Strange metals, black holes, and graphene

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Foundations of quantum many body theory:

1. Ground states connected adiabatically to independent electron states
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2. Boltzmann-Landau theory of quasiparticles
Entangled phases of quantum matter:

1. Ground states disconnected from independent electron states: many-particle entanglement

2. Boltzmann-Landau theory of quasiparticles

Topological quantum matter:

The fractional quantum Hall effect of electrons in two dimensions (e.g. in graphene) in the presence of a strong magnetic field. The ground state is described by Laughlin’s wavefunction, and the excitations are quasiparticles which carry fractional charge.
Entangled phases of quantum matter:

1. Ground states disconnected from independent electron states: many-particle entanglement

2. No quasiparticles

Strange metals with no quasiparticles:

Such metals are found, most prominently, in certain high temperature superconductors.
Entangled phases of quantum matter:

1. Ground states disconnected from independent electron states: many-particle entanglement

2. No quasiparticles

Strange metals with no quasiparticles:

Such metals are found, most prominently, in certain high temperature superconductors.

But how can we be sure that no quasiparticles exist in a given system? Perhaps there are some entangled quasiparticles inaccessible to current experiments……..
Local thermal equilibration or phase coherence time, $\tau_\varphi$:

- There is an *lower bound* on $\tau_\varphi$ in all many-body quantum systems of order $\hbar/(k_B T)$,

$$\tau_\varphi > C \frac{\hbar}{k_B T},$$

and the lower bound is realized by systems *without* quasiparticles.

- In systems *with* quasiparticles, $\tau_\varphi$ is parametrically larger at low $T$;
  e.g. in Fermi liquids $\tau_\varphi \sim 1/T^2$,
  and in gapped insulators $\tau_\varphi \sim e^{\Delta/(k_B T)}$ where $\Delta$ is the energy gap.
A bound on quantum chaos:

- The time over which a many-body quantum system becomes “chaotic” is given by $\tau_L = 1/\lambda_L$, where $\lambda_L$ is the “Lyapunov exponent” determining memory of initial conditions. This Lyapunov time obeys the rigorous lower bound

$$\tau_L \geq \frac{1}{2\pi} \frac{\hbar}{k_B T}$$

A. I. Larkin and Y. N. Ovchinnikov, JETP 28, 6 (1969)

A bound on quantum chaos:

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\[
\tau_L \geq \frac{1}{2\pi} \frac{\hbar}{k_B T}
\]

Quantum matter without quasiparticles
\( \approx \) fastest possible many-body quantum chaos
$\text{BaFe}_2(\text{As}_{1-x}\text{P}_x)_2$

A high temperature superconductor
Superconductivity

\[ \text{Resistivity} \sim \rho_0 + AT^\alpha \]


Resistivity \( \sim \rho_0 + A T^\alpha \)


Strange metals

\[
\frac{1}{\tau} = \alpha \frac{k_B T}{\hbar}
\]

LIGO
September 14, 2015
• Black holes have a “ring-down” time, \( \tau_r \), in which they radiate energy, and stabilize to a ‘featureless’ spherical object. This time can be computed in Einstein’s general relativity theory.
• ‘Featureless’ black holes have a Bekenstein-Hawking entropy, and a Hawking temperature, $T_H$. 
Expressed in terms of the Hawking temperature, the ring-down time is \( \tau_r \sim \frac{\hbar}{k_B T_H} \)!
Expressed in terms of the Hawking temperature, the ring-down time is $\tau_r \sim \frac{\hbar}{k_B T_H}$.

For this black hole $T_H \approx 1$ nK, $\tau_r = 7.7$ milliseconds. (Radius of black hole = 183 km; Mass of black hole = 62 solar masses.)
• Is there a connection between strange metals and black holes?

• Why do they have the same equilibration time $\sim \hbar/(k_B T)$?
The Sachdev-Ye-Kitaev (SYK) model:

- A theory of a strange metal
- Has a dual representation as a black hole
Infinite-range (SYK) model of a strange metal

\[ H = \frac{1}{(2N)^{3/2}} \sum_{i,j,k,\ell=1}^{N} J_{ij;kl} c_i^\dagger c_j^\dagger c_k c_\ell - \mu \sum_i c_i^\dagger c_i \]

\[ c_i c_j + c_j c_i = 0 \quad , \quad c_i c_j^\dagger + c_j^\dagger c_i = \delta_{ij} \]

\[ Q = \frac{1}{N} \sum_i c_i^\dagger c_i \]

\( J_{ij;kl} \) are independent random variables with \( \bar{J}_{ij;kl} = 0 \) and \( |\bar{J}_{ij;kl}|^2 = J^2 \)

\( N \to \infty \) yields critical strange metal.

S. Sachdev and J. Ye, PRL 70, 3339 (1993)

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\[ Q = \frac{1}{N} \sum_i c_i^\dagger c_i \]

A fermion can move only by entangling with another fermion: the Hamiltonian has “nothing but entanglement”.

S. Sachdev and J. Ye, PRL 70, 3339 (1993)
$$H = \frac{1}{(2N)^{3/2}} \sum_{i,j,k,\ell=1}^{N} J_{i;j,k,\ell} c_i^\dagger c_j^\dagger c_k c_\ell$$
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\[ Q = \frac{1}{N} \sum_{i} \langle c_{i}^\dagger c_{i} \rangle. \]

Local fermion density of states

\[ \rho(\omega) \sim \begin{cases} \omega^{-1/2}, & \omega > 0 \\ e^{-2\pi\epsilon |\omega|^{-1/2}}, & \omega < 0. \end{cases} \]

and “conformal” extension to \( T > 0. \)

Known ‘equation of state’ determines \( \mathcal{E} \) as a function of \( Q \)

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S. Sachdev and J. Ye, PRL 70, 3339 (1993)

A. Georges and O. Parcollet PRB 59, 5341 (1999)

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Known ‘equation of state’ determines \( \mathcal{E} \) as a function of \( Q \)

Microscopic zero temperature entropy density, \( S \), obeys
\[ \frac{\partial S}{\partial Q} = 2\pi \mathcal{E} \]

A. Georges, O. Parcollet, and S. Sachdev
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\end{cases}$

and “conformal” extension to $T > 0$.

Known ‘equation of state’ determines $\mathcal{E}$ as a function of $Q$

The SYK strange metal is holographically dual to the gravity theory of the AdS$_2$ near-horizon geometry of charged black holes

S. Sachdev,

A. Georges, O. Parcollet, and S. Sachdev
PRB 63, 134406 (2001)
\[ H = \frac{1}{(2N)^{3/2}} \sum_{i,j,k,\ell=1}^{N} J_{ij;kl} c_i^\dagger c_j^\dagger c_k c_\ell \]

\[ Q = \frac{1}{N} \sum_i \langle c_i^\dagger c_i \rangle. \]

Local fermion density of states
\[ \rho(\omega) \sim \begin{cases} \omega^{-1/2}, & \omega > 0 \\ e^{-2\pi E} |\omega|^{-1/2}, & \omega < 0. \end{cases} \]

and "conformal" extension to \( T > 0 \).

Known ‘equation of state’ determines \( E \) as a function of \( Q \)

Microscopic zero temperature entropy density, \( S \), obeys
\[ \frac{\partial S}{\partial Q} = 2\pi E \]

Einstein-Maxwell theory + cosmological constant

Horizon area \( A_h; \)
AdS\(_2 \times R^d\)
\[ ds^2 = (d\zeta^2 - dt^2)/\zeta^2 + d\vec{x}^2 \]
Gauge field: \( A = (E/\zeta)dt \)

Boundary area \( A_b; \)
charge density \( Q \)

\[ \zeta = \infty \quad \zeta \]

\[ L = \bar{\psi} \Gamma^\alpha D_\alpha \psi + m \bar{\psi} \psi \]
Local fermion density of states
\[ \rho(\omega) \sim \begin{cases} \omega^{-1/2}, & \omega > 0 \\ e^{-2\pi E} |\omega|^{-1/2}, & \omega < 0. \end{cases} \]

and identical conformal extension to \( T > 0 \).
‘Equation of state’ relating \( E \)
and \( Q \) depends upon the geometry of spacetime far from the AdS\(_2\)
Known ‘equation of state’ determines \( E \) as a function of \( Q \). Einstein-Maxwell theory + cosmological constant

\[ H = \frac{1}{(2N)^{3/2}} \sum_{i,j,k,\ell=1}^{N} J_{ij; k\ell} c_i^{\dagger} c_j^{\dagger} c_k c_{\ell} \]

\[ Q = \frac{1}{N} \sum_i \langle c_i^{\dagger} c_i \rangle. \]

Local fermion density of states

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and “conformal” extension to \( T > 0 \).

Known ‘equation of state’ determines \( E \) as a function of \( Q \).

Microscopic zero temperature entropy density, \( S \), obeys

\[ \frac{\partial S}{\partial Q} = 2\pi\epsilon \]

\( \zeta = \infty \)

\( \zeta \)

Black hole thermodynamics (classical general relativity) yields

\[ \frac{\partial S_{\text{BH}}}{\partial Q} = 2\pi\epsilon \]

A. Sen hep-th/0506177; S. Sachdev PRX 5, 041025 (2015)
\[ H = \frac{1}{(2N)^{3/2}} \sum_{i,j,k,\ell=1}^{N} J_{i;j;\ell} \, c_{i}^{\dagger} c_{j}^{\dagger} c_{k} c_{\ell} \]

- The two models have the same symmetry at low energies: reparameterization invariance spontaneously broken to \( \text{SL}(2, \mathbb{R}) \). At low energies they are described by the same effective action (the “Schwarzian”) for a scalar field which is the remnant on the graviton on \( \text{AdS}_{2} \).

Einstein-Maxwell theory + cosmological constant

Boundary area \( \mathcal{A}_b \);

Charge density \( Q \)

Horizon area \( \mathcal{A}_h \);

\[ ds^2 = (d\zeta^2 - dt^2)/\zeta^2 + d\vec{x}^2 \]

Gauge field: \( A = (E/\zeta)dt \)

$H = \frac{1}{(2N)^{3/2}} \sum_{i,j,k,\ell=1}^{N} J_{ij; k\ell} \, \bar{c}_{i}^{\dagger} \bar{c}_{j}^{\dagger} \bar{c}_{k} \bar{c}_{\ell}$

Einstein-Maxwell theory + cosmological constant

Horizon area $A_{h}$; 
AdS$_{2} \times R^{d}$
$ds^{2} = (d\zeta^{2} - dt^{2})/\zeta^{2} + d\vec{x}^{2}$
Gauge field: $A = (E/\zeta)dt$

Boundary area $A_{b}$; charge density $Q$

The scrambling times of the SYK model and of black holes in Einstein gravity saturate the bound on quantum chaos

$\tau_{L} = \frac{1}{2\pi} \frac{\hbar}{k_{B}T}$

• Is there a connection between strange metals and black holes?

• Why do they have the same equilibration time $\sim \hbar/(k_B T)$?
• Is there a connection between strange metals and black holes? Yes, *e.g.* the SYK model.

• Why do they have the same equilibration time $\sim \hbar/(k_B T)$? Strange metals don’t have quasiparticles and thermalize rapidly; Black holes are “fast scramblers”.
Graphene

Strange metal transport
Theoretical predictions inspired by holography
Comparison with experiments
Graphene

Same “Hubbard” model as for ultracold atoms, but for electrons on the honeycomb lattice
Graphene

Electron Fermi surface

$\mu > 0$
Graphene

Hole Fermi surface

Electron Fermi surface

$\mu < 0$

$\mu > 0$
Graphene

\[ T(K) \]

Quantum critical
Dirac liquid

Hole
Fermi liquid

Electron
Fermi liquid

M. Müller, L. Fritz, and S. Sachdev, PRB 78, 115406 (2008)
M. Müller and S. Sachdev, PRB 78, 115419 (2008)
Graphene

\[
\sim 1 \sqrt{n} (1 + \lambda \ln \Lambda \sqrt{n})
\]

\[T(K)\]

Quantum critical Dirac liquid

Hole Fermi liquid

Electron Fermi liquid

Predicted "strange metal" without quasiparticles

M. Müller, L. Fritz, and S. Sachdev, PRB 78, 115406 (2008)
M. Müller and S. Sachdev, PRB 78, 115419 (2008)
Fermi liquids: quasiparticles moving ballistically between impurity (red circles) scattering events
Fermi liquids: quasiparticles moving ballistically between impurity (red circles) scattering events

Strange metals: electrons scatter frequently off each other, so there is no regime of ballistic quasiparticle motion. The electron “liquid” then “flows” around impurities
Thermal and electrical conductivity with quasiparticles

- Wiedemann-Franz law in a Fermi liquid:

\[ L_0 = \frac{\kappa}{\sigma T} \approx \frac{\pi^2 k^2_B}{3e^2} \approx 2.45 \times 10^{-8} \frac{W \cdot \Omega}{K^2}. \]

Transport in Strange Metals

For a strange metal with a “relativistic” Hamiltonian, hydrodynamic, holographic, and memory function methods yield Lorentz ratio

\[ L = \frac{\kappa}{(T\sigma)} = \frac{v_F^2 \mathcal{H} \tau_{\text{imp}}}{T^2 \sigma Q} \left(1 + e^2 v_F^2 Q^2 \tau_{\text{imp}}/(\mathcal{H} \sigma Q)^2\right)^2 \]

\( Q \rightarrow \) electron density; \( \mathcal{H} \rightarrow \) enthalpy density
\( \sigma_Q \rightarrow \) quantum critical conductivity
\( \tau_{\text{imp}} \rightarrow \) momentum relaxation time from impurities.

Note that for a clean system (\( \tau_{\text{imp}} \rightarrow \infty \) first), the Lorentz ratio diverges \( L \sim 1/Q^4 \), as we approach “zero” electron density at the Dirac point.
Graphene

Quantum critical
Dirac liquid

Hole
Fermi liquid

Electron
Fermi liquid

$T(K)$

$\sim \sqrt{n} (1 + \lambda \ln \Lambda \sqrt{n})$

$T(K)$

$10^{12} / m^2$

Predicted strange metal

M. Müller, L. Fritz, and S. Sachdev, PRB 78, 115406 (2008)
M. Müller and S. Sachdev, PRB 78, 115419 (2008)
Graphene

Quantum critical
Disordered

$T(K)$

$\sim \sqrt{n} (1 + \lambda \ln \Lambda \sqrt{n})$

$T(K)$

Hole
Fermi liquid

Electron
Fermi liquid

$\mu > 0$

$\mu < 0$

$n \sim 10^{12} / m^2$

Predicted strange metal

M. Müller, L. Fritz, and S. Sachdev, PRB 78, 115406 (2008)

M. Müller and S. Sachdev, PRB 78, 115419 (2008)
Graphene

Predicted strange metal

Quantum critical

Electron Fermi liquid

Hole Fermi liquid

Impurity scattering dominates

$T(K)$

$\sqrt{n}$

$n / 10^{12} / m^2$

M. Müller, L. Fritz, and S. Sachdev, PRB 78, 115406 (2008)
M. Müller and S. Sachdev, PRB 78, 115419 (2008)
Red dots: data
Blue line: value for $L = L_0$
To minimize disorder, the monolayer graphene samples used in this report are encapsulated in hexagonal boron nitride (hBN). All devices used in this study are representative device (see SM for all samples). From this, all measurements are performed in a cryostat controlling bath temperatures. The residual carrier density at the neutrality point (green) is estimated by the intersection of the minimum conductivity [\(\kappa_{\text{min}}\)] as a function of log(\(\Delta V_g\)).

Red dots: data
Blue line: value for \(L = L_0\)

**FIG. 1.** Temperature and density dependent electrical and thermal conductivity. (A) Thermal conductivity (red points) as a function of (C) gate voltage and (D) bath temperature. (B) Electrical conductivity (blue) as a function of the charge density set by the back gate for different temperatures. At low temperature and/or high doping (Fig. 1B), electron-phonon coupling becomes significant, leading to disorder and an increase in thermal excitations. This results in the sample entering the non-degenerate regime near the charge neutrality point. In this regime, the thermal energy becomes larger than the local chemical potential, and the Wiedemann-Franz law (WF law) is violated. This is a non-trivial check on the quality of our measurement. In the non-degenerate regime (\(\Delta V_g > 0\)), the neutrality point is approached, and thermal excitations become significant at low temperature. The residual carrier density at the neutrality point is estimated by the intersection of the minimum conductivity (green) with the temperature axis (Fig. 1D). Solid black lines correspond to 4 times the mean free path in the half-tone background, representing the thermal conductivity. Formation of the DF is further complicated by phonon scattering at high temperature, which can relax momentum by creating additional inelastic scattering events. This high temperature limit occurs when the thermal energy is larger than the local chemical potential, and even when the sample is globally neutral, it will exhibit finite potential. All data from this figure is taken from the known sample dimensions. At the CNP, the residual carrier density can be estimated by extrapolating a linear fit of log(\(\kappa_{\text{min}}\)) versus \(\Delta V_g\).
two-terminal to keep a well-defined temperature profile.

The breakdown of the WF law can be observed.

Temperatures set the experimental window in which the DF and the electron-electron scattering rate. These two temperatures cause spatial variations in the local chemical potential.

The thermal energy be larger than the local chemical potential.

To minimize disorder, the monolayer graphene samples used in this report are encapsulated in hexagonal boron nitride (hBN).

Formation of the DF is further complicated by phonon scattering at high temperature which can relax momentum by creating additional inelastic scattering.

The thermal conductivity is enhanced and the WF law is violated. Above this high temperature limit occurs when the thermal energy be larger than the local chemical potential.

The neutrality point. In the non-degenerate regime compared to the Wiedemann-Franz law, the minimum density is limited by disorder (charge puddles). However, above the minimum density (green) aligns with the temperature axis to the right. Solid black lines correspond to 4 and 100 K.

The residual carrier density at the neutrality point (green) is estimated by the intersection of the minimum conductivity with a linear fit to log(σ) versus T.

At low temperature and/or high doping (∆Vg), electron-phonon coupling becomes appreciable and begins to dominate thermal transport at all measured gate voltages. All data from this figure is taken from the silicon substrate to tune the charge carrier density.

FIG. 1. Temperature and density dependent electrical and thermal conductivity. (A) Electrical conductivity (blue) as a function of the charge density set by the back gate for different temperatures. Fig. 1B shows the resistance (C) gate voltage and (D) bath temperature.

Red dots: data
Blue line: value for \( L = L_0 \)
Graphene

Quantum critical
Disorder

Hole
Fermi liquid

Electron
Fermi liquid

Impurity scattering dominates

M. Müller, L. Fritz, and S. Sachdev, PRB 78, 115406 (2008)
M. Müller and S. Sachdev, PRB 78, 115419 (2008)
Strange metal in graphene

J. Crossno et al., Science 351, 1058 (2016)

Wiedemann-Franz obeyed
Strange metal in graphene

J. Crossno et al., Science 351, 1058 (2016)

Wiedemann-Franz law violated!
Lorentz ratio $L = \frac{\kappa}{(T\sigma)}$

$$= \frac{v_F^2 H \tau_{\text{imp}}}{T^2 \sigma_Q} \frac{1}{(1 + e^2 v_F^2 Q^2 \tau_{\text{imp}}/(H \sigma_Q))^2}$$

$Q \rightarrow$ electron density; $H \rightarrow$ enthalpy density

$\sigma_Q \rightarrow$ quantum critical conductivity

$\tau_{\text{imp}} \rightarrow$ momentum relaxation time from impurities


J. Crossno et al., Science 351, 1058 (2016)
Comparison to theory with a single momentum relaxation time $\tau_{\text{imp}}$. Best fit of density dependence to thermal conductivity does not capture the density dependence of electrical conductivity.
Figure 1: A comparison of our hydrodynamic theory of transport with the experimental results of \[33\] in clean samples of graphene at \(T = 75\) K. We study the electrical and thermal conductances at various charge densities \(n\) near the charge neutrality point. Experimental data is shown as circular red data markers, and numerical results of our theory, averaged over 30 disorder realizations, are shown as the solid blue line. Our theory assumes the equations of state described in \((27)\) with the parameters \(C_0 \equiv 11\), \(C_2 \equiv 9\), \(C_4 \equiv 200\), \(\rho_0 \equiv 110\), and \(h_0 \equiv 1.7\). The yellow shaded region shows where Fermi liquid behavior is observed and the Wiedemann-Franz law is restored, and our hydrodynamic theory is not valid in or near this regime. We also show the predictions of \((2)\) as dashed purple lines, and have chosen the 3 parameter fit to be optimized for \(L(n)\).


where \(e\) is the electron charge, \(s\) is the entropy density, \(n\) is the charge density (in units of length \(2\)), \(H\) is the enthalpy density, \(\tau\) is a momentum relaxation time, and \(q\) is a quantum critical effect, whose existence is a new effect in the hydrodynamic gradient expansion of a relativistic fluid. Note that up to \(q\), \(n^2\) is simply described by Drude physics. The Lorenz ratio then takes the general form

\[
L(n) = L_{DF}(1 + (n/n_0)^2)^2,
\]

\[(3)\]

Where

\[
L_{DF} = v_F^2 H/\tau T^2 q,
\]

\[(4a)\]

\[
n_0^2 = H q e^2 v_F^2 H/\tau.
\]

\[(4b)\]

\(L(n)\) can be parametrically larger than \(L_{WF}\) (as \(\tau \to 1\) and \(n \tau n_0\)), and much smaller \((n/n_0)\).

Both of these predictions were observed in the recent experiment, and fits of the measured \(L\) to \((3)\) were quantitatively consistent, until large enough \(n\) where Fermi liquid behavior was restored. However, the experiment also found that the conductivity did not grow rapidly away from \(n = 0\) as predicted in \((2)\), despite a large peak in \(\eta\) near \(n = 0\), as we show in Figure 1. Furthermore, the theory of \[25\] does not make clear predictions for the temperature dependence of \(\tau\), which determines \(\eta\).

In this paper, we argue that there are two related reasons for the breakdown of \((2)\). One is that the dominant source of disorder in graphene – fluctuations in the local charge density, commonly referred to as charge puddles \([43, 44, 45, 46]\) – are not perturbatively weak, and therefore a non-perturbative treatment of their effect is necessary.

The second is that the parameter \(\tau\), even when it is sharply defined, is a third parameter. See \([47, 48]\) for a theory of electrical conductivity in charge puddle dominated graphene at low temperatures.

Strange metal in graphene

Negative local resistance due to viscous electron backflow in graphene

Strange metal in graphene

Negative local resistance due to viscous electron backflow in graphene

D. A. Bandurin\textsuperscript{1}, I. Torre\textsuperscript{2,3}, R. Krishna Kumar\textsuperscript{1,4}, M. Ben Shalom\textsuperscript{1,5}, A. Tomadin\textsuperscript{6}, A. Principi\textsuperscript{7}, G. H. Auton\textsuperscript{5}, E. Khestanova\textsuperscript{1,5}, K. S. Novoselov\textsuperscript{5}, I. V. Grigorieva\textsuperscript{1}, L. A. Ponomarenko\textsuperscript{1,4}, A. K. Geim\textsuperscript{1}, M. Polini\textsuperscript{3,6}

\textbf{Figure 1.} Viscous backflow in doped graphene. (a,b) Steady-state distribution of current injected through a narrow slit for a classical conducting medium with zero $\nu$ (a) and a viscous Fermi liquid (b). (c) Optical micrograph of one of our SLG devices. The schematic explains the measurement geometry for vicinity resistance. (d,e) Longitudinal conductivity $\sigma_{xx}$ and $R_V$ for this device as a function of $n$ induced by applying gate voltage. $I = 0.3$ $\mu$A; $L = 1$ $\mu$m. For more detail, see Supplementary Information.
Signature of Navier-Stokes hydrodynamic flow in PdCoO$_2$

Experiment: Successively narrow the channel in factors of 2, measuring the resistance after every step.

Graphene: “a metal that behaves like water”
Entangled quantum matter without quasiparticles

- Is there a connection between strange metals and black holes? Yes, e.g. the SYK model.

- Why do they have the same equilibration time $\sim \hbar/(k_B T)$? Strange metals don’t have quasiparticles and thermalize rapidly; Black holes are “fast scramblers”.

- Theoretical predictions for strange metal transport in graphene agree well with experiments