aligned}$$

Imposing $x'_{0\alpha} = x_\alpha$,

$$x'_{s\alpha} = x_\alpha \exp(sd)$$

Which accounts for dilatations — rescaling — of the coordinate x_α ,

- $b_{[\alpha\mu]}$

$$\begin{aligned}\frac{d}{ds}x'_{s\alpha} &= b_{[\alpha\mu]}x_s'^\mu \\ x'_{s\alpha} &= \exp\left(s\frac{1}{2}b_{[\rho\sigma]}M^{[\rho\sigma]}\right)_{\alpha\mu} x^\mu\end{aligned}$$

This solution may seem taken directly from a hat, but, let's record that this equation we're solving is a set of first order linear ordinary differential equations, and the solution to one of the kind $\dot{\mathbf{x}}(s) = \mathbf{A} \cdot \mathbf{x}(s)$ is always of the form $\mathbf{x}(s) = \exp(s\mathbf{A})\mathbf{x}(0)$, which is exactly what we wrote, apart from introducing a new tensor M and a $\frac{1}{2}$ factor, why you may ask? Let us be honest, we all know what is happening here... But, we have explicit written just two of the four index in M , the other two stands outside the exponential, M is not at all a free parameter, it has to satisfy the consistency condition,

$$\begin{aligned}\frac{d}{ds}x'_{s\alpha} &= \frac{1}{2}b_{[\rho\sigma]}M^{[\rho\sigma]}_{\alpha}{}^{\nu} \exp\left(s\frac{1}{2}b_{[\rho\sigma]}M^{[\rho\sigma]}\right)_{\nu\mu} x^\mu \\ \frac{d}{ds}x'_{s\alpha} &= \frac{1}{2}b_{[\rho\sigma]}M^{[\rho\sigma]}_{\alpha}{}^{\nu} x'_{s\nu} \\ \frac{d}{ds}x'_{s\alpha} &= \frac{1}{2}b_{[\rho\sigma]}M^{[\rho\sigma]}_{\alpha\nu} x_s'^\nu \Rightarrow \frac{1}{2}b_{[\rho\sigma]}M^{[\rho\sigma]}_{\alpha\nu} = b_{[\alpha\nu]}\end{aligned}$$

This is easy to solve,

$$\begin{aligned}
\frac{1}{2}b_{[\rho\sigma]}M^{[\rho\sigma]}_{\alpha\nu} &= b_{[\alpha\nu]} \\
\frac{1}{2}b_{[\rho\sigma]}M^{[\rho\sigma]}_{\alpha\nu} &= \frac{1}{2}b_{[\rho\sigma]}(\eta_\alpha{}^\rho\eta_\nu{}^\sigma - \eta_\nu{}^\rho\eta_\alpha{}^\sigma) \\
M^{[\rho\sigma]}_{[\alpha\nu]} &= (\eta_\alpha{}^\rho\eta_\nu{}^\sigma - \eta_\nu{}^\rho\eta_\alpha{}^\sigma)
\end{aligned}$$

It's clear that these transformations are just the usual Lorentz ones, boosts plus rotations, as M is exactly the generators of such transformations — in a real form —, just for completeness, we compute,

$$\begin{aligned}
x_s'^\alpha x_{s\alpha}' &= \exp\left(s\frac{1}{2}b_{[\rho\sigma]}M^{[\rho\sigma]}\right)^{[\alpha\nu]} x_\nu \exp\left(s\frac{1}{2}b_{[\rho\sigma]}M^{[\rho\sigma]}\right)_{[\alpha\mu]} x^\mu \\
x_s'^\alpha x_{s\alpha}' &= \exp\left(-s\frac{1}{2}b_{[\rho\sigma]}M^{[\rho\sigma]}\right)^{\nu[\alpha]} x_\nu \exp\left(s\frac{1}{2}b_{[\rho\sigma]}M^{[\rho\sigma]}\right)_{[\alpha\mu]} x^\mu \\
x_s'^\alpha x_{s\alpha}' &= x_\nu \exp\left(-s\frac{1}{2}b_{[\rho\sigma]}M^{[\rho\sigma]}\right)^{[\nu\alpha]} \exp\left(s\frac{1}{2}b_{[\rho\sigma]}M^{[\rho\sigma]}\right)_{[\alpha\mu]} x^\mu
\end{aligned}$$

As trivially $b_{[\rho\sigma]}M^{[\rho\sigma]}$ commute with itself — everything commutes with itself, as long as it doesn't have characteristic 2 —, so that BCH gives trivially,

$$\begin{aligned}
x_s'^\alpha x_{s\alpha}' &= x_\nu \exp\left(-s\frac{1}{2}b_{[\rho\sigma]}M^{[\rho\sigma]} + s\frac{1}{2}b_{[\rho\sigma]}M^{[\rho\sigma]}\right)^\nu_\mu x^\mu \\
x_s'^\alpha x_{s\alpha}' &= x_\nu \exp(\mathbf{0})^\nu_\mu x^\mu \\
x_s'^\alpha x_{s\alpha}' &= x_\nu \eta^\nu_\mu x^\mu \\
x_s'^\alpha x_{s\alpha}' &= x_\mu x^\mu
\end{aligned}$$

Just exactly what a Lorentz transformation should do. Then, the $b_{[\alpha\mu]}$ generates the Lorentz transformations.

- a_α

$$\dot{x}'_{s\alpha} = 2a \cdot x'_s x'_{s\alpha} - a_\alpha x'_s \cdot x'_s \quad (1.13)$$

This is the last, hardest and most interesting one to solve, we proceed with the usual magic hat tricks and divine inspiration/intervention,

$$\begin{aligned}
\frac{\dot{x}'_{s\alpha}}{x'_s \cdot x'_s} &= \frac{2(a \cdot x'_s)x'_{s\alpha}}{x'_s \cdot x'_s} - a_\alpha \\
\frac{\dot{x}'_{s\alpha}}{x'_s \cdot x'_s} - \frac{2(a \cdot x'_s)x'_{s\alpha}}{x'_s \cdot x'_s} &= -a_\alpha \\
\frac{\dot{x}'_{s\alpha}}{x'_s \cdot x'_s} - \frac{2(a \cdot x'_s)(x'_s \cdot x'_s)x'_{s\alpha}}{(x'_s \cdot x'_s)^2} &= -a_\alpha
\end{aligned} \quad (1.14)$$

Now multiply equation 1.13 by $x_s'^\alpha$, which gives,

$$\begin{aligned}
x'_s \cdot \dot{x}'_s &= 2(a \cdot x'_s)(x'_s \cdot x'_s) - (a \cdot x'_s)(x'_s \cdot x'_s) \\
x'_s \cdot \dot{x}'_s &= (a \cdot x'_s)(x'_s \cdot x'_s)
\end{aligned}$$

Substituting this back in equation 1.14,

$$\begin{aligned}
\frac{\dot{x}'_{s\alpha}}{x'_s \cdot x'_s} - \frac{2(x'_s \cdot \dot{x}'_s)x'_{s\alpha}}{(x'_s \cdot x'_s)^2} &= -a_\alpha \\
\frac{1}{x'_s \cdot x'_s} \frac{d}{ds} x'_{s\alpha} - x'_{s\alpha} \frac{1}{(x'_s \cdot x'_s)^2} \frac{d}{ds} (x'_s \cdot x'_s) &= -a_\alpha \\
\frac{1}{x'_s \cdot x'_s} \frac{d}{ds} x'_{s\alpha} + x'_{s\alpha} \frac{d}{ds} \frac{1}{x'_s \cdot x'_s} &= -a_\alpha \\
\frac{d}{ds} \left(\frac{x'_{s\alpha}}{x'_s \cdot x'_s} \right) &= -a_\alpha \\
\frac{x'_{s\alpha}}{x'_s \cdot x'_s} &= -sa_\alpha + \frac{x_\alpha}{x \cdot x}
\end{aligned}$$

What remains is a mere algebraic equation which is straightforward to solve, first, contract with x'^{α}_s ,

$$\begin{aligned}
\frac{x'_s \cdot x'_s}{x'_s \cdot x'_s} &= 1 = -sx'_s \cdot a + \frac{x'_s \cdot x}{x \cdot x} \\
-sx'_s \cdot a + \frac{x'_s \cdot x}{x \cdot x} &= 1
\end{aligned}$$

Which gives one constraint, now, contract with $-sa^\alpha + \frac{x^\alpha}{x \cdot x}$,

$$\frac{1}{x'_s \cdot x'_s} \left(-sx'_s \cdot a + \frac{x'_s \cdot x}{x \cdot x} \right) = \left(-sa^\alpha + \frac{x^\alpha}{x \cdot x} \right) \left(-sa_\alpha + \frac{x_\alpha}{x \cdot x} \right)$$

Using the constraint derived,

$$\frac{1}{x'_s \cdot x'_s} = \left(-sa^\alpha + \frac{x^\alpha}{x \cdot x} \right) \left(-sa_\alpha + \frac{x_\alpha}{x \cdot x} \right)$$

Which gives the final constraint, substituting this back on our original equation gives,

$$\frac{x'_{s\alpha}}{x'_s \cdot x'_s} = -sa_\alpha + \frac{x_\alpha}{x \cdot x} \tag{1.15}$$

$$\begin{aligned}
x'_{s\alpha} \left(-sa^\alpha + \frac{x^\alpha}{x \cdot x} \right) \left(-sa_\alpha + \frac{x_\alpha}{x \cdot x} \right) &= -sa_\alpha + \frac{x_\alpha}{x \cdot x} \\
x'_{s\alpha} &= \frac{-sa_\alpha + \frac{x_\alpha}{x \cdot x}}{\left(-sa^\alpha + \frac{x^\alpha}{x \cdot x} \right) \left(-sa_\alpha + \frac{x_\alpha}{x \cdot x} \right)} \\
x'_{s\alpha} &= \frac{-sa_\alpha(x \cdot x) + x_\alpha}{(-sa^\alpha(x \cdot x) + x^\alpha) \left(-sa_\alpha + \frac{x_\alpha}{x \cdot x} \right)} \\
x'_{s\alpha} &= \frac{x_\alpha - sa_\alpha(x \cdot x)}{1 - 2sa \cdot x + s^2 a \cdot a(x \cdot x)} \tag{1.16}
\end{aligned}$$

These are the special conformal transformations, from 1.15 it's easy to interpret geometrically what they do. First, apply an inversion, $x^\mu \rightarrow \frac{x^\mu}{x \cdot x}$, second, do a translation $\frac{x^\mu}{x \cdot x} \rightarrow \frac{x^\mu}{x \cdot x} - sa^\mu$, and lastly perform a second inversion, $\frac{x^\mu}{x \cdot x} - sa^\mu \rightarrow \frac{\frac{x^\mu}{x \cdot x} - sa^\mu}{\left(\frac{x^\mu}{x \cdot x} - sa^\mu\right) \cdot \left(\frac{x^\mu}{x \cdot x} - sa_\mu\right)}$. As opposed to a first impression, they're indeed well behaved at $x \cdot x = 0$, as can be seen from 1.16, but sadly, they aren't well behaved for the denominator in 1.16 to be zero, that is,

$$1 - 2\tilde{s}a \cdot x + \tilde{s}^2(a \cdot a)(x \cdot x) = 0 \Rightarrow \tilde{s} = \frac{2(a \cdot x) \pm \sqrt{4(a \cdot x)^2 - 4(a \cdot a)(x \cdot x)}}{2(a \cdot a)(x \cdot x)}$$

Which only has a real solution for $x^\mu \propto a^\mu$, but, for this point away from the origin the transformation needs not to be well behaved, as the charts x, x'_s themselves don't need to be defined from the whole manifold, nevertheless, if we're working in a compact manifold we can make sense of this divergence $x'^\mu_s \rightarrow \infty$ as this is indeed a point of a compact manifold.

Problem 2

2.A)

First, let's state a few remarks, we have at our disposal the following algebras,

$$\begin{aligned} [L_m, L_n] &= (m - n)L_{n+m} + \frac{c}{12}m(m^2 - 1)\delta_{m,-n} \\ [\alpha_m^\mu, \alpha_n^\nu] &= mg^{\mu\nu}\delta_{m,-n} \\ [L_m, \alpha_n^\nu] &= -n\alpha_{n+m}^\nu \end{aligned}$$

Which should be interpreted only as elements of a Lie Algebra. But, we seek representations of this algebra as linear maps in our physical Hilbert space, as any Hilbert space, it'll contain a notion of inner product, to which can be defined adjoint of linear maps. Let's formally manipulate the adjoints of the algebra,

$$\begin{cases} [L_m, L_n]^\dagger &= (m - n)L_{n+m}^\dagger + \frac{c}{12}m(m^2 - 1)\delta_{m,-n} \\ [\alpha_m^\mu, \alpha_n^\nu]^\dagger &= mg^{\mu\nu}\delta_{m,-n} \\ [L_m, \alpha_n^\nu]^\dagger &= -n(\alpha_{n+m}^\nu)^\dagger \end{cases}$$

$$\begin{cases} [L_n^\dagger, L_m^\dagger] &= (-n - (-m))L_{n+m}^\dagger + \frac{c}{12}(-n)((-n)^2 - 1)\delta_{-n,-(-m)} \\ [(\alpha_n^\nu)^\dagger, (\alpha_m^\mu)^\dagger] &= -ng^{\mu\nu}\delta_{-n,-(-m)} \\ [(\alpha_n^\nu)^\dagger, L_m^\dagger] &= -[L_m^\dagger, (\alpha_n^\nu)^\dagger] = -n(\alpha_{n+m}^\nu)^\dagger \end{cases}$$

In this format it's clear to see that the algebra span by $L_n^\dagger, (\alpha_n^\mu)^\dagger$ is isomorphic to the algebra span by L_n, α_n^μ , in particular, the isomorphism is given by, $L_n^\dagger \rightarrow L_{-n}, (\alpha_n^\mu)^\dagger \rightarrow \alpha_{-n}^\mu$, as can be seen directly from the algebra by making the substitutions,

$$\begin{aligned} [L_{-n}, L_{-m}] &= (-n - (-m))L_{-n-m} + \frac{c}{12}(-n)((-n)^2 - 1)\delta_{-n,-(-m)} \\ [\alpha_{-n}^\nu, \alpha_{-m}^\mu] &= -ng^{\mu\nu}\delta_{-n,-(-m)} \\ [L_{-m}, \alpha_{-n}^\nu] &= n\alpha_{-n-m}^\nu \end{aligned}$$

Which is exactly the right algebra. We'll suppose that in our representation we have exactly $L_n^\dagger = L_{-n}, (\alpha_n^\mu)^\dagger = \alpha_{-n}^\mu$ as hinted above. A few comments on the inner product of our Hilbert space, it is non degenerate and non negative. About the non negative part, as we have a Lorentzian metric in the target space, at least one of the components of X — the time-like one — has a wrong sign in the kinetic term in the action, that is, some of the modes will inevitably 'create' negative norm states, but, these cannot be physical, and therefore must/can be removed. One way of doing such is fixing the gauge a priori, which isn't the case, the other way is restricting the polarization vectors to which the modes, of ∂X can couple to. That is, as example, if α_n^μ when acting in the vacuum creates a negative norm state, this by itself is not enough to guarantee that it annihilate the vacuum, we have to show that no matter what choice of ξ_μ is done, $\xi_\mu \alpha_n^\mu$ always is a negative norm state, then we can conclude α_n^μ annihilate the vacuum. Closing these initial remarks, we start by defining the vacuum state, Ψ_0 , as the one for which,

$$L_m \Psi_0 = 0, \quad m \geq -1$$

Let's analyze the action of $\xi_\mu \alpha_n^\mu$ in it,

$$\begin{aligned} L_m \xi_\mu \alpha_n^\mu \Psi_0 &= \xi_\mu (L_m \alpha_n^\mu - \alpha_n^\mu L_m) \Psi_0, \quad m \geq -1 \\ L_m \xi_\mu \alpha_n^\mu \Psi_0 &= \xi_\mu [L_m, \alpha_n^\mu] \Psi_0, \quad m \geq -1 \\ L_m \xi_\mu \alpha_n^\mu \Psi_0 &= -n \xi_\mu \alpha_{n+m}^\mu \Psi_0, \quad m \geq -1 \end{aligned} \quad (2.1)$$

The first step was add $0 = \xi_\mu \alpha_n^\mu L_m \Psi_0$, as L_m , $m \geq -1$ annihilate the vacuum. In the special case of $n = 0$, we have,

$$L_m \xi_\mu \alpha_0^\mu \Psi_0 = 0, \quad m \geq -1$$

That implies $\xi_\mu \alpha_0^\mu \Psi_0$ has exactly all the symmetries of the vacuum, hence, it must be proportional to it, $\xi_\mu \alpha_0^\mu \Psi_0 = \xi \Psi_0$ for some constant ξ , now we have to evoke that we do have an inner product (\cdot, \cdot) , so that the special case of $n = -m$ in 2.1 is,

$$\begin{aligned} L_m \xi_\mu \alpha_{-m}^\mu \Psi_0 &= m \xi_\mu \alpha_0^\mu \Psi_0, \quad m \geq -1 \\ L_m \xi_\mu \alpha_{-m}^\mu \Psi_0 &= m \xi \Psi_0, \quad m \geq -1 \\ (\Psi_0, L_m \xi_\mu \alpha_{-m}^\mu \Psi_0) &= (\Psi_0, m \xi \Psi_0), \quad m \geq -1 \\ (L_m^\dagger \Psi_0, \xi_\mu \alpha_{-m}^\mu \Psi_0) &= m \xi (\Psi_0, \Psi_0), \quad m \geq -1 \\ (L_{-m} \Psi_0, \xi_\mu \alpha_{-m}^\mu \Psi_0) &= m \xi (\Psi_0, \Psi_0), \quad m \geq -1 \end{aligned}$$

For $m = \pm 1$,

$$\begin{aligned} (L_{\mp 1} \Psi_0, \xi_\mu \alpha_{\mp 1}^\mu \Psi_0) &= \pm \xi (\Psi_0, \Psi_0) \\ (0, \xi_\mu \alpha_{\mp 1}^\mu \Psi_0) &= \pm \xi (\Psi_0, \Psi_0) \\ 0 &= \pm \xi (\Psi_0, \Psi_0) \end{aligned}$$

Of course $(\Psi_0, \Psi_0) \neq 0$ otherwise it would be the null state, hence, $\xi = 0 \Rightarrow \xi_\mu \alpha_0^\mu \Psi_0 = \xi \Psi_0 = 0$. As this is valid for any choice of ξ_μ , we just proved that $\alpha_0^\mu \Psi_0 = 0$.

2.B)

Problem 3

3.A)

Our action is,

$$S = \frac{1}{4\pi} \int d^2z \psi \bar{\partial} \psi$$

To obtain the equation of motion is simple, first we set up the path integral of a total derivative, which is zero, the argument of the total derivative we set to $\exp(-S)\mathcal{O}$, where \mathcal{O} is any combination of local fields that does not contain $\psi(z_1, \bar{z}_1)$, then,

$$\begin{aligned} 0 &= \int \mathcal{D}\psi \frac{\delta}{\delta\psi(z_1, \bar{z}_1)} [\exp(-S)\mathcal{O}] \\ 0 &= \int \mathcal{D}\psi \frac{\delta}{\delta\psi(z_1, \bar{z}_1)} [\exp(-S)] \mathcal{O} \end{aligned}$$

As $\frac{\delta}{\delta\psi(z_1, \bar{z}_1)} \mathcal{O}$ is zero except for $\mathcal{O} = \psi(z_1, \bar{z}_1)$, to which it's $\delta^{(2)}(0)$, but, that's not the case,

$$\begin{aligned} 0 &= \int \mathcal{D}\psi \frac{\delta}{\delta\psi(z_1, \bar{z}_1)} [\exp(-S)] \mathcal{O} \\ 0 &= -\frac{1}{4\pi} \int \mathcal{D}\psi \frac{\delta}{\delta\psi(z_1, \bar{z}_1)} \left[\int d^2z_2 \psi(z_2, \bar{z}_2) \partial_{\bar{z}_2} \psi(z_2, \bar{z}_2) \right] \mathcal{O} \exp(-S) \\ 0 &= -\frac{1}{4\pi} \int \mathcal{D}\psi \int d^2z_2 [\delta^{(2)}(z_2 - z_1) \partial_{\bar{z}_2} \psi(z_2, \bar{z}_2) - \psi(z_2, \bar{z}_2) \partial_{\bar{z}_2} \delta^{(2)}(z_2 - z_1)] \mathcal{O} \exp(-S) \end{aligned}$$

As long as we're dealing with the closed fermion, we need not to worry about the boundary conditions, as there is no boundary, so,

$$\begin{aligned} 0 &= -\frac{1}{4\pi} \int \mathcal{D}\psi \int d^2z_2 [\delta^{(2)}(z_2 - z_1) \partial_{\bar{z}_2} \psi(z_2, \bar{z}_2) + \partial_{\bar{z}_2} \psi(z_2, \bar{z}_2) \delta^{(2)}(z_2 - z_1)] \mathcal{O} \exp(-S) \\ 0 &= -\frac{1}{2\pi} \int \mathcal{D}\psi \partial_{\bar{z}_1} \psi(z_1, \bar{z}_1) \mathcal{O} \exp(-S) \end{aligned}$$

What in the Operator formalism would account for some radially ordered expectation value under some state, which is specified with the boundary conditions on the path integral itself,

$$\partial_{\bar{z}_1} \langle \psi(z_1, \bar{z}_1) \mathcal{O} \rangle = 0$$

As both the operator \mathcal{O} and the state are arbitrary, as long as there is no insertion of $\psi(z_1, \bar{z}_1)$, we conclude the equation of motion in the operator form is just,

$$\bar{\partial} \psi = 0 \Rightarrow \psi(z, \bar{z}) \equiv \psi(z)$$

That is, ψ is, at least, a meromorphic operator/function. Now, let's repeat this to obtain the two point function, this time we include in \mathcal{O} a single factor of $\psi(z_1, \bar{z}_1)$ disguised as $\psi(z_2, \bar{z}_2)$, this also can be phrased as we being interested in the limit $z_2 \rightarrow z_1$,

$$\begin{aligned}
0 &= \int \mathcal{D}\psi \frac{\delta}{\delta\psi(z_1, \bar{z}_1)} [\exp(-S)\psi(z_2, \bar{z}_2)\mathcal{O}] \\
0 &= \int \mathcal{D}\psi \frac{\delta}{\delta\psi(z_1, \bar{z}_1)} [\exp(-S)\psi(z_2, \bar{z}_2)]\mathcal{O} \\
0 &= \int \mathcal{D}\psi \left\{ \frac{\delta}{\delta\psi(z_1, \bar{z}_1)} [\exp(-S)]\psi(z_2, \bar{z}_2) + \frac{\delta}{\delta\psi(z_1, \bar{z}_1)} [\psi(z_2, \bar{z}_2)] \exp(-S) \right\} \mathcal{O} \\
0 &= \int \mathcal{D}\psi \left\{ \frac{\delta}{\delta\psi(z_1, \bar{z}_1)} [-S]\psi(z_2, \bar{z}_2) + \delta^{(2)}(z_2 - z_1) \right\} \exp(-S)\mathcal{O}
\end{aligned}$$

As we derived before, $\frac{\delta}{\delta\psi(z_1, \bar{z}_1)} [-S] = -\frac{1}{2\pi} \partial_{\bar{z}_1} \psi(z_1, \bar{z}_1)$,

$$0 = \int \mathcal{D}\psi \left\{ -\frac{1}{2\pi} \partial_{\bar{z}_1} \psi(z_1, \bar{z}_1) \psi(z_2, \bar{z}_2) + \delta^{(2)}(z_2 - z_1) \right\} \exp(-S)\mathcal{O}$$

Which is, translating to the operator formalism,

$$\begin{aligned}
\partial_{\bar{z}_1} \langle \psi(z_1) \psi(z_2) \mathcal{O} \rangle &= \langle 2\pi \delta^{(2)}(z_2 - z_1) \mathcal{O} \rangle \\
\partial_{\bar{z}_1} \psi(z_1) \psi(z_2) &= 2\pi \delta^{(2)}(z_2 - z_1), \quad |z_1| \geq |z_2|
\end{aligned}$$

So that we can interpret this as a operator equality. The condition $|z_1| \geq |z_2|$ is due to the implicit radial ordering in the expectation value, and also notice that in the path integral formulation $\psi = \psi(z_1, \bar{z}_1)$, but in the operator formalism $\psi = \psi(z_1)$, as every operator is always ‘*on-shell*’, while the path integral integrands aren’t. Now we integrate this two point function, pick any compact closed region R in the complex plane which contains z_2 and not any other of the points of the insertions \mathcal{O} , with also the boundary being a continuous curve, for all $m \in \mathbb{N}$ the following is true,

$$\begin{aligned}
\partial_{\bar{z}_1} \langle \psi(z_1) \psi(z_2) \mathcal{O} \rangle &= 2\pi \delta^{(2)}(z_2 - z_1) \langle \mathcal{O} \rangle \\
(z_1 - z_2)^m \partial_{\bar{z}_1} \langle \psi(z_1) \psi(z_2) \mathcal{O} \rangle &= 2\pi (z_1 - z_2)^m \delta^{(2)}(z_2 - z_1) \langle \mathcal{O} \rangle, \quad m \in \mathbb{N} \\
\int_R d^2 z_1 (z_1 - z_2)^m \partial_{\bar{z}_1} \langle \psi(z_1) \psi(z_2) \mathcal{O} \rangle &= 2\pi \int_R d^2 z_1 (z_1 - z_2)^m \delta^{(2)}(z_2 - z_1) \langle \mathcal{O} \rangle, \quad m \in \mathbb{N} \\
\int_R d^2 z_1 \partial_{\bar{z}_1} \{ (z_1 - z_2)^m \langle \psi(z_1) \psi(z_2) \mathcal{O} \rangle \} &= 2\pi (z_1 - z_2)^m \Big|_{z_1=z_2} \langle \mathcal{O} \rangle, \quad m \in \mathbb{N}
\end{aligned}$$

In the last line we used the fact ‘*trivial*’ fact that $\partial_{\bar{z}_1} (z_1 - z_2)^m = 0$, which is valid for $m \in \mathbb{N}$, but isn’t for negative non-integer values.

$$\begin{aligned}
\int_R d^2 z_1 \partial_{\bar{z}_1} \{ (z_1 - z_2)^m \langle \psi(z_1) \psi(z_2) \mathcal{O} \rangle \} &= 2\pi (z_1 - z_2)^m \Big|_{z_1=z_2} \langle \mathcal{O} \rangle, \quad m \in \mathbb{N} \\
-i \int_{\partial R} dz_1 (z_1 - z_2)^m \langle \psi(z_1) \psi(z_2) \mathcal{O} \rangle &= 2\pi \delta_{m,0} \langle \mathcal{O} \rangle, \quad m \in \mathbb{N}
\end{aligned}$$

This is simply the complex version of the divergence theorem,

$$\begin{aligned}
-i \int_{\partial R} dz_1 (z_1 - z_2)^m \langle \psi(z_1) \psi(z_2) \mathcal{O} \rangle &= 2\pi \delta_{m,0} \langle \mathcal{O} \rangle, \quad m \in \mathbb{N} \\
\int_{\partial R} \frac{dz_1}{2\pi i} (z_1 - z_2)^m \langle \psi(z_1) \psi(z_2) \mathcal{O} \rangle &= \delta_{m,0} \langle \mathcal{O} \rangle, \quad m \in \mathbb{N}
\end{aligned}$$

Well, the left-hand side of this last equation picks up the pole of m -th order of the expression $\langle \psi(z_1) \psi(z_2) \mathcal{O} \rangle$ under $z_1 \rightarrow z_2$, but, the right-hand side is only non-zero for $m = 0$, this is telling us that $\langle \psi(z_1) \psi(z_2) \mathcal{O} \rangle$ doesn't have any pole besides the first order one with residue $\langle \mathcal{O} \rangle$, so we can write,

$$\langle \psi(z_1) \psi(z_2) \mathcal{O} \rangle = \frac{1}{z_1 - z_2} \langle \mathcal{O} \rangle + \text{regular}$$

Or, in the operator formalism,

$$\begin{aligned}
\psi(z_1) \psi(z_2) &= \frac{1}{z_1 - z_2} + \text{regular}, \quad |z_1| \geq |z_2| \\
\psi(z_1) \psi(z_2) &= \frac{1}{z_1 - z_2} + : \psi(z_1) \psi(z_2) :, \quad |z_1| \geq |z_2|
\end{aligned}$$

Which is the OPE.

3.B)

The energy momentum tensor has two components, a meromorphic one $T(z)$, and an anti-meromorphic one $\bar{T}(\bar{z})$, as we have at our disposal only a weight $(\frac{1}{2}, 0)$ field, it's impossible to construct a anti-meromorphic energy momentum tensor, but, for the meromorphic one, let us see what kind of combinations have the right weights. First, we remember that our energy momentum tensor must be normal ordered, so, all the possible normal ordered weight $(2, 0)$ combinations of ψ and derivatives are,

$$: \psi \partial \psi : (z), : (\partial \psi) \psi : (z), : \partial(\psi \psi) : (z), : \psi \psi \psi \psi : (z)$$

The third and fourth are zero due to the fermionic statistic, and the first and second are linear dependent also because the fermionic statistic. Hence, up to a unknown constant α , the meromorphic energy momentum tensor is,

$$T(z) = \alpha : \psi \partial \psi : (z)$$

Let's use this expression to compute the following,

$$: T(z_1) T(z_2) : = T(z_1) T(z_2) + \alpha^2 \text{ contractions } \{ : (\psi \partial \psi)(z_1) (\psi \partial \psi)(z_2) : \}, \quad |z_1| \geq |z_2|$$

Where the contraction part means exchanging pairs of ψ with different z by $-(z_1 - z_2)^{-1}$, as cause of the fermion statistic, we have to pay attention to the signs, also, we'll omit for now the $|z_1| \geq |z_2|$, but it will still be imposed,

$$\begin{aligned}
& : T(z_1)T(z_2) : = T(z_1)T(z_2) + \alpha^2 : \overbrace{(\psi\partial\psi)(z_1)(\psi\partial\psi)(z_2)} : + \alpha^2 : \overbrace{(\psi\partial\psi)(z_1)(\psi\partial\psi)(z_2)} : \\
& \quad + \alpha^2 : \overbrace{(\psi\partial\psi)(z_1)(\psi\partial\psi)(z_2)} : + \alpha^2 : \overbrace{(\psi\partial\psi)(z_1)(\psi\partial\psi)(z_2)} : \\
& \quad + \alpha^2 : \overbrace{(\psi\partial\psi)(z_1)(\psi\partial\psi)(z_2)} : + \alpha^2 : \overbrace{(\psi\partial\psi)(z_1)(\psi\partial\psi)(z_2)} : \\
& : T(z_1)T(z_2) : = T(z_1)T(z_2) + \frac{1}{z_1 - z_2} \alpha^2 : \partial\psi(z_1)\partial\psi(z_2) : - \partial_{z_2} \left(\frac{1}{z_1 - z_2} \right) \alpha^2 : \partial\psi(z_1)\psi(z_2) : \\
& \quad - \alpha^2 \partial_{z_1} \left(\frac{1}{z_1 - z_2} \right) : \psi(z_1)\partial\psi(z_2) : + \alpha^2 \partial_{z_1} \partial_{z_2} \left(\frac{1}{z_1 - z_2} \right) : \psi(z_1)\psi(z_2) : \\
& \quad + \alpha^2 \frac{1}{z_1 - z_2} \partial_{z_1} \partial_{z_2} \left(\frac{1}{z_1 - z_2} \right) - \alpha^2 \partial_{z_2} \left(\frac{1}{z_1 - z_2} \right) \partial_{z_1} \left(\frac{1}{z_1 - z_2} \right) \\
& : T(z_1)T(z_2) : = T(z_1)T(z_2) + \frac{\alpha^2}{z_1 - z_2} : \partial\psi(z_1)\partial\psi(z_2) : - \frac{\alpha^2}{(z_1 - z_2)^2} : \partial\psi(z_1)\psi(z_2) : \\
& \quad + \frac{\alpha^2}{(z_1 - z_2)^2} : \psi(z_1)\partial\psi(z_2) : + \partial_{z_1} \left(\frac{\alpha^2}{(z_1 - z_2)^2} \right) : \psi(z_1)\psi(z_2) : \\
& \quad + \frac{\alpha^2}{z_1 - z_2} \partial_{z_1} \left(\frac{1}{(z_1 - z_2)^2} \right) - \frac{\alpha^2}{(z_1 - z_2)^2} \partial_{z_1} \left(\frac{1}{z_1 - z_2} \right) \\
& : T(z_1)T(z_2) : = T(z_1)T(z_2) + \frac{\alpha^2}{z_1 - z_2} : \partial\psi(z_1)\partial\psi(z_2) : - \frac{\alpha^2}{(z_1 - z_2)^2} : \partial\psi(z_1)\psi(z_2) : \\
& \quad + \frac{\alpha^2}{(z_1 - z_2)^2} : \psi(z_1)\partial\psi(z_2) : - 2 \frac{\alpha^2}{(z_1 - z_2)^3} : \psi(z_1)\psi(z_2) : \\
& \quad - 2 \frac{\alpha^2}{z_1 - z_2} \frac{1}{(z_1 - z_2)^3} + \frac{\alpha^2}{(z_1 - z_2)^2} \frac{1}{(z_1 - z_2)^2} \\
& : T(z_1)T(z_2) : = T(z_1)T(z_2) + \frac{\alpha^2}{z_1 - z_2} : \partial\psi(z_1)\partial\psi(z_2) : - \frac{\alpha^2}{(z_1 - z_2)^2} : \partial\psi(z_1)\psi(z_2) : \\
& \quad + \frac{\alpha^2}{(z_1 - z_2)^2} : \psi(z_1)\partial\psi(z_2) : - 2 \frac{\alpha^2}{(z_1 - z_2)^3} : \psi(z_1)\psi(z_2) : \\
& \quad - \frac{\alpha^2}{(z_1 - z_2)^4}
\end{aligned}$$

To proceed further we have to Taylor expand both $\psi(z_1)$, $\partial\psi(z_1)$, of course this will generate regular terms in our expansion, as for example, the second term in the right-hand side,

$$\begin{aligned}
\frac{1}{z_1 - z_2} : \partial\psi(z_1)\partial\psi(z_2) : &= \frac{1}{z_1 - z_2} : \partial\psi(z_2)\partial\psi(z_2) : + \sum_{n=1}^{\infty} \frac{1}{z_1 - z_2} \frac{1}{n!} (z_1 - z_2)^n : \partial^n \psi(z_2)\partial\psi(z_2) : \\
\frac{1}{z_1 - z_2} : \partial\psi(z_1)\partial\psi(z_2) : &= \frac{1}{z_1 - z_2} : \partial\psi\partial\psi : (z_2) + \sum_{n=1}^{\infty} \frac{(z_1 - z_2)^{n-1}}{n!} : \partial^n \psi\partial\psi : (z_2)
\end{aligned}$$

It's clear that the sum in the right-hand side is of only regular terms, so that,

$$\frac{1}{z_1 - z_2} : \partial\psi(z_1)\partial\psi(z_2) : = \frac{1}{z_1 - z_2} : \partial\psi\partial\psi : (z_2) + \text{regular}$$

But $: \partial\psi\partial\psi :$ is zero by statistics, hence,

$$\frac{1}{z_1 - z_2} : \partial\psi(z_1)\partial\psi(z_2) : = \text{regular}$$

We do this procedure for all the terms in the expansion,

$$\begin{aligned} : T(z_1)T(z_2) : &= T(z_1)T(z_2) - \frac{\alpha^2}{(z_1 - z_2)^2} : \partial\psi\psi : (z_2) - \frac{\alpha^2}{(z_1 - z_2)} : \partial^2\psi\psi : (z_2) \\ &+ \frac{\alpha^2}{(z_1 - z_2)^2} : \psi\partial\psi : (z_2) + \frac{\alpha^2}{(z_1 - z_2)} : \partial\psi\partial\psi : (z_2) \\ &- 2\frac{\alpha^2}{(z_1 - z_2)^3} : \psi\psi : (z_2) - 2\frac{\alpha^2}{(z_1 - z_2)^2} : \partial\psi\psi : (z_2) - \frac{\alpha^2}{(z_1 - z_2)} : \partial^2\psi\psi : (z_2) \\ &- \frac{\alpha^2}{(z_1 - z_2)^4} + \text{regular} \end{aligned}$$

Removing the terms which are zero by statistics and grouping the others,

$$\begin{aligned} : T(z_1)T(z_2) : &= T(z_1)T(z_2) + \frac{4\alpha^2}{(z_1 - z_2)^2} : \psi\partial\psi : (z_2) + \frac{2\alpha^2}{(z_1 - z_2)} : \psi\partial^2\psi : (z_2) \\ &- \frac{\alpha^2}{(z_1 - z_2)^4} + \text{regular} \end{aligned}$$

As $: \partial\psi\partial\psi : \equiv 0$ we can add it as we please,

$$\begin{aligned} : T(z_1)T(z_2) : &= T(z_1)T(z_2) + \frac{4\alpha^2}{(z_1 - z_2)^2} : \psi\partial\psi : (z_2) + \frac{2\alpha^2}{(z_1 - z_2)} : \psi\partial^2\psi + \partial\psi\partial\psi : (z_2) \\ &- \frac{\alpha^2}{(z_1 - z_2)^4} + \text{regular} \\ : T(z_1)T(z_2) : &= T(z_1)T(z_2) + \frac{4\alpha^2}{(z_1 - z_2)^2} : \psi\partial\psi : (z_2) + \frac{2\alpha^2}{(z_1 - z_2)} \partial : \psi\partial\psi : (z_2) \\ &- \frac{\alpha^2}{(z_1 - z_2)^4} + \text{regular} \\ : T(z_1)T(z_2) : &= T(z_1)T(z_2) + \frac{4\alpha}{(z_1 - z_2)^2} T(z_2) + \frac{2\alpha}{(z_1 - z_2)} \partial T(z_2) \\ &- \frac{\alpha^2}{(z_1 - z_2)^4} + \text{regular} \end{aligned}$$

Well, $: T(z_1)T(z_2) :$ itself is regular, then,

$$T(z_1)T(z_2) = -\frac{4\alpha}{(z_1 - z_2)^2} T(z_2) - \frac{2\alpha}{(z_1 - z_2)} \partial T(z_2) + \frac{\alpha^2}{(z_1 - z_2)^4} + \text{regular}$$

But we know the general form of the TT OPE,

$$T(z_1)T(z_2) = \frac{2}{(z_1 - z_2)^2} T(z_2) + \frac{1}{(z_1 - z_2)} \partial T(z_2) + \frac{c}{2(z_1 - z_2)^4} + \text{regular}$$

From where is easy to read $\alpha = -\frac{1}{2}$, and also, $c = 2\alpha^2 = 2\frac{1}{4} = \frac{1}{2}$, that is, the meromorphic component of the energy momentum tensor, and the central charge are,

$$T(z) = -\frac{1}{2} : \psi \partial \psi : (z), \quad c = \frac{1}{2}$$

3.C)

Now let ψ have an additional index, ψ^i which takes two values $i = 1, 2$. This is a realization of a internal $SO(2)$ symmetry, to why, let $\psi'^i = \Lambda^i_j \psi^j$, the change in the action is,

$$\begin{aligned} S &= \frac{1}{4\pi} \int d^2z \delta_{ij} \psi^i \bar{\partial} \psi^j \\ S' &= \frac{1}{4\pi} \int d^2z \delta_{ij} \psi'^i \bar{\partial} \psi'^j \\ S' &= \frac{1}{4\pi} \int d^2z \delta_{ij} \Lambda^i_k \Lambda^j_l \psi^k \bar{\partial} \psi^l \end{aligned}$$

As long as $\delta_{ij} \Lambda^i_k \Lambda^j_l = \delta_{kl}$, this transformation is a symmetry. But that's exactly the definition of Λ belonging to the $O(2)$ group, but as we're just interested in elements continuously connected to the identity, we set $SO(2)$. The group definition under an infinitesimal transformation,

$$\begin{aligned} \delta_{kl} &= \delta_{ij} \Lambda^i_k \Lambda^j_l \\ \delta_{kl} &= \delta_{ij} (\delta^i_k + \omega^i_k + \mathcal{O}(\omega^2)) (\delta^j_l + \omega^j_l + \mathcal{O}(\omega^2)) \\ \delta_{kl} &= \delta_{kl} + \omega_{lk} + \omega_{kl} + \mathcal{O}(\omega^2) \\ \omega_{lk} &= -\omega_{kl} \end{aligned}$$

We can use this information to write,

$$\begin{aligned} \Lambda^i_j &= \delta^i_j + \omega^i_j + \mathcal{O}(\omega^2) \\ \Lambda^i_j &= \delta^i_j + \frac{1}{2} \omega_{lk} (\delta^{il} \delta^k_j - \delta^{ik} \delta^l_j) + \mathcal{O}(\omega^2) \\ \Lambda^i_j &= \delta^i_j - \frac{i}{2} \omega_{lk} i (\delta^{il} \delta^k_j - \delta^{ik} \delta^l_j) + \mathcal{O}(\omega^2) \\ \Lambda^i_j &= \delta^i_j - \frac{i}{2} \omega_{[lk]} T^{[lk]i}_j + \mathcal{O}(\omega^2) \end{aligned}$$

Where we defined the generator of the group, $T^{[lk]i}_j = i(\delta^{il} \delta^k_j - \delta^{ik} \delta^l_j)$. The sign of the definition is arbitrary. A finite transformation is then,

$$\psi^i \rightarrow \Lambda^i_j \psi^j = \exp \left(-\frac{i}{2} \omega_{[lk]} T^{[lk]} \right)^i_j \psi^j$$

To obtain the conserved current of this symmetry we make a little trick, set ω as a function $\omega(z, \bar{z})$, so that the infinitesimal variation of the action is,

$$\begin{aligned}
S &= \frac{1}{4\pi} \int d^2z \delta_{ij} \psi^i \bar{\partial} \psi^j \\
S' &= \frac{1}{4\pi} \int d^2z \delta_{ij} \psi'^i \bar{\partial} \psi'^j \\
S' &= \frac{1}{4\pi} \int d^2z \delta_{ij} \Lambda^i_k \psi^k \bar{\partial} (\Lambda^j_l \psi^l) \\
S' &= \frac{1}{4\pi} \int d^2z \delta_{ij} \Lambda^i_k \psi^k \{ \Lambda^j_l \bar{\partial} \psi^l + \psi^l \bar{\partial} \Lambda^j_l \} \\
S' &= \frac{1}{4\pi} \int d^2z \delta_{ij} \Lambda^i_k \psi^k \left\{ \Lambda^j_l \bar{\partial} \psi^l - \frac{i}{2} T^{[ab]m}_l \psi^l \Lambda^j_m \bar{\partial} \omega_{[ab]} \right\} \\
S' &= \frac{1}{4\pi} \int d^2z \Lambda_{jk} \psi^k \Lambda^j_l \bar{\partial} \psi^l - \frac{1}{4\pi} \frac{i}{2} \int d^2z \Lambda_{jk} \psi^k T^{[ab]m}_l \psi^l \Lambda^j_m \bar{\partial} \omega_{[ab]} \\
S' &= \frac{1}{4\pi} \int d^2z \delta_{kl} \psi^k \bar{\partial} \psi^l - \frac{1}{4\pi} \frac{i}{2} \int d^2z \delta_{km} \psi^k T^{[ab]m}_l \psi^l \bar{\partial} \omega_{[ab]} \\
S' &= S - \frac{1}{4\pi} \frac{i}{2} \int d^2z \psi^k T^{[ab]}_{kl} \psi^l \bar{\partial} \omega_{[ab]} \\
\delta S &= -\frac{1}{4\pi} \frac{i}{2} \int d^2z \bar{\partial} \left\{ \psi^k T^{[ab]}_{kl} \psi^l \omega_{[ab]} \right\} + \frac{1}{4\pi} \frac{i}{2} \int d^2z \bar{\partial} \left\{ \psi^k T^{[ab]}_{kl} \psi^l \right\} \omega_{[ab]}
\end{aligned}$$

With suitable boundary conditions on $\omega_{[ab]}$, or, with no boundaries, we conclude,

$$\delta S = \frac{1}{4\pi} \frac{i}{2} \int d^2z \bar{\partial} \left\{ \psi^k T^{[ab]}_{kl} \psi^l \right\} \omega_{[ab]}$$

Now back to $\omega_{[ab]}$ being a constant — so that $\delta S = 0$ —, using also that there is just one two index anti-symmetric tensor in two dimensions, the Levi-Civita ϵ_{ab} , $\epsilon_{12} = 1$, we conclude that $\omega_{[ab]} = \omega \epsilon_{ab}$,

$$\begin{aligned}
0 &= \frac{1}{4\pi} \frac{i}{2} \omega \epsilon_{ab} T^{[ab]}_{kl} \int d^2z \bar{\partial} \{ \psi^k \psi^l \} \\
0 &= -\frac{\omega \epsilon_{12}}{2\pi i} \frac{1}{2} T^{12}_{kl} \int d^2z \bar{\partial} \{ \psi^k \psi^l \}
\end{aligned}$$

So there is just one current, as there is just one linear independent generator. The classical current can be read from the last equation as,

$$j(z) = \frac{1}{2} T^{12}_{kl} \psi^k \psi^l(z)$$

We factored out the $-\frac{\omega \epsilon_{12}}{2\pi i}$ as is usual to obtain an interpretation of the residue of the j OPE. The quantum version of it will require normal ordering,

$$\begin{aligned}
j(z) &= \frac{1}{2} T^{12}_{kl} : \psi^k \psi^l : (z) \\
j(z) &= \frac{1}{2} i (\delta_k^1 \delta_l^2 - \delta_k^2 \delta_l^1) : \psi^k \psi^l : (z) \\
j(z) &= \frac{i}{2} (: \psi^1 \psi^2 : (z) - : \psi^2 \psi^1 : (z))
\end{aligned}$$

$$j(z) = i : \psi^1 \psi^2 : (z)$$

Now, to get it's OPE we proceed as usual, computing the following normal ordered operator,

$$\begin{aligned} : j(z_1) j(z_2) : &= j(z_1) j(z_2) - : \overbrace{\psi^1 \psi^2 : (z_1)} : \overbrace{\psi^1 \psi^2 : (z_2)} - : \overbrace{\psi^1 \psi^2 : (z_1)} : \overbrace{\psi^1 \psi^2 : (z_2)} \\ &\quad - : \overbrace{\psi^1 \psi^2 : (z_1)} : \overbrace{\psi^1 \psi^2 : (z_2)} \\ : j(z_1) j(z_2) : &= j(z_1) j(z_2) - \frac{1}{z_1 - z_2} \psi^2(z_1) \psi^2(z_2) - \frac{1}{z_1 - z_2} \psi^1(z_1) \psi^1(z_2) \\ &\quad - \frac{1}{(z_1 - z_2)^2} \\ : j(z_1) j(z_2) : &= j(z_1) j(z_2) - \frac{1}{z_1 - z_2} \psi^2(z_2) \psi^2(z_2) - \frac{1}{z_1 - z_2} \psi^1(z_1) \psi^1(z_1) \\ &\quad - \frac{1}{(z_1 - z_2)^2} + \text{regular} \end{aligned}$$

In the first to second line we used the $\psi\psi$ OPE, remembering that the $\psi^1\psi^2$ OPE has no poles, from the second line to third, we expanded every term around z_2 . And now we set the terms which are zero by statistics,

$$\begin{aligned} : j(z_1) j(z_2) : &= j(z_1) j(z_2) - \frac{1}{(z_1 - z_2)^2} + \text{regular} \\ j(z_1) j(z_2) &= \frac{1}{(z_1 - z_2)^2} + \text{regular} \end{aligned}$$

A rather simple OPE.

3.D)

Much of what we did can be recycled here. First, we start with N real fermions ψ^i , $i = 1, \dots, N$, and by similar arguments analyzing the action,

$$\begin{aligned} S &= \frac{1}{4\pi} \int d^2 z \delta_{ij} \psi^i \bar{\partial} \psi^j \\ S' &= \frac{1}{4\pi} \int d^2 z \delta_{ij} \psi^i \bar{\partial} \psi'^j \\ S' &= \frac{1}{4\pi} \int d^2 z \delta_{ij} \Lambda_k^i \Lambda_l^j \psi^k \bar{\partial} \psi^l \end{aligned}$$

We obtain the constraint of this being a symmetry, $\delta_{ij} \Lambda_k^i \Lambda_l^j = \delta_{kl}$, which is the defining property of the $O(N)$ group. Of course, the relevant part for generators and Lie Algebras is just the part connected to the identity, $SO(N)$, the same reasoning of infinitesimal transformations allows us to write,

$$\begin{aligned} \delta_{kl} &= \delta_{ij} \Lambda_k^i \Lambda_l^j \\ \delta_{kl} &= \delta_{ij} (\delta_k^i + \omega_k^i + \mathcal{O}(\omega^2)) (\delta_l^j + \omega_l^j + \mathcal{O}(\omega^2)) \\ \delta_{kl} &= \delta_{kl} + \omega_{lk} + \omega_{kl} + \mathcal{O}(\omega^2) \end{aligned}$$

$$\omega_{lk} = -\omega_{kl}$$

We can use this information to write,

$$\begin{aligned}\Lambda^i_j &= \delta^i_j + \omega^i_j + \mathcal{O}(\omega^2) \\ \Lambda^i_j &= \delta^i_j + \frac{1}{2}\omega_{lk}(\delta^{il}\delta^k_j - \delta^{ik}\delta^l_j) + \mathcal{O}(\omega^2) \\ \Lambda^i_j &= \delta^i_j - \frac{i}{2}\omega_{lk}(\delta^{il}\delta^k_j - \delta^{ik}\delta^l_j) + \mathcal{O}(\omega^2) \\ \Lambda^i_j &= \delta^i_j - \frac{i}{2}\omega_{[lk]}T^{[lk]i}_j + \mathcal{O}(\omega^2) \\ \Lambda^i_j &= \exp\left(-\frac{i}{2}\omega_{[lk]}T^{[lk]i}_j\right)_j\end{aligned}$$

To obtain the current we do again the same trick of promoting $\omega_{[lk]}$ to a function and computing the change in the action,

$$\begin{aligned}S &= \frac{1}{4\pi} \int d^2z \delta_{ij} \psi^i \bar{\partial} \psi^j \\ S' &= \frac{1}{4\pi} \int d^2z \delta_{ij} \psi'^i \bar{\partial} \psi'^j \\ S' &= \frac{1}{4\pi} \int d^2z \delta_{ij} \Lambda^i_k \psi^k \bar{\partial} (\Lambda^j_l \psi^l) \\ S' &= \frac{1}{4\pi} \int d^2z \delta_{ij} \Lambda^i_k \psi^k \{ \Lambda^j_l \bar{\partial} \psi^l + \psi^l \bar{\partial} \Lambda^j_l \} \\ S' &= \frac{1}{4\pi} \int d^2z \delta_{ij} \Lambda^i_k \psi^k \left\{ \Lambda^j_l \bar{\partial} \psi^l - \frac{i}{2} T^{[ab]m}_l \psi^l \Lambda^j_m \bar{\partial} \omega_{[ab]} \right\} \\ S' &= \frac{1}{4\pi} \int d^2z \Lambda_{jk} \psi^k \Lambda^j_l \bar{\partial} \psi^l - \frac{1}{4\pi} \frac{i}{2} \int d^2z \Lambda_{jk} \psi^k T^{[ab]m}_l \psi^l \Lambda^j_m \bar{\partial} \omega_{[ab]} \\ S' &= \frac{1}{4\pi} \int d^2z \delta_{kl} \psi^k \bar{\partial} \psi^l - \frac{1}{4\pi} \frac{i}{2} \int d^2z \delta_{km} \psi^k T^{[ab]m}_l \psi^l \bar{\partial} \omega_{[ab]} \\ S' &= S - \frac{1}{4\pi} \frac{i}{2} \int d^2z \psi^k T^{[ab]}_{kl} \psi^l \bar{\partial} \omega_{[ab]} \\ \delta S &= -\frac{1}{4\pi} \frac{i}{2} \int d^2z \bar{\partial} \left\{ \psi^k T^{[ab]}_{kl} \psi^l \omega_{[ab]} \right\} + \frac{1}{4\pi} \frac{i}{2} \int d^2z \bar{\partial} \left\{ \psi^k T^{[ab]}_{kl} \psi^l \right\} \omega_{[ab]}\end{aligned}$$

With suitable boundary conditions on $\omega_{[ab]}$, or, with no boundaries, we conclude,

$$\delta S = \frac{1}{4\pi} \frac{i}{2} \int d^2z \bar{\partial} \left\{ \psi^k T^{[ab]}_{kl} \psi^l \right\} \omega_{[ab]}$$

Demoting $\omega_{[ab]}$ back to being a constant — so that $\delta S = 0$ —, we obtain,

$$0 = -\frac{\omega_{[ab]}}{2\pi i} \frac{1}{2} T^{[ab]}_{kl} \int d^2z \bar{\partial} \{ \psi^k \psi^l \}$$

Again, the usual normalization is done by removing $-\frac{\omega_{[ab]}}{2\pi i}$, so that the current is,

$$j^{[ab]}(z) = \frac{1}{2} T_{kl}^{[ab]} : \psi^k \psi^l : (z)$$

This is not a very orthodox way of displaying, all of our index here have the range $1, \dots, N$, but notice the generators are indexed by an anti-symmetric pair of index, $[ab]$, so there are $\frac{N(N-1)}{2}$ such, this is only useful if we work in the fundamental representation, so, we'll change now to a more common notation, instead of labeling the generators by $\frac{N(N-1)}{2}$ anti-symmetric pairs of index, we'll label them by just one index with range $A = 1, \dots, \frac{N(N-1)}{2}$, so that the same equation reads,

$$j^A(z) = \frac{1}{2} T_{kl}^A : \psi^k \psi^l : (z)$$

This is more aesthetic for Lie Algebras in general, but not always desired — Lorentz group as example —. Now we proceed by computing the OPE associated with this collection of currents,

$$: j^A(z_1) j^B(z_2) : =$$

3.E)

3.F)

Problem 4

4.A)

4.B)

4.C)

4.D)

4.E)

Problem 5

5.A)

5.B)

Problem 6

6.A)

6.B)

6.C)

6.D)

A Bosonic CFT

Here we're going to derive some results of a bosonic CFT of X^μ ,

$$S = \frac{1}{4\pi\alpha'} \int d^2\sigma \partial_a X^\mu \partial^a X_\mu$$

Where it's understood a flat '*world-sheet*' metric δ_{ab} , we make a change of variables from σ^1, σ^2 to $\sigma^z \equiv z = \sigma^1 + i\sigma^2, \sigma^{\bar{z}} \equiv \bar{z} = \sigma^1 - i\sigma^2$, so that the new metric reads,

$$\begin{aligned} g^{zz} &= \frac{\partial z}{\partial \sigma^1} \frac{\partial z}{\partial \sigma^1} \delta^{11} + \frac{\partial z}{\partial \sigma^2} \frac{\partial z}{\partial \sigma^2} \delta^{22} \\ g^{z\bar{z}} &= 1 - 1 = 0 \\ g^{\bar{z}\bar{z}} &= \frac{\partial \bar{z}}{\partial \sigma^1} \frac{\partial \bar{z}}{\partial \sigma^1} \delta^{11} + \frac{\partial \bar{z}}{\partial \sigma^2} \frac{\partial \bar{z}}{\partial \sigma^2} \delta^{22} \\ g^{z\bar{z}} &= 1 + 1 = 2 = g^{\bar{z}z} \\ g^{\bar{z}\bar{z}} &= \frac{\partial \bar{z}}{\partial \sigma^1} \frac{\partial \bar{z}}{\partial \sigma^1} \delta^{11} + \frac{\partial \bar{z}}{\partial \sigma^2} \frac{\partial \bar{z}}{\partial \sigma^2} \delta^{22} \\ g^{\bar{z}\bar{z}} &= 1 - 1 = 0 \end{aligned}$$

So that our action in these new coordinates is,

$$\begin{aligned} S &= \frac{1}{4\pi\alpha'} \int d\sigma^1 d\sigma^2 \sqrt{|\text{Det}[\delta]|} \partial_a X^\mu \partial^a X_\mu \\ S &= \frac{1}{4\pi\alpha'} \int dz d\bar{z} \sqrt{|\text{Det}[g]|} \partial_a X^\mu \partial^a X_\mu \\ S &= \frac{1}{4\pi\alpha'} \frac{1}{2} \int dz d\bar{z} (g^{z\bar{z}} \partial_z X^\mu \partial_{\bar{z}} X_\mu + g^{\bar{z}z} \partial_{\bar{z}} X^\mu \partial_z X_\mu) \\ S &= \frac{1}{4\pi\alpha'} \frac{4}{2} \int dz d\bar{z} \partial_z X^\mu \partial_{\bar{z}} X_\mu \\ S &= \frac{1}{2\pi\alpha'} \int d^2z \partial X^\mu \bar{\partial} X_\mu \end{aligned}$$

Where we defined $\partial_z = \partial, \partial_{\bar{z}} = \bar{\partial}$ and $d^2z = dz d\bar{z}$