Holographic non-Fermi-liquid fixed points

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Techniques arising from string theory can be used to study assemblies of strongly interacting fermions. Via this ‘holographic duality’, various strongly coupled many-body systems are solved using an auxiliary theory of gravity. Simple holographic realizations of finite density exhibit single-particle spectral functions with sharp Fermi surfaces, of a form distinct from those of the Landau theory. The self-energy is given by a correlation function in an infrared (IR) fixed-point theory that is represented by a two-dimensional anti de Sitter space (AdS2) region in the dual gravitational description. Here, we describe in detail the gravity calculation of this IR correlation function.

Keywords: gauge/gravity duality; non-Fermi liquids; black holes

1. Introduction

The metallic states that we understand well are described by Landau Fermi-liquid theory. This is a free stable renormalization group (RG) fixed point [1–3] (modulo the Bardeen–Cooper–Schrieffer (BCS) instability that sets in at parametrically low temperatures). Landau quasi-particles manifest themselves as (a surface of) poles in single-fermion Green’s functions at $k_\perp \equiv |k| - k_F = 0$,

$$G_R(\omega, k) = \frac{Z}{\omega - v_F k_\perp + i\Gamma} + \cdots,$$

where the dots represent incoherent contributions. Landau quasi-particles are long lived: their width is $\Gamma \sim \omega^*\omega$, where $\omega^*(k)$ is the real part of the location of the pole. The residue $Z$, their overlap with the external electron, is finite on Fermi surfaces, and the spectral density becomes arbitrarily sharp there,

$$A(\omega, k) \equiv \frac{1}{\pi} \text{Im} \, G_R(\omega, k) \xrightarrow{k_\perp \to 0} Z\delta(\omega - v_F k_\perp).$$

Thermodynamical and transport behaviour of the system can be characterized in terms of these long-lived quasi-particles.

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Non-Fermi-liquid metals exist, but are mysterious; the ‘strange-metal’ phase of optimally doped cuprates is a notorious example. For example, data from angle-resolved photoemission (ARPES) in this phase (see [4] and references therein) indicate gapless modes—the spectral function \( A(\omega, k) \) exhibits non-analyticity at \( \omega \sim 0, k \sim k_F \)—around some \( k_F \) with width \( \Gamma(\omega_*), \omega_* \), and vanishing residue \( Z \rightarrow 0 \). Furthermore, the system’s resistivity exhibits a linear temperature dependence in sharp contrast to the quadratic dependence of a Fermi liquid. These anomalies, along with others, suggest that in the strange-metal phase, there is still a sharp Fermi surface, but no long-lived quasi-particles.

Known field-theoretical examples of non-Fermi liquids that also exhibit the phenomenon of a Fermi surface without long-lived quasi-particles include Luttinger liquids in 1+1 dimensions [5], and a free fermion gas coupled to some gapless bosonic excitations, which can be either a transverse gauge field or certain-order parameter fluctuations near a quantum critical point (see [6–25]). Neither of these, however, is able to explain the behaviour of a strange-metal phase. The former is specific to \((1+1)\)-\(d\) kinematics. In the latter class of examples, the influence of gapless bosons is mostly along the forward direction, and is not enough, for example, to account for the linear temperature dependence of the resistivity.

Here, we summarize recent findings of a class of non-Fermi liquids, some of which share similar low-energy behaviour to those of a strange-metal phase, using the holographic approach ([26–30]; see also [31–33]). One of the most intriguing aspects of these systems is that their low energy behaviour is controlled by an infrared (IR) fixed point, which exhibits non-analytical behaviour only in the time direction. In particular, single-particle spectral function and charge transport can be characterized by the scaling dimension of the fermionic operator in the IR fixed point.

We should emphasize that, unlike efforts to understand strange metals, which begin from the Hubbard model such as those of Leigh & Phillips [34] and Leigh et al. [35], at a microscopic level, the systems we consider differ very much from the electronic systems underlying strange metals: they are translationally invariant and spherically symmetric; they typically involve a large number of fermions, scalars and gauge fields, characterized by a parameter \( N \), and we have to work with the large-\( N \) limit; at short distances, the systems approach a relativistic conformal field theory (CFT). Nevertheless, the similarity of their low-energy behaviour to that of a strange metal is striking and may not be an accident. After all, the key to our ability to characterize many-body systems has often been universalities of low-energy physics among systems with different microscopics.

We should also mention that the holographic non-Fermi liquids we describe here probably reflect some intermediate-scale physics rather than genuine ground states. The systems in which they are embedded can have various superconducting [36–38] or magnetic [41] instabilities, just as in real-life condensed-matter systems.

In §2, we give a lightning introduction to holographic duality. Sections 3 and 4 describe the structure of the fermionic response in the simplest holographic realization of finite density. Section 5 is the core of this paper, where we describe

\[1\] For a review, see [39,40].
in detail the calculation of the self-energy. In the final sections, we comment on transport, the superconducting state and a useful cartoon of the mechanism that kills the quasi-particles.

2. Holographic duality

The study of black holes has taught us that a quantum theory of gravity has a number of degrees of freedom that is sub-extensive. In general, one expects a gravitational system in some volume to be describable in terms of an ordinary quantum system living on the boundary. For a certain class of asymptotic geometries, these boundary degrees of freedom have been precisely identified \cite{42–44}. For a review of this holographic duality in the present spirit, see \cite{45–49}.

The basic example of the duality is the following. A theory of gravity in \((d+1)\)-dimensional anti de Sitter space \((\text{AdS}_{d+1})\) is a CFT in \(d\) spacetime dimensions. Here, by ‘is’ we mean that the observables are in one-to-one correspondence. AdS space,

\[
    ds^2 = \frac{r^2}{R^2} (-dt^2 + d\mathbf{x}^2) + R^2 \frac{dr^2}{r^2},
\]

can be viewed as a collection of copies of Minkowski space \(\mathbb{R}^{d,1}\) (whose isometries are the Poincaré group) of varying ‘size’, parametrized by a coordinate \(r\). The isometries of \(\text{AdS}_{d+1}\) are precisely the conformal group in \(d\) spacetime dimensions. The ‘boundary’ where the dual field theory lives is at \(r = \infty\) in these coordinates (figure 1).

The extra (‘radial’) dimension can be considered as corresponding to the resolution scale for the dual theory. The evolution of the geometry (and the fields propagating therein) along this radial direction represent the RG flow of the dual field theory. The ‘AdS radius’, \(R\) in equation (2.1), is a dimensionful parameter that, roughly, encodes the rate of RG flow.

A practical consequence of this duality is the following equation for the field-theory partition function:

\[
    Z_{\text{QFT}}[\text{sources}] = Z_{\text{quantum gravity}}[\text{boundary condition at } r \to \infty] \\
    \approx e^{-N^2 S_{\text{bulk}}[\text{boundary condition at } r \to \infty]|_{\text{extremum of } S_{\text{bulk}}}}. \tag{2.2}
\]

The r.h.s. of the first line is a partition function of quantum gravity, a general understanding of which is still lacking. The boundary conditions at \(r \to \infty\) specify the sources in the quantum field theory (QFT) partition function. In the second line, the classical approximation to the quantum gravity partition function has been made via a saddle-point approximation. The dimensionless parameter \(N^2\), which makes such an approximation possible, is the inverse of the Newton constant in units of curvature radius \(R\), and is large in the classical limit. Through the duality, \(N^2\) is mapped to the number of degrees of freedom per point in the dual free-field theory. In the best-understood examples, the classical nature of the gravity theory arises from large-\(N\) factorization in the ‘t Hooft limit of a gauge theory.
Fields in AdS\(_{d+1}\) correspond to operators in CFT; the mass of the field determines the scaling dimension of the operator. The boundary conditions on a bulk field specifies the coefficient of the corresponding operator in the field-theory action. For example, since the bulk theory is a gravitational theory, the bulk spacetime metric \(d s^2 = g_{\mu\nu} \, dx^\mu \, dx^\nu\) is dynamical. The boundary value of bulk metric \(\lim_{r \to \infty} g_{\mu\nu}\) is the source for the field-theory stress-energy tensor \(T^{\mu\nu}\). So, we can rewrite equation (2.2) as

\[
\langle e^{\phi^{(0)} f} \rangle \approx e^{-N^2 S_{\text{bulk}}(\phi_{\text{a}} \overset{r \to \infty}{\longrightarrow} \phi_{\text{a}}^{(0)})} \big|_{\text{extremum of } S_{\text{bulk}}}. \tag{2.3}
\]

We can change the field-theory action just by changing the boundary conditions on the bulk fields. Different couplings in the bulk action then correspond to entirely different field theories. For example, changing the bulk Newton constant (in units of curvature radius) is accomplished by changing the parameter \(N^2\); but this is the number of degrees of freedom per point of the field theory.

In addition to large \(N^2\), useful calculation requires the background geometry to have small curvature. This is because otherwise we cannot reliably approximate the bulk action as

\[
S = \int d^{d+1}x \sqrt{g} \left( \mathcal{R} + \frac{d(d + 1)}{R^2} + \cdots \right), \tag{2.4}
\]

where \(\mathcal{R}\) is the Ricci scalar, and the ellipsis indicates terms with more powers of the curvature (and hence more derivatives). One necessary condition for this is that the AdS radius \(R\) be large compared to the energy scale set by the string tension. Given how it appears in equation (2.4), this requirement is a bulk version of the cosmological constant problem—the vacua with large cosmological constant require more information about the full string theory. The largeness of the dual geometry implies that the dual QFT is strongly coupled. Circumstantial evidence for this statement is that in QFTs that are weakly coupled, we can calculate and tell that there is not a large extra dimension sticking out. The fact that certain strongly coupled field theories can be described by classical gravity on some auxiliary spacetime is extremely powerful, once we believe it.
before proceeding to apply this machinery to non-Fermi-liquid metals, we
pause here to explain the reasons that give us enough confidence in these rather
odd statements to try to use them to do physics. The reasons fall roughly into
three categories:

— Many detailed checks have been performed in special examples. These
checks have been done mainly in relativistic gauge theories (where the fields
are $N \times N$ matrices) with extra symmetries (conformal invariance and
supersymmetry). The checks involve the so-called ‘Bogolmnyi–Prasad–
Sommerfeld quantities’ (which are the same at weak coupling and strong
coupling or which can be computed as a function of the coupling),
integrable techniques, and more recently some numerics. We will not
discuss any of these kinds of checks here because they involve calculations
of quantities that only exist in specific models; neither the quantities nor
the models are of interest here.

— The holographic correspondence unfailingly gives sensible answers for
physics questions. This includes rediscoveries of many known physical
phenomena, some of which are quite hard to describe otherwise, e.g.
colour confinement, chiral symmetry breaking, thermo, hydro, thermal
screening, entanglement entropy, chiral anomalies, superconductivity, etc.
The gravity limit, when valid, says which are the correct variables, and
gives immediate answers to questions about thermodynamics, transport,
RG flow, etc. in terms of geometric objects.

— If we are bold, the third class of reasons can be called ‘experimental checks’.
These have arisen from applications to the quark-gluon plasma (QGP)
produced at the relativistic heavy ion collider. Holographic calculations
have provided a benchmark value for the viscosity of such a strongly
interacting plasma, and have provided insight into the behaviour of hard
probes of the medium, and of the approach to equilibrium.

As an illustration of the manner in which the correspondence solves hard
problems by simple pictures, we offer the following (figure 2). The bulk geometry
is a spectrograph separating the theory by energy scales. The geometry dual to
a general Poincaré-invariant state of a QFT takes the form

$$ds^2 = w(r)^2(-dt^2 + d\mathbf{x}^2) + R^2 \frac{dr^2}{r^2}.$$  (2.5)

For the gravity dual of a CFT, the bulk geometry goes on forever, and the ‘warp
factor’ $w(r) = r/R \rightarrow 0$. On the other hand, in the gravity dual of a model with a
gap, the geometry ends smoothly, warp factor $w(r)$ has a non-zero minimum
value. If the IR region of the geometry is missing, there are no low-energy
excitations and hence an energy gap in the dual field theory.

(b) Finite density

A basic question for the holographic description is how to describe a finite
density. To approach this question, we introduce a minimal set of necessary
ingredients in the bulk model. The fact that any local QFT has a stress tensor
$T^{\mu\nu}$ means that we should have a dynamical metric $g_{\mu\nu}$ in bulk. As a proxy for

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Figure 2. (a) A comparison of the geometries associated with a CFT and (b) with a system with a mass gap. (Online version in colour.)

Consider any relativistic CFT with a gravity dual and a conserved $U(1)$ symmetry. The discussion goes through for any $d > 1 + 1$, but we focus on $d = 2 + 1$. The gravity dual is described by

$$S = \frac{1}{2\kappa^2} \int d^4x \sqrt{g} \left( R + \frac{6}{R^2} - \frac{2\kappa^2}{g_s^2} F_{\mu\nu} F^{\mu\nu} + \cdots \right). \quad (2.6)$$

The ellipsis indicates fields that vanish in the ground state, and more irrelevant couplings. This is the action we would guess based on Wilsonian naturalness, and it is what comes from string theory when we can compute it.

As a warm up, let us discuss the holographic description of finite temperature. The canonical ensemble of a QFT at temperature $T$ is described by putting the QFT on a space with periodic Euclidean time of radius $T^{-1}$. The spacetime on which the dual QFT lives is the boundary of the bulk geometry at $r \to \infty$. So, the bulk description of finite temperature is the extremum of the bulk action whose Euclidean time has a period that approaches $T^{-1}$ at the boundary. For many bulk actions, including equation (2.4), this saddle point is a black hole in AdS. This is a beautiful application of Hawking’s observation that black holes have a temperature (figure 3).

Similarly, the grand canonical ensemble is described by a periodic Euclidean time with a Wilson line $e^{i\mathcal{C}} A^{(0)} = e^{-i\partial \mu}$, where $\mathcal{C}$ is the thermal circle and $A^{(0)}$ is the background gauge field coupling to the current in question. The source $A^{(0)}$ is the boundary value of the bulk gauge field. For the minimal bulk field content introduced above, the bulk solution with these boundary conditions is the Reissner–Nördstrom (RN) black hole in AdS. We will comment below on how this conclusion is modified if we include other bulk fields, such as scalars, in equation (2.6).

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For zero temperature, the solution of equation (2.6) describing a finite density of $U(1)$ charge is

$$ds^2 = \frac{r^2}{R^2}(-f dt^2 + d\mathbf{x}^2) + R^2 \frac{dr^2}{r^2 f} \quad \text{and} \quad A = \mu \left(1 - \frac{1}{r}\right) dt,$$

(2.7)

where $f(r) = 1 + 3/r^4 - 4/r^3$ is the zero-temperature limit. This geometry has a horizon at $r = 1$ and $\mu$ is the chemical potential for the $U(1)$ symmetry.

(c) Strategy to find a Fermi surface

To look for a Fermi surface, we look for sharp features in fermionic Green’s functions at finite momentum and small frequency, following [32]. Assume that amongst the $\cdots$ in the bulk action (2.6) is

$$S_{\text{probe}}[\psi] = \int d^{d+1}x \sqrt{g} (\bar{\psi} (D - m) \psi + \text{interactions}).$$

(2.8)

The dimension and charge of the boundary fermion operator are determined by

$$\Delta = \frac{d}{2} + mR \quad \text{and} \quad q = q.$$

(2.9)

Here, we can see a kind of ‘bulk universality’: for two-point functions, the interaction terms do not matter. We can describe many CFTs (many universality classes!) by a single bulk theory. The results only depend on $q$ and $\Delta$. Some comments on the strategy are as follows.

— There are many string-theory vacua with these ingredients [50–54]. In specific examples of dual pairs (e.g. M2-branes $\leftrightarrow$ M-theory on AdS$_4 \times S^7$), the interactions and the parameters $q$ and $m$ are specified. Which sets $\{q, m\}$ are possible and what correlations there are, are not clear, and, as an expedient, we treat them as parameters.

The derivative $\mathcal{D} = D^M D_M$ contains the coupling to both the spin connection and the gauge field $D_M \equiv \partial_M + \frac{1}{4} \omega_{MAB} F^{AB} - i q_\Psi A_M$. 

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— This is a large complicated system (with a density $\rho \sim N^2$), of which we are probing a tiny part (the fermion density can be seen to scale like $\rho_f \sim N^0$).
— In general, both bosons and fermions of the dual field theory are charged under the $U(1)$ current: this is a Bose–Fermi mixture. The relative density of bosons and fermions, and whether the bosons will condense, is a complicated dynamical question. Fortunately, the gravity theory solves this problem for us; let us see what happens.

(d) Anti de Sitter space/conformal field theory prescription for spinors

To compute the retarded single-fermion Green’s function $G_R$, we must solve the Dirac equation $(\not{D} + m)\psi = 0$ in the black-hole geometry, and impose infalling boundary conditions at the horizon [55,56]. Like retarded response, falling into the black hole is something that happens, rather than unhappens. Translation invariance in $x$ and $t$ turns the Dirac equation into an ordinary differential equation in $r$. Rotation invariance allows us to set $k_i = \delta_i^1 k_1$; we can then choose a basis of gamma matrices in which the Dirac equation is block diagonal and real.\(^3\) Near the boundary, solutions behave in this basis as

$$\Phi \overset{r \to \infty}{\approx} a_\alpha r^m \begin{pmatrix} 0 \\ 1 \end{pmatrix} + b_\alpha r^{-m} \begin{pmatrix} 1 \\ 0 \end{pmatrix}. \tag{2.10}$$

From this data, we can extract a matrix of Green’s functions, which has two independent eigenvalues,

$$G_\alpha(\omega, k) = \frac{b_\alpha}{a_\alpha}, \quad \alpha = 1, 2. \tag{2.11}$$

The label $\alpha = 1, 2$ indexes a multiplicity that arises in the boundary theory as a consequence of the short-distance Lorentz invariance and will not be important for our purposes. The equation depends on $q$ and $\mu$ only through $\mu_q \equiv \mu q$. We emphasize that all frequencies $\omega$ appearing below are measured from the effective chemical potential, $m_q$. All dimensionful quantities below are quoted in units of the chemical potential.

3. A Fermi surface

The system is rotation invariant, $G$ depends on $k = |k|$ and $\omega$. Green’s function satisfies the following:

— The spectral density is positive for all $\omega$ as required by unitarity.
— The only non-analyticity in $G_R$ (for real $\omega$) occurs at $\omega = 0$. This could be anticipated on physical grounds, and is a consequence of the following argument. Consider a mode that is normalizable at the boundary, i.e. $a = 0$ in equation (2.10). This is a real boundary condition. Thus, if the equations of motion are real (which they are for real $\omega$), the wave must be real at the other end, i.e. at the horizon. Thus, it must be a sum of both infalling and

\(^3\)It is also convenient to redefine the independent variable by $\tilde{\psi} = (-\det gg^{rr})^{-1/4} \Phi$. 

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Figure 4. (a) Momentum distribution of the spectral density (MDC) for the two spin components, for \( q = 1, m = 0 \); a Fermi surface is visible in \( \text{Im } G_2 \) at \( k_F \approx 0.918528499 \). (b) The energy distribution curve (EDC) near the Fermi surface (upper curve); the real part of \( G_2 \) is also shown (lower curve). (a) MDC: \( \text{Im } G_{1,2}(\omega = -10^{-3}, k) \); (b) EDC: \( G_2(\omega, k = 0.9) \). (Online version in colour.)

Figure 5. Three-dimensional plots of (a) \( \text{Im } G_1(\omega, k) \) and (b) \( \text{Im } G_2(\omega, k) \) for \( m = 0 \) and \( q = 1 (\mu_q = \sqrt{3}) \). In (b), the ridge at \( k \gg \mu_q \) corresponds to the smoothed-out peaks at finite density of the divergence at \( \omega = k \) in the vacuum. As one decreases \( k \) to a value \( k_F \approx 0.92 < \mu_q \), the ridge develops into an (infinitely) sharp peak indicative of a Fermi surface. (Online version in colour.)

...outgoing waves at the horizon. Thus, it cannot be an infalling solution. The exception is at \( \omega = 0 \) where the infalling and outgoing conditions are each real.

— The result approaches the CFT behaviour at higher energies (\( |\omega| \gg \mu \)).

At \( T = 0 \), we find numerically a quasi-particle peak near \( k = k_F \approx 0.9185 \). The peak moves with dispersion relation \( \omega \sim k_z^z \) with \( z = 2.09 \) for \( q = 1, \Delta = 3/2 \) and \( z = 5.32 \) for \( q = 0.6, \Delta = 3/2 \). There is a scaling behaviour near the Fermi surface,

\[
G_R(\lambda k, \lambda^2 \omega) = \lambda^{-\alpha} G_R(k, \omega), \quad \alpha = 1.
\]  

It is not a Fermi liquid. The above scaling should be contrasted with the scaling in a Landau Fermi liquid, \( z = \alpha = 1 \). Finally, in the cases shown here, the residue vanishes at the Fermi surface (figures 4 and 5).
4. Low-frequency behaviour

So far, AdS/CFT is a black box that produces consistent spectral functions. To understand better where these numbers come from, we need to take apart the black box a bit.

(a) Emergent quantum criticality from geometry

At $T = 0$, the ‘emblackening factor’ in equation (2.7) behaves as $f \approx 6(r - 1)^2$, and this means that the near-horizon geometry is $\text{AdS}_2 \times \mathbb{R}^{d-1}$. Recall that AdS means conformal symmetry. The conformal invariance of this metric is emergent. We broke the microscopic conformal invariance when we put in the finite density. AdS/CFT suggests that the low-energy physics is governed by the dual IR CFT. More precisely, there is such a zero-dimensional CFT for each $k$. At small temperature $T \ll \mu$, the geometry is a black hole in $\text{AdS}_2$ times the space directions. The bulk geometry is a picture of the RG flow from the $d$-dimensional CFT to this non-relativistic CFT (figure 6).

(b) Matching calculation

At high temperatures (i.e. not small compared to the chemical potential), the retarded Green function $G_R(\omega)$ is analytic near $\omega = 0$, and can therefore be conveniently computed in series expansion [56]. In the limit of interest to us, $T \ll \mu$, expanding the wave equation in $\omega$ is delicate. This is because the $\omega$-term dominates near the horizon. We proceed using the method of matched asymptotic expansions: we find the solution (in an $\omega$-expansion) in two regions of black-hole geometry (IR and ultraviolet (UV)), and fix integration constants by matching their behaviour in the region of overlap. The region of overlap is large when $\omega, T \ll \mu$. This technique is familiar to string theorists from the brane absorption calculations [57] that led to the discovery of holographic duality. Here, this ‘matching’ can be interpreted in the QFT as RG matching between UV and IR CFTs. Note that we need only assume the existence of the IR CFT; the gravity dual lets us compute.

First we discuss the IR boundary condition: this is associated with the near-horizon region of the bulk geometry, which is $\text{AdS}_2 \times \mathbb{R}^2$. Wave equations for charged fields in $\text{AdS}_2$ turn out to be solvable. Near the boundary of $\text{AdS}_2$, the solutions for the mode with momentum $k$ are power laws with exponents $\pm \nu_k$ where

$$\nu_k \equiv R_2 \sqrt{\frac{m^2 + k^2 - q^2}{2}}. \quad (4.1)$$

This exponent determines the scaling dimension of the IR CFT operator $O_k$ to which the spinor operator with momentum $k$ flows: $\delta_k = \frac{1}{2} + \nu_k$. The retarded two-point function of the operator $O_k$ will play an important role, and is

$$G_k(\omega) = c(k) \omega^{2\nu_k}, \quad (4.2)$$

where $c(k)$ is a complex function whose calculation and explicit form is described in the next section.

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Next, we consider the low-frequency expansion in the near-boundary region of the geometry, which is associated with the UV of the field theory. The outer region solutions can be expanded in powers of $u$, and we can work in a basis where the Dirac equation is completely real. The UV data are therefore real and analytic in $u$. A basis of solutions at $u=0$ is

$$F(0)^{\pm} \approx v^{\pm}(r-1)^{\mp}, \quad (4.3)$$

where $v^{\pm}$ are certain constant spinors that will be specified in §5. Two solutions can then be constructed perturbatively in $u$,

$$F^{(0)^{\pm}} = F(0)^{\pm} + \omega F^{(1)^{\pm}} + \omega^2 F^{(2)^{\pm}} + \ldots \quad \text{and} \quad F^{(n)^{\pm}} \approx \begin{pmatrix} \frac{b^{(n)^{\pm}}}{a^{(n)^{\pm}}} + \frac{m}{r} \\ \frac{b^{(n)^{\pm}}}{a^{(n)^{\pm}}} + \frac{m}{r} \end{pmatrix}. \quad (4.4)$$

Matching these solutions to the leading and subleading solutions in the near-horizon region gives

$$\psi_a = \psi_a^+ + G(\omega) \psi_a^-, \quad (4.5)$$

where $G$ is the IR CFT Green function defined above in equation (4.2).

For any $k$, this produces a Green’s function of the form

$$G_R(\omega, k) = K \frac{b^{(0)}_+ + \omega b^{(1)}_+ + O(\omega^2) + G_k(\omega)(b^{(0)}_- + \omega b^{(1)}_- + O(\omega^2))}{a^{(0)}_+ + \omega a^{(1)}_+ + O(\omega^2) + G_k(\omega)(a^{(0)}_- + \omega a^{(1)}_- + O(\omega^2))}. \quad (4.6)$$

The overall factor $K$ will not be important in the following and we set it to unity. For generic $k$, the UV coefficient $a^{(0)}_+(k)$ is non-zero. The low-frequency expansion gives

$$G_R(\omega, k) = \frac{b^{(0)}_+}{a^{(0)}_+} + r_1 \omega + r_2 G_k(\omega) + \cdots. \quad (4.7)$$

Since $a_\pm$ and $b_\pm$ are all real, we see that the non-analytic behaviour and dissipation are controlled by IR CFT.

(c) Consequences for Fermi surfaces

Suppose there is some $k = k_F$ such that $a^{(0)}_+(k_F) = 0$ in equation (4.6). This happens if there exists a zero-energy bound state of the outer-region
Dirac equation; it could be called ‘inhomogeneous fermionic hair’ on the black hole. Near such a value of $k$,

$$G_R(\omega, k) = \frac{h_1}{k^\perp - \frac{1}{v_F} \omega - h_2 \omega 2 \nu_F}.$$  (4.8)

The coefficients $h_{1,2}$ and $v_F$ are real, UV data. This form of Green’s function correctly fits numerics near the Fermi surface. The expression (4.8) contains a lot of information about the behaviour near the Fermi surface. The result depends on the value of $\nu$ compared to $\frac{1}{2}$.

First, suppose $\frac{1}{2} > \nu > 1$; in this case, $O_{k_F}$ is irrelevant and the linear term dominates the dispersion,

$$G_R(\omega, k) = \frac{h_1}{k^\perp + \frac{1}{v_F} \omega + c_k \omega 2 \nu_F} \omega_*(k) \sim v_F k^\perp \quad \text{and} \quad Z \sim k^\perp \rightarrow 0.$$  (4.9)

The excitation at the Fermi surface is not a stable quasi-particle,

$$\frac{\Gamma(k)}{\omega_*(k)} \xrightarrow{k_\perp \rightarrow 0} \text{const.} \quad \text{and} \quad Z \propto k_{\perp}^{(2 \nu_F - 1)} k_{\perp} \rightarrow 0.$$

Finally, suppose $\nu = \frac{1}{2}$; $O_{k_F}$ is marginal. The two frequency-dependent terms compete. $v_F \propto \nu - \frac{1}{2} \rightarrow 0$, while $c(k_F)$ has a pole; they cancel and leave behind a logarithm,

$$G_R \approx \frac{h_1}{k_{\perp} + \tilde{c}_1 \ln \omega + \tilde{c}_1 \omega}, \quad \tilde{c}_1 \in \mathbb{R}, \quad c_1 \in \mathbb{C}, \quad \text{and} \quad Z \sim \frac{1}{|\ln \omega_*|} k_{\perp} \rightarrow 0.$$  (4.13)

This is the form of Green’s function proposed in a well-named phenomenological model of the electronic excitations of the cuprates near optimal doping [58].

The case $\nu = 1$ appears similar to a Landau Fermi liquid [31], though there are also logarithms in this case, and the physical origin of the quasi-particle decay is quite distinct from electron–electron interactions.

(d) Ultraviolet data: where are the Fermi surfaces?

Above, we assumed $a^{(0)}_+(k_F) = 0$. This happens at $k = k_F$ such that there exists a normalizable, incoming solution at $\omega = 0$. From a relativist’s viewpoint, the black hole acquires inhomogeneous fermionic hair. By a change of variables,
Figure 7. Shown here is a sequence of Schrödinger potentials for the scalar wave equation in the RN black hole. The horizontal axis is a ‘tortoise’ coordinate $s$ that makes the wave equation into a one-dimensional Schrödinger problem. The role of the energy in the Schrödinger problem is played by $-k^2$. The dashed curve is a cartoon of the bound-state wave function at $\omega = 0$ with energy $-k_F^2$; the solid curve that becomes horizontal at large negative $s$ (the IR region) is the associated Schrödinger potential for $\omega = 0$. As $\omega$ increases from zero, the potential develops a well in the IR region (the other solid curves), into which the bound state can tunnel. The width of the barrier is $\Delta s \sim -2\ln|\omega|$, and the height is $n^2_kF$; hence, the tunnelling amplitude that determines the decay rate of the Fermi-surface bound state is $e^{-\text{area}} \sim \omega^{2\nu_F}$. (Online version in colour.)

This problem can be translated into a bound-state problem in one-dimensional quantum mechanics. The relevant Schrödinger potentials are shown in figure 7. A few cases are worth noting. For $k > qe_d$, the potential is always positive and there can be no bound state; there is therefore no Fermi surface in this regime. Below a certain momentum, $k < k_{\text{osc}} \equiv \sqrt{(qe_d)^2 - m^2}$, the potential develops a singular well near the horizon $V(x) \sim \alpha/\tau^2$, with $\alpha < -\frac{1}{4}$. We refer to this as the ‘oscillatory regime’; it is associated with particle production in the AdS2 region of the geometry. In this case, the exponent $\nu$ is imaginary, and hence $G_R$ periodic in log $\omega$. We note that this behaviour is quite independent of the Fermi surface behaviour.

In the intermediate regime $k \in (qe_d, k_{\text{osc}})$, the potential develops a well, indicating the possible existence of a zero-energy bound state. The locations of these bound states can easily be determined numerically and are shown in figure 8. Recently, such states have been shown to exist also for fields near black holes in asymptotically flat space [59].

(e) Summary

The location of the Fermi surface is determined by short-distance physics analogous to band structure—one must find a normalizable solution of the $\omega = 0$ Dirac equation in full black holes. The low-frequency scaling behaviour near the Fermi surface however is universal; it is determined by a near-horizon region and in particular the IR CFT $G$. Depending on the dimension of the operator in the IR CFT, we find Fermi-liquid behaviour (but not Landau) or non-Fermi-liquid behaviour. From the bulk point of view, the quasi-particles decay by falling into
Figure 8. The values of \( q \) and \( k \) at which poles of Green’s function \( G_2 \) occur are shown by solid lines for \( m = -0.4, 0, 0.4 \). The dotted lines represent zeros of \( G_2 \), and hence poles of Green’s function \( \tilde{G}_1 \) for the opposite sign of the mass. The shaded areas indicate the oscillatory region, where \( \nu \in \mathbb{R} \). (Online version in colour.)

Figure 9. As we vary the mass and charge of the spinor field, we find the following behaviour. In the white region, there is no Fermi surface. In the region \( \nu < \frac{1}{2} \), we find non-Fermi-liquid behaviour. In the remainder of the parameter space, there is a stable quasi-particle. (Online version in colour.)

the black hole. The rate at which they fall in is determined by their effective mass (which by the correspondence determines the exponent \( \nu \)) in the near-horizon region of the geometry (figure 9).

5. Correlation functions in the two-dimensional anti de Sitter space infrared conformal field theory

Here, we give a derivation of the retarded Green function \( \mathcal{G}(\omega) \) for a charged scalar and spinor in AdS_2. We then describe the finite-temperature case. Then, we briefly discuss the generalization to other IR geometries.

The AdS_2 background metric and gauge field are given by

\[
\mathrm{d}s^2 = \frac{R_2^2}{\zeta^2} (-\mathrm{d}r^2 + \mathrm{d}\zeta^2) \quad \text{and} \quad A = \frac{e_d}{\zeta} \mathrm{d}r. \quad (5.1)
\]
We describe below how the IR CFT correlators relevant to the calculations described above can be extracted from the behaviour of fields in this background. One may worry that the non-zero profile of the gauge field breaks the conformal invariance. In fact, conformal invariance is restored by acting simultaneously with a gauge transformation, as in [60].

(a) Scalar

We consider the following quadratic scalar action:

\[-d^2 x \sqrt{-g} [g^{ab} (\partial_a + iqA_a) \phi^* (\partial_b - iqA_b) \phi + m^2 \phi^* \phi] \]  

in the background (5.1). In frequency space \( \phi(\tau, \zeta) = e^{-i\omega \tau} \phi(\omega, \zeta) \), the equation of motion for \( \phi \) can be written as

\[-\partial_\zeta^2 \phi + V(\zeta) \phi = 0, \]  

with

\[V(\zeta) = \frac{m^2 R^2}{\zeta^2} - \left( \omega + \frac{q e_d}{\zeta} \right)^2. \]  

Note that \( \omega \) can be scaled away by redefining \( \zeta \), reflecting the scaling symmetry of the background solution. Equation (5.3) can be solved exactly; the two linearly independent solutions are

\[\phi = c_{\text{out}} W_{\nu,\mu}(z(-2i \omega \zeta)) + c_{\text{in}} W_{-\nu,\mu}(2i \omega \zeta), \]  

where \( W_{\lambda,\mu}(z) \) is the Whittaker function, and for scalars \( \nu \equiv \sqrt{m^2 R^2 - q^2 e_d^2 + \frac{1}{4}} \). Of these, the function multiplying \( c_{\text{in}}, W_{\nu,\mu}(z(-2i \omega \zeta)) \sim e^{i \omega \zeta} \zeta^{i q e_d} \), is ingoing at the horizon.

Near the boundary of the AdS\(_2\) region, the scalar solution behaves as

\[\phi \approx A \zeta^{-\nu+1/2} + B \zeta^{\nu+1/2}. \]  

The retarded scalar function of the IR CFT is then\(^5\)

\[G_R(\omega) = \frac{B}{A} e^{-i \omega \tau} \Gamma(-2\nu) \Gamma\left(\frac{1}{2} + \nu - i e_d\right) (2\omega)^{2\nu}. \]  

The advanced function is given by

\[G_A(\omega) = e^{i \omega \tau} \Gamma(-2\nu) \Gamma\left(\frac{1}{2} + \nu + i e_d\right) (2\omega)^{2\nu}. \]  

Now consider a charged scalar field on AdS\(_2 \times \mathbb{R}^{d-1}\). In momentum space, the equation of motion reduces to equation (5.4) with \( m^2 \) replaced by

\[m_k^2 \equiv k^2 \frac{R^2}{r_*^2} + m^2, \quad k^2 = |k^2|. \]  

\(^5\)Note that we define the scalar Green function to be \( B/A \) without the prefactor of \( 2\nu \) first emphasized in [61].
Thus, the retarded function is given by equation (5.7) with \( m^2 \) replaced by \( m_k^2 \).

\[ \text{(b) Spinor} \]

We consider the following quadratic action for a spinor field \( \psi \) in the geometry (5.1):

\[
S_{\text{spinor}} = \int d^2x \sqrt{-g} \left( \bar{\psi} \Gamma^a D_a \psi - m \bar{\psi} \psi + i \tilde{m} \bar{\psi} \Gamma \psi \right);
\]

the last term in this action is a parity-violating mass term that, in our application, will be related to momentum in \( \mathbb{R}^{d-1} \). We make the following choice of Gamma matrices, chosen to be compatible with the choice made in [26], with equation (5.1) arising as the near-horizon limit:

\[ \Gamma^z = \sigma^3, \quad \Gamma^x = i \sigma^1 \quad \text{and} \quad \Gamma = -\sigma^2. \]

Then, the equations of motion for \( \psi \) can be written as

\[
(\zeta \partial_\zeta - i \sigma^3(\zeta \omega + q e_d)) \tilde{\Phi} = R_2(m \sigma^2 + \tilde{m} \sigma^1) \tilde{\Phi}.
\]

Here,

\[
\tilde{\Phi} \equiv \left( \frac{\bar{y}}{z} \right) \equiv \frac{1}{\sqrt{2}} (1 + i \tilde{\sigma}^1)(-gg^{\nu \nu})^{-1/4} \psi.
\]

The general solution to this equation is

\[
\tilde{\Phi}(\zeta) = \zeta^{-1/2} \left[ c_{\text{out}} W_{-(\sigma^3/2)-i q e_d, \nu}(2i \omega \zeta) \begin{pmatrix} \tilde{m} - im \\ -1 \end{pmatrix} \\
+ c_{\text{in}} W_{(\sigma^3/2)+i q e_d, \nu}(-2i \omega \zeta) \begin{pmatrix} -1 \\ m + im \end{pmatrix} \right],
\]

where again \( W \) is a Whittaker function, and where, for spinors, \( \nu = \sqrt{m^2 R_2^2 - q^2 e_d^2} \).

The notation \( \sigma^3 \) in the index of the Whittaker function indicates \( \pm 1 \) when acting on the top/bottom component of the spinor.

The field that matches to the spinor field in the outer region is \( \Phi = \frac{1}{\sqrt{2}} (1 - i \sigma^1) \tilde{\Phi} \). Near the boundary of AdS_2, \( \zeta \rightarrow 0 \), the AdS_2 Dirac equation in this basis becomes

\[
\zeta \partial_\zeta \Phi = U \Phi, \quad U = \begin{pmatrix} m R_2 & \tilde{m} R_2 - q e_d \\ \tilde{m} R_2 + q e_d & -m R_2 \end{pmatrix}.
\]

The asymptotic behaviour of \( \Phi \) near the boundary of the AdS_2 region is as in equation (4.3),

\[
\Phi = A v_-(\zeta^{-p}(1 + O(\zeta)) + B v_+ \zeta^p(1 + O(\zeta)),
\]

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where $v_{\pm}$ are real eigenvectors of $U$ with eigenvalues $\pm\nu$. The relative normalization of $v_+$ and $v_-$ is a convention, which affects the normalization of the AdS$_2$ Green functions. We will choose $v_\pm$ to be given by

$$v_{\pm} = \left( \frac{mR_2 \pm \nu}{\tilde{m}R_2 + qe_d} \right). \tag{5.17}$$

With the normalization convention in equation (5.17), the bottom components of $v_\pm$ are equal, and so $A$ and $B$ can be extracted from the asymptotics of $z$. The relative normalization of the eigenspinors in equation (5.17) affects the answer for the AdS$_2$ Green function; $v_{\pm} \to \lambda_{\pm} v_{\pm}$ takes $G_R \to (\lambda_+/\lambda_-) G_R$. This rescaling does not, however, affect the full Green function $G_R$ computed by the matching procedure.

One then finds that the retarded Green function of the operator coupling to $y$ (the upper component of $\Phi$) can be written as

$$G_R(\omega) = e^{-i\pi\nu} \frac{\Gamma(-2\nu)\Gamma(1 + \nu - iqe_d)}{\Gamma(2\nu)\Gamma(1 - \nu - iqe_d)} \cdot \frac{(m - i\tilde{m})R_2 - iqe_d - \nu}{(m - i\tilde{m})R_2 - iqe_d + \nu} (2\omega)^{2\nu}. \tag{5.18}$$

The advanced function is given by

$$G_A(\omega) = e^{i\pi\nu} \frac{\Gamma(-2\nu)\Gamma(1 + \nu + iqe_d)}{\Gamma(2\nu)\Gamma(1 - \nu + iqe_d)} \cdot \frac{(m + i\tilde{m})R_2 + iqe_d - \nu}{(m + i\tilde{m})R_2 + iqe_d + \nu} (2\omega)^{2\nu}. \tag{5.19}$$

Now consider a charged fermion field on AdS$_2 \times \mathbb{R}^{d-1}$. Taking only $k_1 = k \neq 0$ dimensional reduction to equation (5.4), we find the action given by equation (5.10) with the identification $\tilde{m} = -(kr_*/R)(-1)^a$.

(c) Finite-temperature generalization

The above discussion can be generalized to finite temperature. In this case, the near-horizon region is a (charged) black hole in AdS$_2$ (times space),

$$ds^2 = \frac{R_2^2}{\xi^2} \left( -\left(1 - \frac{\xi^2}{\xi_0^2}\right) d\tau^2 + \frac{d\xi^2}{1 - \frac{\xi^2}{\xi_0^2}} \right) + \frac{\tau^2}{R^2} d\vec{x}^2, \quad A = e_d \frac{1 - \xi}{\xi_0} \left(1 - \frac{\xi}{\xi_0}\right) d\tau \tag{5.20}$$

and a temperature (with respect to $\tau$) $T = 1/2\pi\xi_0$. Locally, this black hole in AdS$_2$ is related by a coordinate transformation (combined with a gauge transformation) to the vacuum AdS$_2$, that is, it can be obtained by some global identifications. This fact can be used to infer the finite-temperature correlators from their zero-temperature limit presented above; they can also be determined directly as we do next.

The equation of motion for a minimally coupled scalar is

$$\frac{d^2}{d\xi^2} \phi + \frac{2\xi}{\xi^2 - \xi_0^2} \frac{d}{d\xi} \phi + \left( -\frac{m^2 R_2^2}{\xi^2 - \xi_0^2} + \frac{\omega + qe_d \left( \frac{1}{\xi} - \frac{1}{\xi_0} \right)^2}{1 - \frac{2\xi^2}{\xi_0^2} + \frac{\xi^2}{\xi_0^2}} \right) \phi = 0.$$
Two linearly independent solutions are given by

$$\phi(\zeta) \sim \left( \frac{1}{\zeta} - \frac{1}{\zeta_0} \right)^{-1/2+i\nu} \left( \frac{\zeta_0 + \zeta}{\zeta_0 - \zeta} \right)^{(i\omega_0/2) - iq_c} \times \mathbb{F}_1 \left( \frac{1}{2} \pm \nu + i\omega_0 - iq_c, 1/2 \pm \nu - i\omega, 1 \pm 2\nu; \frac{2\zeta}{\zeta - \zeta_0} \right),$$

where $\nu = \frac{1}{\sqrt{4 + m^2 R_2^2 - q_c^2}}$ and the top sign gives the ingoing solution.

Using $T = 1/2\pi \zeta_0$, the scalar retarded Green function has the following form:

$$G_R(\omega) = (4\pi T)^{2\nu} \frac{\Gamma(-2\nu) \Gamma \left( \frac{1}{2} + \nu - \frac{i\omega}{2\pi T} + iq_c \right) \Gamma \left( \frac{1}{2} + \nu - iq_c \right)}{\Gamma(2\nu) \Gamma \left( \frac{1}{2} - \nu - \frac{i\omega}{2\pi T} + iq_c \right) \Gamma \left( \frac{1}{2} - \nu - iq_c \right)}.$$

The spinor equation of motion at finite temperature is

$$\left( \partial_\gamma - i\tilde{\sigma}^3 \omega + qA_\gamma \right) \tilde{\phi} = \frac{R_2}{\zeta \sqrt{\tilde{f}}} (m\tilde{\sigma}^2 + \tilde{m}\tilde{\sigma}^1) \tilde{\phi},$$

(5.21)

where $\tilde{f} \equiv 1 - \zeta^2/\zeta_0^2$ is the emblacking factor in the AdS$_2$ black-hole metric in equation (5.20). It can similarly be solved in terms of hypergeometric functions. The general solution for the upper component $\tilde{y}$ of $\tilde{\phi}$ is

$$\tilde{y}(\zeta) = \left( \frac{\zeta + \zeta_0}{\zeta} \right)^{(1/2) + (i(qc + x_0\omega)/2)} \left[ c_{in} \left( -1 + \frac{\zeta_0}{\zeta} \right)^{-i\omega_0/2} \mathbb{F}_1 \left( \frac{1}{2} + iq_c - \nu - i\zeta_0 \omega, \frac{1}{2} + iq_c + \nu - i\zeta_0 \omega, \frac{3}{2} + i\zeta_0 \omega, \frac{\zeta - \zeta_0}{2\zeta} \right) + c_{out} \left( -1 + \frac{\zeta_0}{\zeta} \right)^{(i\omega_0/2) + (1/2)} \right].$$

(5.22)

The retarded function is then

$$G_R(\omega) = (4\pi T)^{2\nu} \frac{(m - i\tilde{m})R_2 + iq_c + \nu}{(m - i\tilde{m})R_2 + iq_c - \nu} \times \frac{\Gamma(-2\nu) \Gamma \left( \frac{1}{2} + \nu - \frac{i\omega}{2\pi T} + iq_c \right) \Gamma(1 + \nu - iq_c)}{\Gamma(2\nu) \Gamma \left( \frac{1}{2} - \nu - \frac{i\omega}{2\pi T} + iq_c \right) \Gamma(1 - \nu - iq_c)}.$$

(5.23)

Note that the branch point at $\omega = 0$ of zero temperature now disappears and the branch cut is replaced at finite temperature by a line of poles parallel to the next imaginary axis. In the zero-temperature limit, the pole line becomes a branch cut. Similar phenomena have been observed previously [62].

(d) Finite-temperature correlators and conformal invariance

In this subsection, we describe how conformal invariance of the IR CFT actually completely fixes the frequency dependence of the finite-temperature correlator derived above. The key point here, described in [63], is that the AdS$_2$ black-hole metric equation (5.20) is actually related to the $T = 0$ AdS$_2$ metric.
via a coordinate transformation that acts as a conformal transformation on the boundary of AdS$_2$. To see this explicitly, consider the following coordinate transformation to a new set of coordinates $(\sigma, t)$:

$$\tau \pm \xi_0 \tanh^{-1} \left( \frac{\xi}{\xi_0} \right) = \xi_0 \log \left( \frac{t \pm \sigma}{\xi_0} \right). \quad (5.24)$$

Working out the metric (5.20) in these new coordinates we find eventually

$$ds^2 = R_2^2 \left( \frac{dt^2}{\sigma^2} + \frac{d\sigma^2}{\sigma^2} \right). \quad (5.25)$$

This is just the $T = 0$ AdS$_2$ metric, as claimed above. To better understand what happened, consider the effect of the coordinate transformation (5.24) at the AdS$_2$ boundary $\sigma = \xi = 0$

$$t = \frac{1}{2\pi T} \exp(2\pi T \tau), \quad (5.26)$$

which is exactly the transformation that generates Rindler space in the $\tau$-coordinate from the vacuum in the $t$-coordinate. This is simply the statement that a coordinate choice which defines a black hole in the AdS$_2$ bulk is equivalent to a coordinate choice which puts the field theory at a finite temperature. Because of conformal invariance, this coordinate change is actually a symmetry operation in the IR CFT.

We note however that the gauge field adds a new subtlety; in particular, going through the same procedure with the zero-temperature gauge field (5.1) does not give us the finite-temperature gauge field (5.20). Let us go through the same steps as previously, starting this time with the zero-temperature configuration appropriate to the metric (5.25),

$$A = \frac{e_d}{\sigma} dt. \quad (5.27)$$

Writing this in terms of the finite-temperature coordinates, we obtain after some algebra

$$A = \frac{e_d}{\xi} d\tau + e_d d\left( \xi_0 \tanh^{-1} \left( \frac{\xi}{\xi_0} \right) \right). \quad (5.28)$$

Compare this to the gauge-field configuration in equation (5.20). It is not the same, but the difference is pure gauge; indeed the gauge field $\tilde{A}$ defined by

$$\tilde{A} = A + dA, \quad A = -2\pi T e_d \tau - e_d \left( \xi_0 \tanh^{-1} \left( \frac{\xi}{\xi_0} \right) \right)$$

is exactly the gauge field appearing in the finite-temperature solution (5.20). Thus, we have shown that the charged AdS$_2$ finite-temperature black hole is precisely the same as a coordinate transform plus a gauge transform of the vacuum AdS$_2$. Note that this gauge transform does not vanish at the AdS$_2$ boundary $\xi = 0$; the $d\tau$ part remains non-zero and corresponds to putting the IR CFT field theory at an (extra) constant value of $A_\tau$. 

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Now the transformation of the IR CFT correlators under each of these manipulations is known; thus, it should be possible to determine the finite-temperature correlators from the $T=0$ conformal results. For the conformal transformation (5.26), we use the usual CFT transformation law

$$\langle O^\dagger(t)O(t') \rangle = \left( \frac{dt}{d\tau} \right)^{\delta} \left( \frac{dt'}{d\tau'} \right)^{\delta} \langle O^\dagger(\tau)O(\tau') \rangle,$$

where $\delta = \frac{1}{2} + \nu$ is the dimension of the IR CFT operator.

The operation (5.29) is different and corresponds to turning on a source for the boundary gauge field that is pure gauge: generalizing slightly, in higher dimensions, it would correspond to $A_\mu = v_\mu L$, with $\mu$ running over all field-theory directions (in our case, only $t$). Thus, with the insertion of this source, we are now computing the field-theory correlator

$$\langle O^\dagger(x)O(x') \rangle_A = \exp \left( -i \int dy \, A(y) \partial_\mu j^\mu(y) \right) \langle O^\dagger(x)O(x') \rangle.$$ 

Putting these ingredients together and using the fact that in the $(t, \sigma)$-coordinate system, we have just

$$\langle O^\dagger(t)O(t') \rangle \sim (t - t')^{-2\delta},$$

we find

$$\langle O^\dagger(t)O(t') \rangle \sim \left( \frac{\pi T}{\sinh(\pi T(\tau - \tau'))} \right)^{2\delta} \exp(-2\pi i Te_d q(\tau - \tau')).$$

Here, the first factor is simply the usual expression for a finite-temperature correlator in the chiral half of a two-dimensional CFT in a thermal ensemble, where the role of the coordinate of the ‘chiral half’ is being played by the time coordinate. The extra factor is a new ingredient arising from the non-trivial sourcing of the gauge field. Upon Fourier transformation, we should find the usual expression from two-dimensional CFT at finite temperature (e.g. [55,64]), except with a shift in the frequency from the gauge-field contribution,

$$G(\omega) \sim (2\pi T)^{2\delta-1} \frac{\Gamma(\delta - (i/2\pi T)(\omega - 2\pi qe_d T))}{\Gamma(1 - \delta - (i/2\pi T)(\omega - 2\pi qe_d T))}.$$ 

This frequency dependence agrees with that of the correlator derived directly from the bulk wave equation above.

6. Finite temperature

At small non-zero temperature, the near-horizon geometry is a black hole in AdS2. Using the results for $G(\omega, T)$ from §5c for Green’s functions resulting from this
IR geometry, the fermion self-energy becomes

$$\Sigma(\omega, T) = T^{2p} g\left(\frac{\omega}{T}\right) = (4\pi T)^{2p} \frac{\Gamma\left(\frac{1}{2} + \nu - \left(i\omega/2\pi T\right) + i q e_d\right)}{\Gamma\left(\frac{1}{2} - \nu - \left(i\omega/2\pi T\right) + i q e_d\right)} \xrightarrow{T \to 0} c_k \omega^{2r}. \quad (6.1)$$

We first note from this formula (6.1) that what at $T = 0$ was a branch cut for $\omega^{2p}$ has become a line of discrete poles of the Gamma function.

An interesting phenomenon which occurs at $T = 0$ is that as $k \to k_F$, the quasiparticle pole sometimes moves under the branch cut and escapes onto another sheet of the complex $\omega$-plane; this leads to an extreme particle–hole asymmetry (visible in fig. 2c of [28]). More precisely, when the pole hits the branch point of $G(\omega)$ at $\omega = 0$, it changes its ‘velocity’ $d\omega/dk$ according to $v_F(|k| - k_F) = G(\omega)$. Depending on the phase of $G$, this can put the pole on another sheet of the complex $\omega$-plane.

This raises an interesting question: what happens to the pole that at $T = 0$ would have gone under the branch cut? The answer is that it joins the line of poles approximating the branch cut; their combined motion mimics the effects of the motion of the pole on the second sheet. This effect is visible in the sequence of pictures in figure 10.

At finite $T$, the pole no longer hits the real axis; the distance of closest approach is of order $(\pi T)^{2r}$. This is the thermal smearing of the Fermi surface.
7. Discussion: charged anti de Sitter space black holes and frustration

The state we have been studying has a large low-lying density of states. The entropy density of the black hole is

$$s(T = 0) = \frac{1}{V_{d-1}} \frac{A}{4G_N} = 2\pi e_d \rho.$$  \hspace{1cm} (7.1)

At leading order in \(1/N^2\), this is a large ground-state degeneracy. The state we are studying is not supersymmetric, and so we expect this degeneracy to be lifted at finite \(N\). The most probable way in which the third law of thermodynamics will be enforced is by an instability to a superconducting state, the implementation and effects of which are discussed below in §9. In the absence of the necessary ingredients for a holographic superconductor, one way in which the degeneracy can be lifted was pointed out in [65]. The ground state of the spinor field in the RN black hole is a Fermi sea of filled negative-energy states (a related discussion for a different black hole appears in [59]). An arbitrarily small density of matter close enough to the extremal horizon will produce an order-one back reaction on the bulk geometry. According to [65], the density of spinor particles themselves modify the far-IR of the AdS\(_2 \times \mathbb{R}^2\) to a Lifshitz geometry [66] with dynamical exponent \(z \propto N^2\). This gravitating object, supported by the degeneracy pressure of charged fermions, has been called an ‘electron star’ [67]. The modification occurs out to a very small distance from the horizon that scales like \(e^{N^2}\), and does not change the features of our calculations that we have emphasized above. (More concretely, the results are unchanged down to temperatures scaling like \(e^{-N^2}\).) If one considers instead spinor fields whose mass and charge grow with \(N\) (e.g. like \(q \sim N\)), then the back reaction of a finite density can modify the geometry out to a larger radius of order \(1/N\) to a Lifshitz background with a dynamical exponent that is finite in the large-\(N\) limit [67].

We should comment in more detail on the effects on \(G_R\) of changes to the near-horizon geometry resulting from the gravitational back reaction of the bulk fermion density [65,68]. The geometry AdS\(_2 \times \mathbb{R}^2\) discussed above is the \(z \to \infty\) limit of the following family of metrics with Lifshitz scaling \(t \to \lambda^z t, x \to \lambda x\):

$$ds^2 = \frac{R^2}{\xi^2} (d\tau^2 + d\xi^2) + \frac{\tau^2}{R^2} \xi^{2/z} dx^2.$$ \hspace{1cm} (7.2)

Any finite \(z\) modifies the non-analytic behaviour of the self-energy [68]: the Lifshitz scaling implies that the IR CFT scaling function takes the form \(G_R(\omega, k) = \omega^\alpha F(\omega/k^z)\); note that it is \(k\) and not \(k_\perp\) that appears in the scaling function—the IR CFT knows nothing of \(k_F\). Lifshitz wave equations for general \(z\) seem not to be exactly solvable, but at small frequency, a Wentzel–Kramers–Brillouin analysis can be used to study the scaling function more explicitly [68]; the self-energy goes like \(\exp(-1/\omega^{1/(z-1)})\) for frequencies less than the scale at which the RN ground-state degeneracy is split. We emphasize that, nevertheless, the non-Fermi-liquid behaviour persists down to parametrically low energies (of order \(\mu e^{-z}\)). There exist other geometries in which the ground entropy vanishes. For example, the inclusion of other light bulk modes (e.g. neutral scalars) [69–71] can have an important effect on the ground state. A systematic exploration of the fermion response in the general holographic description of finite density will be valuable.
In this brief subsection, we observe that some features of the electron star are already visible in the fermion correlators computed in the RN black hole at leading order in $1/N$.

The electronic states in the curved background correspond to the solutions of the bulk Dirac equation with normalizable boundary conditions in the UV and regular boundary conditions (ingoing in a Lorentzian signature) in the IR. The transverse momenta can be chosen continuously, $k_x, k_y \in \mathbb{R}$. In the radial $r$ direction, there is no translational invariance, so in this direction, we describe the wave function in real space. The radial quantum number $n$ is discrete because boundary conditions have been specified. Finally, after fixing the transverse momenta, the frequency should be adjusted such that the solution satisfies the boundary conditions. This happens for $\omega$ values where the retarded Green function has a pole.

Instead of labelling the states by $\{k_x, k_y, n\}$, where $n$ is the discrete radial quantum number, we can label them by $\{k_x, k_y, \omega\}$, where $\omega$ is chosen from the discrete set of poles of the retarded correlator with fixed $k_x$ and $k_y$. The states will be automatically filled for $\text{Re} \, \omega < 0$ and empty for $\text{Re} \, \omega > 0$.

If there were translational invariance in the radial direction, then $n$ would be a continuous parameter. When labelling the states by $\omega$, this means that there is a continuum of poles in the correlator from which we can choose. (In practice, the poles would probably form a branch cut.) In the bulk, deep inside the electron star, there is locally a three-dimensional Fermi surface. Since AdS/CFT should describe the excitations of the star, we expect infinitely many poles in Green’s functions for some values of the momenta.

Surprisingly, something like this already happens in the oscillatory region ($k < k_o$) in the RN case. This is a regime of momenta where Green’s functions are observed to have log-periodic behaviour in frequency, and where the zero-frequency spectral weight is non-vanishing [30]. The basic observation is that although the poles do not yet form a branch cut, they do accumulate near $\omega = 0$. In tortoise coordinates, one can see that, exactly in this oscillatory region, the wave function ‘falls into the black hole’, i.e. it is localized in the IR.

We note further that in examples without an oscillatory region (as is the case for the alternative quantization, see the right part of fig. 5 of [26]), there is no bulk Fermi sea near the horizon, and therefore no such modification of the geometry. We must observe, however, that in known examples that have a Fermi surface but no oscillatory region, there is a relevant operator, perturbation by which removes the Fermi surface. It would be interesting to find a stable example without an oscillatory region.

8. Transport

The most prominent mystery of the strange-metal phase is the linear-in-$T$ electrical resistivity. Electron–electron scattering (combined with umklapp or impurities) produces $\rho \sim T^2$, electron–phonon scattering produces $\rho \sim T^5$. A van Hove singularity requires fine tuning of the Fermi level. No simple, robust effective field theory gives $\rho \sim T$.

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AdS/CFT techniques are well adapted for studying transport. Holographic conductivity is computed by solving Maxwell’s equations in the bulk. The answer is of the form

$$\sigma_{DC}^{\text{DC}} = \lim_{\omega \to 0} \text{Im} \frac{1}{\omega} (j^x j^x)(\omega, 0) = O(N^2) + N^0 \sigma_{\text{from spinor}} + \cdots . \quad (8.1)$$

The $O(N^2)$ contribution does not know about $k_F$ as it only depends on the black-hole geometry, and not on the dynamics of the fermions. The contribution to the conductivity from the holographic Fermi surface is down by $N^{-2}$: its extraction requires a (spinor) loop in the bulk, as in [72–74].

The result is very similar to the calculation of Fermi-liquid conductivity (e.g. [75]), with extra integrals over $r$, and no vertex corrections. The only process that contributes to leading order is depicted in figure 11. The key step is to relate the bulk spinor spectral function to that of the boundary fermion operator,

$$\text{Im} S_{\alpha\beta}(Q, k; r_1, r_2) = \frac{\psi^b_\alpha(Q, k, r_1) \tilde{\psi}^\beta_\beta(Q, k, r_2)}{W_{ab}} A(Q, k), \quad (8.2)$$

where $A = (1/\pi)\text{Im} G_R$ is the spectral weight computed above. The conductivity is

$$\sigma_{DC}^{\text{FS}} = S_{d} \int dk \frac{k^{d-2}}{k^{d-2}} \int d\omega \frac{df}{d\omega} A^2(k, \omega) A(\omega, k)^2, \quad (8.3)$$

where $f$ is the Fermi function and $A \sim q \int_0^\infty dr \sqrt{g} g^{xx} A_x(r, 0)(\tilde{\psi}^b(r) \Gamma^x \psi^h(r))/W_{ab}$ is a vertex factor, encoding data analogous to the UV coefficients $v_F, h_{1,2}$ in equation (4.8),

$$\sigma_{\text{from FS}}^{\text{DC}} = Y^2 k_F^{d-2} \int d\omega \frac{f'(\omega)}{\text{Im} G(\omega)} \sim T^{-2\nu}. \quad (8.4)$$

The factor $Y$ is determined by bound-state wavefunctions and $k_F^{d-2}$ is the volume of the Fermi surface. In the last step, we have used the scaling relation $\text{Im} G(T \to 0, \omega/ T \text{fixed}) = T^{2\nu} g(\omega/ T)$.

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In the marginal Fermi-liquid case where \( n = \frac{1}{2} \), this indeed gives
\[
\rho = (\sigma^{\text{DC}})^{-1} \sim T. \tag{8.5}
\]

In contrast to the models of non-Fermi liquids that arise by coupling a Fermi surface to a gapless bosonic mode, the transport lifetime and the single-particle lifetime have the same temperature dependence. This is possible because the quasi-particle decay is mediated by the IR CFT, which contains low-energy modes for non-zero (indeed, for any) momentum.

Although the direct current (DC) resistivity is not sensitive to the stability of the quasi-particles, the behaviour of the optical conductivity does change dramatically at \( n = \frac{1}{2} \) [28,29].

9. The superconducting state

A useful test of any model of the normal state is whether it can incorporate the transition to superconductivity. It is indeed possible to describe superconductivity holographically by including charged scalar fields in the black-hole background. At low temperature, they can condense, spontaneously breaking the \( U(1) \) symmetry, changing the background [36–38].

The problem of the fermion response in various possible holographic superconducting phases has been studied [24,76–81]. The superconducting condensate can open a gap in the fermion spectrum around the chemical potential [27] if a suitable bulk coupling between spinor and scalar is included. The bulk action we consider for the fermion is
\[
S[\psi] = \int d^{d+1}x \sqrt{-g} \left[ i \dot{\psi} (i_M D_M - m_\psi) \psi + \eta^*_5 \phi^* \psi^T C \Gamma^5 \psi + \eta_5 \phi \bar{\psi} C \Gamma^5 \bar{\psi}^T \right], \tag{9.1}
\]
where \( \phi \) is the scalar field whose condensation spontaneously breaks the \( U(1) \) symmetry, \( C \) is the charge conjugation matrix and \( \Gamma^5 \) is the chirality matrix, \( \{ \Gamma^5, \Gamma^M \} = 0 \). The superconducting order parameter in these studies is s-wave, but is not BCS.

Interestingly, in the condensed phase, one finds stable quasi-particles, even when \( n \leq \frac{1}{2} \). The modes into which the quasi-particle decays in the normal state are lifted by the superconducting condensate. The quasi-particles are stable in a certain kinematical regime, similar to Landau’s critical velocity for drag in a superfluid: these holographic superconductor groundstates do have gapless excitations (in fact a relativistic CFT worth), but the Fermi surface can occur outside their lightcone (visible as the dashed line in figure 12).

10. Conclusions

The holographic calculation of the spectral function can be described by a simple low-energy effective theory [26,28,68,83]. Consider a Fermi liquid (with creation
Figure 12. The effect of the superconducting order on the fermion spectral density (from Faulkner et al. [27]). Shown are plots of $A(k, \omega)$ at various $k \in [0.81, 0.93]$ for $q_\psi = \frac{1}{2}, m_\psi = 0$ in a low-temperature background of a scalar with $q_\psi = 1, m_{\phi}^2 = -1$ (first constructed in [82]), with $\eta_5 = 0.025$. The dashed line indicates the boundary of the region in which the incoherent part of the spectral density is completely suppressed, and the lifetime of the quasi-particle is infinite. The dashed line with filled circles indicates the location of the peak. (Online version in colour.)

operator $\psi$) mixing with a bath of critical fermionic fluctuations with a large dynamical exponent,

$$L = \bar{\psi}(\omega - v_F k)\psi + \bar{\psi}\chi + \psi\bar{\chi} + \bar{\chi}G^{-1}\chi,$$  \hspace{1cm} (10.1)

where $\chi$ is an operator in the IR CFT; the $k$-independence of its correlations suggests that it arises from localized critical degrees of freedom. The corrected $\psi$ Green function is given by the geometric series

$$\langle \bar{\psi}\psi \rangle = \frac{1}{\omega - v_F k - G}G = \langle \bar{\chi}\chi \rangle = c(k)\omega^{2

Note that according to the free fermion scaling, $\chi$ has dimension zero; therefore for $\nu \leq \frac{1}{2}$, the $\bar{\psi}\chi$ coupling is a relevant perturbation.

Above, we have described a certain class of fixed points that have some features in common with various non-Fermi-liquid metals. We find Fermi surfaces with vanishing quasi-particle residue. The single-fermion self-energy is a power law in frequency, independent of momentum, as in dynamical mean-field theory and some slave-particle descriptions. Whether the states we describe can arise from any model of electrons with short-range interactions is an important open question.7

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7See [84,85] for some recent ideas in this direction.

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References


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