University of California at San Diego – Department of Physics – Prof. John McGreevy

Physics 215B QFT Winter 2026
Assignment 3 – Solutions

Due 11:59pm Monday, January 26, 2026

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1. **Brain-warmer.**

   Use the Clifford algebra to show that in 3+1 dimensions

   $$\gamma^\mu \left(x\not{p} + m\right) \gamma_\mu = -2x\not{p} + 4m$$

   where as usual $\not{p} \equiv p^\mu \gamma_\mu$. This identity will be useful in the numerator of the electron self-energy.

2. **An example of renormalization in classical physics.**

   Consider a classical scalar field in $D+2$ spacetime dimensions coupled to an *impurity* (or defect or brane) in D dimensions, located at $X = (x^\mu, 0, 0)$. Suppose the field has a self-interaction which is localized on the defect. For definiteness and calculability, we'll consider the simple (quadratic) action

   $$S[\phi] = \int d^{D+2}X \left(\frac{1}{2}\partial_M \phi(X) \partial^M \phi(X) + \frac{1}{2} g \delta^2(\vec{x}_\perp)\phi^2(X) \right).$$

   Here $X^M = (x^\mu, x^i_\perp)$, $\mu = 0..D-1$, $i = 1, 2$, *i.e.* $x_\perp$ are coordinates transverse to the impurity.

   This example is from this paper by Goldberger and Wise.

   (a) What is the mass dimension of the coupling g? This is why I picked a codimension[1]-two defect.

   (b) Find the equation of motion for ϕ. Where have you seen an equation like this before?

   It's the Schrödinger equation for a particle in a 2d delta function potential.

   (c) We will study the propagator for the field in a mixed representation:

   $$G_k(x_\perp, y_\perp) \equiv \langle \phi(k, x_\perp)\phi(-k, y_\perp)\rangle = \int d^D z\, e^{\mathbf{i}k_\mu z^\mu} \langle \phi(z, x_\perp)\phi(0, y_\perp)\rangle$$

   – *i.e.* we go to momentum space in the directions in which translation symmetry is preserved by the defect. (I will drop the $\perp$ subscripts in what

[1] An object whose position requires specification of p coordinates has codimension p.

follows.) Find and evaluate the diagrams contributing to $G_k(x,y)$ in terms of the free propagator $D_k(x,y) \equiv \langle \phi(k,x)\phi(-k,y)\rangle_{g=0}$. (We will not need the full form of $D_k(x,y)$.) Note that there are no loop diagrams, and in this sense, all the physics here is classical. Sum the series.

I found it convenient to do this problem in Euclidean spacetime, so G and D are Euclidean propagators.

The euclidean path integral is of the form $\int D\phi e^{-S_0} e^{-V}$ where S_0 is the kinetic term and $V = \int d^{D+2}x \delta^2(x_\perp)\frac{1}{2} g\phi^2$. If we work in real time, the interaction vertex would be a factor of $-\mathbf{i}g\delta^{(2)}(x)$. If we work in euclidean time, the two-point vertex is $-g\delta^{(2)}(x)$, and no $\mathbf{i}$s will appear. From the sum of diagrams of the form (just as if we had done perturbation theory in the mass)

— + —x— + —x—x— + —x—x—x— ...

we find a geometric series

$$
\begin{aligned}
G_k(x,y) = D_k(x,y) &- g\int d^2z_1 D_k(x,z_1)\delta^{(2)}(z_1)D_k(z_1,y) \\
&+ (-g)^2 \int d^2z_1 \int d^2z_2 D_k(x,z_1)\delta^{(2)}(z_1)D_k(z_1,z_2)\delta^{(2)}(z_2)D_k(z_2,y) + \cdots \\
&= D_k(x,y) - gD_k(x,0)D_k(0,y) + (-g)^2 D_k(x,0)D_k(0,0)D_k(0,y) \\
&\quad + (-g)^3 D_k(x,0)D_k(0,0)^2 D_k(0,y) + \cdots \\
&= D_k(x,y) - gD_k(x,0)\left(1 - gD_k(0,0) + (-g)^2 D_k(0,0)^2 + \cdots\right) D_k(0,y) \\
&= D_k(x,y) - \frac{g}{1+gD_k(0,0)} D_k(x,0)D_k(0,y).
\end{aligned}
$$

(d) You should find that your answer to part 2c depends on $D_k(0,0)$, which is divergent. This divergence arises from the fact that we are treating the defect as infinitely thin, as a pointlike object – the δ^2-function in the interaction involves arbitrarily short wavelengths. In general, as usual, we must really be agnostic about the short-distance structure of things. To reflect this, we introduce a regulator. For example, we can replace the Fourier representation of (the Euclidean) $D_k(0,0)$ with the cutoff version

$$
D_k(0,0;\Lambda) = \int_0^{\Lambda} đ^2 q \frac{e^{\mathbf{i}q\cdot 0}}{k^2+q^2}. \tag{1}
$$

Do the integral.

Note that the formula (1) would need an extra factor of $\mathbf{i}$ if we were working in real time (in which case the interaction vertex would be $-\mathbf{i}g\delta^2(x)$, and the $\mathbf{i}$s would eat each other).

$$D_k(0,0;\Lambda) = \int_0^\Lambda đ^2 q \frac{e^{\mathbf{i}q\cdot 0}}{k^2+q^2} = \frac{1}{4\pi}\log\frac{\Lambda^2+k^2}{k^2} \stackrel{\Lambda \gg k}{=} \frac{1}{4\pi}\log\frac{\Lambda^2}{k^2}.$$

These dimensions whose momenta we're integrating here are spacelike, so there's no need for any Wick rotation.

(e) Now we renormalize. We will let the *bare coupling* g (the one which appears in the Lagrangian, and in the series from part 2c) depend on the cutoff $g = g(\Lambda)$. We wish to eliminate $g(\Lambda)$ in our expressions in favor of some measurable quantity. To do this, we impose a renormalization condition: choose some reference scale μ, and demand that[2]

$$G_\mu(x,y) \stackrel{!}{=} D_\mu(x,y) - g(\mu) D_\mu(x,0) D_\mu(0,y). \tag{2}$$

This equation defines $g(\mu)$, which we regard as a physical quantity. Show that (2) is satisfied if we let the bare coupling be $g(\Lambda) = g(\mu)Z$, with

$$Z = \frac{1}{1 - \frac{g(\mu)}{4\pi}\ln\left(\frac{\Lambda^2}{\mu^2}\right)}.$$

(f) Find the beta function for g,

$$\beta_g(g) \equiv \mu \frac{dg(\mu)}{d\mu},$$

and solve the resulting RG equation for $g(\mu)$ in terms of some initial condition $g(\mu_0)$. Does the coupling get weaker or stronger in the UV?

You may be bothered that we previously defined the beta function as $\Lambda\partial_\Lambda g(\Lambda)$, in terms of the cutoff dependence. In a classically scale-invariant theory, the dependence on Λ and μ is very closely tied together, since there are no other scales in the problem.

Solving for $g(\Lambda)$ gives

$$g(\Lambda) = \frac{g(\mu)}{1 - \frac{g(\mu)}{4\pi}\log\frac{\Lambda^2}{\mu^2}}.$$

Then

$$\beta_g(g) = \frac{g(\Lambda)}{\left(1 - \frac{g(\Lambda)}{4\pi}\log\frac{\Lambda^2}{\mu^2}\right)^2}\frac{g(\Lambda)}{2\pi} = \frac{g^2(\mu)}{2\pi} = \frac{g^2}{2\pi}.$$

The solution is

$$g(\mu) = \frac{g(\mu_0)}{1 - \frac{g(\mu_0)}{2\pi}\log\frac{\mu}{\mu_0}}$$

[2] Note that if we worked in real time, there would be an extra $\mathbf{i}$ in front of the second term on the RHS.

which grows with μ. Something bad happens when the denominator vanishes:

$$1 = \frac{g(\mu_0)}{2\pi} \log \frac{\mu_\star}{\mu_0}.$$

This scale $\mu_\star$ where the coupling blows up is called a *Landau pole*.

3. **Scale invariance in QFT in $D = 0 + 0$, part 1.** [I got this problem from Frederik Denef.]

   A nice realization of QFT in $0 + 0$ dimensions is the statistical mechanics of a collection of non-interacting particles. The canonical partition function for a single particle (moving in one dimension) is

   $$Z = \int đP dX e^{-\beta H} \propto \sqrt{T} Z_V(T) \tag{3}$$

   with $H = \frac{P^2}{2} + V(X)$ and $T = 1/\beta$. The momentum integral is Gaussian and we can just do it. The partition function of N non-interacting indistinguishable particles is then $Z^N/N!$, which just multiplies the energy $U = T^2 \partial_T \log Z$ by a factor of N, so we don't miss anything by focussing on the single particle.

   Let's consider the case

   $$V(X) = aX^2 + bX^4 + cX^6 \tag{4}$$

   and figure out the important features of the temperature dependence of the thermodynamic quantities by scaling arguments.

   (a) Assuming $a \neq 0, b \neq 0, c \neq 0$, find the behavior of the thermal energy U and the heat capacity $C = \partial_T U$ in the limit $T \to 0$ and in the limit $T \to \infty$ using scaling arguments. Which parts of the potential determine the respective limiting behaviors?

   First, to understand the low-temperature behavior, let $x \equiv X/\sqrt{T}$, so that

   $$Z_V = \int dX e^{-V(X)/T} = T^{1/2} \int dx e^{-(ax^2+bx^4T+cx^6T^2)} = T^{1/2} \underbrace{\int dx e^{-ax^2}}_{\text{some number}} \underbrace{e^{-(bx^4T+cx^6T^2)}}_{\stackrel{T\to 0}{\to} 1}. \tag{5}$$

   Therefore, $Z \stackrel{T\to 0}{\propto} T^{1/2+1/2}$. In this case the quadratic term is most important. If $Z \propto T^\alpha$ then $U = \alpha T$ and $C = \alpha$, so here $C = 1$. To understand the high-temperature behavior, let $y \equiv X/T^{1/6}$ so that

   $$Z_V = \int dX e^{-V(X)/T} = T^{1/6} \int dy e^{-(cy^6+ax^2/T^{2/3}+bx^4/T^{2/3})} = T^{1/6} \int dy e^{-cy^6} \underbrace{e^{-(ax^2/T^{2/3}+bx^4/T^{2/3})}}_{\stackrel{T\to\infty}{\to} 1}. \tag{6}$$

So at high temperatures $C \to \frac{1}{2} + \frac{1}{6}$. At high temperature, the particle can explore the whole potential and the highest power in the potential is what matters.

(b) If some of the couplings a, b, c vanish, the low or high temperature scaling behavior may change. For example, what is the heat capacity at low temperature when $a = 0, b \neq 0$?

In this case, the quartic term dominates and $Z_V \sim T^{1/4}$ and $C = 3/4$.

A word about notation: the symbol $\sim$ is often used by physicists to indicate a scaling relationship, where the constant prefactors are neglected. The relation we derive here for C however is an equality in the relevant regime of temperatures – the constant is the thing that matters.

(c) When b is sufficiently large (and $a \neq 0, c \neq 0$), there will be an intermediate temperature regime over which the heat capacity is again constant, but different from the low- and high-temperature limits. What is this heat capacity?

Same as the previous part.

(d) In general, we can think of the change of C with T as a kind of classical renormalization group (RG) flow, interpolating between 'fixed points' where C becomes constant. In general, these fixed points correspond to potentials $V(X)$ with a scaling symmetry $V(\lambda^\Delta X) = \lambda V(X)$ for some choice of scaling dimension Δ of X. What is the heat capacity for a fixed point with scaling dimension Δ for X?

$$Z_V = \int dX e^{-V(X)/T} = \int dX e^{-V(T^{-\Delta}X)} = T^\Delta \underbrace{\int dx e^{-V(x)}}_{\text{indep of } T} \tag{7}$$

with $x \equiv T^{-\Delta}X$. So $Z \propto T^{1/2+\Delta}$ and

$$C = \Delta + \frac{1}{2}. \tag{8}$$

(e) Borrowing more language of the renormalization group, we can classify deformations $\delta V(X) = \epsilon X^m$ of a fixed point $V(X) \propto X^{2n}$ as irrelevant, marginal, or relevant, depending on whether the deformation becomes dominant or negligible in the IR limit, *i.e.* in the limit of low T. Here and below ϵ can take on any value, not necessarily small. Restricting to deformations with an $X \to -X$ symmetry, what are the relevant and irrelevant deformations of $V(X) = X^{2n}$? (Note that a deformation $\delta V = \epsilon X^{2n}$ can be absorbed into a redefinition of X, which does not change the heat capacity.)

Lower powers than $2n$ are relevant, higher powers are irrelevant.

(f) The T-dependence of correlation functions (here, expectation values of powers of X) at fixed points is also determined by the scaling properties. What is the T-dependence of $\langle X^k \rangle$ at a fixed point $V(X) = X^{2n}$?

$$\langle X^k \rangle = \frac{\int dX X^k e^{-V(X)/T}}{Z_V} = \frac{T^{\Delta(1+k)} \int dx x^k e^{-V(x)}}{T^\Delta} \propto T^{\Delta k}$$

where $\Delta = \frac{1}{2n}$ is the scaling dimension of X.

(g) Non-polynomial $V(X)$ can be considered as well. For example, what is the heat capacity at small and large T for $V(X) = (1 + X^2)^{1/n}$?

Since this function still grows at large X, the high-temperature behavior is dominated by the large-X behavior where $V(X) \sim X^{2/n}$, so $\Delta = n/2$ and $C = \frac{n+1}{2}$. At low temperature, we Taylor expand in small X to find $V(X) \sim 1 + X^2/n$ and find $\Delta = 1/2$ and $C = 1$, where we used (8).