Functional integral theories of low-dimensional quantum Heisenberg models

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(Received 25 January 1988)
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Fluctuations on a bipartite lattice:
- U(1) gauge theories and monopoles
- Monopole Berry phases
- Valence bond solids
- Deconfined criticality
- Monopole scaling dimensions
Functional integral theories of low-dimensional quantum Heisenberg models

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Fluctuations on a non-bipartite lattice:
- $Z_2$ spin liquid: maps to “odd” $Z_2$ gauge theory
- Visons
- First quantum state with anyons (“topological order”) and time-reversal symmetry
- Toric code and protected quantum memory
Quantum matter without quasiparticles: random fermion models, black holes, and graphene

Interacting Electrons and Quantum Magnetism
Technion, Haifa
June 21, 2016

Subir Sachdev

Talk online: sachdev.physics.harvard.edu
Quantum matter without quasiparticles:

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

2. No quasiparticles

- Superfluid-insulator transition of ultracold bosonic atoms in an optical lattice
- Solvable random fermion Sachdev-Ye-Kitaev (SYK) model
- Charged black hole horizons in anti-de Sitter space
- Graphene

Note: Most states with long-range entanglement, like the fractional quantum Hall states, do have quasiparticles
Quantum matter without quasiparticles:

• Superfluid-insulator transition of ultracold bosonic atoms in an optical lattice
• Solvable random fermion Sachdev-Ye-Kitaev (SYK) model
• Charged black hole horizons in anti-de Sitter space
• Graphene
Superfluid-insulator transition

Ultracold $^{87}$Rb atoms - bosons

$T$  

$T_{KT}$  

Quantum critical  

Superfluid  

Insulator  

$0$  

$\lambda_c$  

$\lambda \sim U/t$
"Boltzmann" theory of Nambu-Goldstone phonons and vortices
Conformal field theory (CFT3) at $T>0$

Quantum matter without quasiparticles
Conformal field theory (CFT3) at $T>0$

Shortest possible “phase coherence” or local thermal equilibration time

$$\sim \frac{\hbar}{k_B T}$$

Local thermal equilibration or phase coherence time, $\tau_\varphi$:

- As $T \to 0$, there is a 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 (the “butterfly effect”):

  $$D(t) = \langle W(t)V(0)W(t)V(0) \rangle \sim c_0 - \epsilon c_1 e^{\lambda_L t},$$

  where we make a (system-dependent) choice to arrange $\epsilon \ll 1$. As $T \to 0$, this Lyapunov time is argued to obey the lower bound

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

- Theories holographically dual to Einstein gravity have the shortest possible $\tau_L = \hbar/(2\pi k_B T)$

A.I. Larkin and Y. N. Ovchinnikov, JETP 28, 6 (1969)
S. H. Shenker and D. Stanford, arXiv:1306.0622
A bound on quantum chaos:

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Quantum matter without quasiparticles

$\approx$ fastest possible many-body quantum chaos
Quantum matter without quasiparticles:

- Superfluid-insulator transition of ultracold bosonic atoms in an optical lattice
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- Graphene
Infinite-range model with quasiparticles

\[ H = \frac{1}{(N)^{1/2}} \sum_{i,j=1}^{N} t_{ij} c_i^\dagger c_j + \ldots \]

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

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

\[ t_{ij} \] are independent random variables with \( \overline{t_{ij}} = 0 \) and \( |t_{ij}|^2 = t^2 \)

Fermions occupying the eigenstates of a \( N \times N \) random matrix
**Infinite-range model with quasiparticles**

Feynman graph expansion in $t_{ij}$, and graph-by-graph average, yields exact equations in the large $N$ limit:

$$G(i\omega) = \frac{1}{i\omega + \mu - \Sigma(i\omega)}, \quad \Sigma(\tau) = t^2 G(\tau)$$

$$G(\tau = 0^-) = Q.$$ 

$G(\omega)$ can be determined by solving a quadratic equation.

Fermi liquid state: Two-body interactions lead to a scattering time of quasiparticle excitations from in (random) single-particle eigenstates which diverges as $\sim T^{-2}$ at the Fermi level.
Infinite-range (SY) model without quasiparticles

\[ H = \frac{1}{(NM)^{1/2}} \sum_{i,j=1}^{N} \sum_{\alpha,\beta=1}^{M} J_{ij} c_{i\alpha}^\dagger c_{i\beta}^\dagger c_{j\beta} c_{j\alpha} \]

\[ c_{i\alpha} c_{j\beta} + c_{j\beta} c_{i\alpha} = 0 \quad , \quad c_{i\alpha} c_{j\beta}^\dagger + c_{j\beta}^\dagger c_{i\alpha} = \delta_{ij} \delta_{\alpha\beta} \]

\[ \frac{1}{M} \sum_{\alpha} c_{i\alpha}^\dagger c_{i\alpha} = Q \]

\( J_{ij} \) are independent random variables with \( \overline{J_{ij}} = 0 \) and \( \overline{J_{ij}^2} = J^2 \)

\( N \to \infty \) at \( M = 2 \) yields spin-glass ground state.

\( N \to \infty \) and then \( M \to \infty \) yields critical strange metal

SYK model without quasiparticles

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

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

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

\( J_{i;j;k,\ell} \) are independent random variables with \( \bar{J}_{i;j;k,\ell} = 0 \) and \( |\bar{J}_{i;j;k,\ell}|^2 = J^2 \)

\( N \rightarrow \infty \) yields same critical strange metal; simpler to study numerically

SYK model without quasiparticles

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

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

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

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

SYK model without quasiparticles

Feynman graph expansion in $J_{ij}$, and graph-by-graph average, yields exact equations in the large $N$ limit:

\[
G(i\omega) = \frac{1}{i\omega + \mu - \Sigma(i\omega)} , \quad \Sigma(\tau) = -J^2 G^2(\tau) G(-\tau) \\
G(\tau = 0^-) = Q.
\]

Low frequency analysis shows that the solutions must be gapless and obey

\[
\Sigma(z) = \mu - \frac{1}{A} \sqrt{z} + \ldots \ , \quad G(z) = \frac{A}{\sqrt{z}}
\]

for some complex $A$. The ground state is a non-Fermi liquid, with a continuously variable density $Q$. 

SYK model without quasiparticles

A better understanding of the above facts can be reached from the perspective of symmetry-protected topological (SPT) phases. As shown recently in Ref.\[14\], the complex SYK model can be thought of as the boundary of a 1D SPT system in the symmetry class AIII. The periodicity of 4 in $N$ arises from the fact that we need to put 4 chains to gap out the boundary degeneracy without breaking the particle-hole symmetry. In the Majorana SYK case, the symmetric Hamiltonian can be constructed as a symmetric matrix in the $\text{Cl}_{0,N}$ Clifford algebra, and the Bott periodicity in the real representation of the Clifford algebra gives rise to a $\mathbb{Z}_8$ classification\[14\]. Here, for the complex SYK case, we can similarly construct the Clifford algebra by dividing one complex fermion into two Majorana fermions, and then we will have a periodicity of 4.

A. Green's function

From the above definition of retarded Green's function, we can relate them to the imaginary time Green's function as defined in Eq. (16),

$$G_R(\tau) = G(i\tau^+) + i\eta.$$  

In Fig. 3, we show a comparison between the imaginary part of the Green's function from large $N$, and from the exact diagonalization computation. The spectral function from ED is particle-hole symmetric for all $N$.

We identify the infinite time limit of $G_B$ as the Edward-Anderson order parameter $q_{EA}$, which can characterize long-time memory of spin-glass:

$$q_{EA} = \lim_{t \to \infty} G_B(t).$$

Then $q_{EA} \neq 0$ indicates that $G_B(\tau) \propto \eta$. This is quite different from the fermionic case, where we have $G_F(z) \propto 1/p_z$; this inverse square-root behavior also holds in the bosonic case without spin glass order\[1\]. Fig. 10 is our result from ED, with a comparison between $G_B$ with $G_F$. It is evident that the behavior of $G_B$ is qualitatively different from $G_F$, and so an inverse square-root behavior is ruled out. Instead, we can clearly see that, as system size gets larger, $G_B$'s peak value increases much faster than the $G_F$'s peak value. This supports the presence of spin glass order.

Large $N$ solution of equations for $G$ and $\Sigma$ agree well with exact diagonalization of the finite $N$ Hamiltonian.

However, exact diagonalization of the same model with hard-core bosons indicates the presence of spin-glass order in the ground state.

W. Fu and S. Sachdev, arXiv: 1603.05246
SYK model without quasiparticles

Local fermion density of states

$$\rho(\omega) = -\text{Im } G(\omega) \sim \begin{cases} \omega^{-1/2}, & \omega > 0 \\ e^{-2\pi \mathcal{E}} |\omega|^{-1/2}, & \omega < 0. \end{cases}$$

$\mathcal{E}$ encodes the particle-hole asymmetry

While $\mathcal{E}$ determines the low energy spectrum, it is determined by the total fermion density $Q$:

$$Q = \frac{1}{4} (3 - \tanh(2\pi \mathcal{E})) - \frac{1}{\pi} \tan^{-1} (e^{2\pi \mathcal{E}}).$$

SYK model without quasiparticles

At non-zero temperature, $T$, the Green’s function also fully determined by $\mathcal{E}$.

$$G^R(\omega) = \frac{-iCe^{-i\theta}}{(2\pi T)^{1-2\Delta}} \frac{\Gamma \left( \Delta - \frac{i\omega}{2\pi T} + i\mathcal{E} \right)}{\Gamma \left( 1 - \Delta - \frac{i\omega}{2\pi T} + i\mathcal{E} \right)}$$

where $\Delta = 1/4$ and \( e^{2\pi\mathcal{E}} = \frac{\sin(\pi\Delta + \theta)}{\sin(\pi\Delta - \theta)} \).

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A. Georges and O. Parcollet PRB 59, 5341 (1999)
SYK model without quasiparticles

The entropy per site, $S$, has a non-zero limit as $T \to 0$. This is not due to an extensive degeneracy, but due to an energy level spacing $\sim e^{-aN}$ in the entire many-body spectrum all the way down to the ground state. The $T \to 0$ limit of $S$ obeys

$$
\left( \frac{\partial S}{\partial Q} \right)_T = - \left( \frac{\partial \mu}{\partial T} \right)_Q = 2\pi \mathcal{E}
$$

Note that $S$ and $\mathcal{E}$ involve low-lying states, while $Q$ depends upon all states, and details of the UV structure

$$
S(Q, T \to 0) = 2\pi \int_{-\infty}^{f^{-1}(Q)} dx \, x f'(x), \quad f(x) = \frac{3 - \tanh(2\pi x)}{4} - \frac{\tan^{-1}(e^{2\pi x})}{\pi}
$$

Infinite-range (SYK) model without quasiparticles

After integrating the fermions, the partition function can be written as a path integral with an action $S$ analogous to a Luttinger-Ward functional

$$Z = \int \mathcal{D}G(\tau_1, \tau_2)\mathcal{D}\Sigma(\tau_1, \tau_2) \exp(-NS)$$

$$S = \ln \det [\delta(\tau_1 - \tau_2)(\partial_{\tau_1} + \mu) - \Sigma(\tau_1, \tau_2)]$$

$$+ \int d\tau_1 d\tau_2 \Sigma(\tau_1, \tau_2) [G(\tau_2, \tau_1) + (J^2/2)G^2(\tau_2, \tau_1)G^2(\tau_1, \tau_2)]$$

At frequencies $\ll J$, the time derivative in the determinant is less important, and without it the path integral is invariant under the reparametrization and gauge transformations

$$\tau = f(\sigma)$$

$$G(\tau_1, \tau_2) = [f'(\sigma_1)f'(\sigma_2)]^{-1/4} \frac{g(\sigma_1)}{g(\sigma_2)} G(\sigma_1, \sigma_2)$$

$$\Sigma(\tau_1, \tau_2) = [f'(\sigma_1)f'(\sigma_2)]^{-3/4} \frac{g(\sigma_1)}{g(\sigma_2)} \Sigma(\sigma_1, \sigma_2)$$

where $f(\sigma)$ and $g(\sigma)$ are arbitrary functions.
Infinite-range (SYK) model without quasiparticles

Let us write the large $N$ saddle point solutions of $S$ as

$$G_s(\tau_1 - \tau_2) \sim (\tau_1 - \tau_2)^{-1/2}, \quad \Sigma_s(\tau_1 - \tau_2) \sim (\tau_1 - \tau_2)^{-3/2}.$$  

These are not invariant under the reparametrization symmetry but are invariant only under a SL(2,R) subgroup under which

$$f(\tau) = \frac{a\tau + b}{c\tau + d}, \quad ad - bc = 1.$$  

So the (approximate) reparametrization symmetry is spontaneously broken.

**Reparametrization zero mode**

Expand about the saddle point by writing

$$G(\tau_1, \tau_2) = [f'(\tau_1)f'(\tau_2)]^{1/4}G_s(f(\tau_1) - f(\tau_2))$$

(and similarly for $\Sigma$) and obtain an effective action for $f(\tau)$. This action does not vanish because of the time derivative in the determinant which is not reparameterization invariant.

J. Maldacena and D. Stanford, arXiv:1604.07818

See also A. Kitaev, unpublished, and J. Polchinski and V. Rosenhaus, arXiv:1601.06768
Infinite-range (SYK) model without quasiparticles

However the effective action must vanish for SL(2,R) transformations because $G_s, \Sigma_s$ are invariant under it. In this manner we obtain the effective action as a Schwarzian

$$NS_{\text{eff}} = -\frac{N\gamma}{4\pi^2} \int d\tau \{ f, \tau \} , \quad \{ f, \tau \} \equiv \frac{f'''}{f'} - \frac{3}{2} \left( \frac{f''}{f'} \right)^2,$$

where the co-efficient $\gamma$ determines the specific heat, $C$

$$C = T \frac{\partial S}{\partial T} = N\gamma T$$

The Schwarzian describes fluctuations of the energy operator with scaling dimension $h = 2$. 

J. Maldacena and D. Stanford, arXiv:1604.07818

See also A. Kitaev, unpublished, and J. Polchinski and V. Rosenhaus, arXiv:1601.06768
Infinite-range (SYK) model without quasiparticles

The Schwarzian effective action implies that the SYK model saturates the lower bound on the Lyapunov time

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

Apart from the energy operator associated with the Schwarzian, there are an infinite number of other scalar operators with irrational scaling dimensions given by the roots of

\[ \tan \left( \frac{\pi (2h - 1)}{4} \right) = \frac{1 - 2h}{3} \]

\[ \Rightarrow \quad h = 3.77354 \ldots, \ 5.67946 \ldots, \ 7.63197 \ldots, \ 9.60396 \ldots, \ldots \]
Holographic Metals and the Fractionalized Fermi Liquid

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(Received 23 June 2010; published 4 October 2010)

We show that there is a close correspondence between the physical properties of holographic metals near charged black holes in anti–de Sitter (AdS) space, and the fractionalized Fermi liquid phase of the lattice Anderson model. The latter phase has a “small” Fermi surface of conduction electrons, along with a spin liquid of local moments. This correspondence implies that certain mean-field gapless spin liquids are states of matter at nonzero density realizing the near-horizon, $\text{AdS}_2 \times \mathbb{R}^2$ physics of Reissner-Nordström black holes.
The non-zero $T \to 0$ entropy density, $S$, matches the Bekenstein-Hawking-Wald entropy density of extremal $\text{AdS}_2$ horizons, and the dependence of the fermion Green’s function on $\omega$, $T$, and $\mathcal{E}$, matches that of a Dirac fermion in $\text{AdS}_2$ (as computed by T. Faulkner, Hong Liu, J. McGreevy, and D. Vegh, PRD 83, 125002 (2011)).

S. Sachdev, PRL 105, 151602 (2010)

More recently, it was noted that the relation $\left(\partial S / \partial Q\right)_T = 2\pi \mathcal{E}$ also matches between SYK and gravity, where $\mathcal{E}$, the electric field on the horizon, also determines the spectral asymmetry of the Dirac fermion.

S. Sachdev, PRX 5, 041025 (2015)
The same Schwarzian effective action describes low energy fluctuations on the boundary theory of gravity theories with AdS$_2$ near-horizon geometries (including the AdS-Reissner-Nordstrom solution of Einstein-Maxwell theory in 4 space-time dimensions). And the co-efficient of the Schwarzian, $N\gamma/4\pi^2$, determines the specific heat $C = N\gamma T$. 

The Schwarzian effective action implies that both the SYK model and the AdS$_2$ theories *saturate* the lower bound on the Lyapunov time

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

This is additional evidence for an AdS$_2$ dual of the SYK model.

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Quantum matter without quasiparticles:

- Superfluid-insulator transition of ultracold bosonic atoms in an optical lattice
- Solvable random fermion Sachdev-Ye-Kitaev (SYK) model
- Charged black hole horizons in anti-de Sitter space
- Graphene
Graphene

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

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

**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:

\[ L_0 = \frac{\kappa}{\sigma T} \approx \frac{\pi^2 k_B^2}{3e^2} \approx 2.45 \times 10^{-8} \frac{W \cdot \Omega}{K^2}. \]
Transport in Strange Metals

For a strange metal with a “relativistic” Hamiltonian, hydrodynamic, holographic, and memory function methods yield Lorentz ratio \( L = \kappa/(T\sigma) \)

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

\( Q \rightarrow \) 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.

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

Predicted strange metal

Quantum critical

Dirac liquid

Electron
Fermi liquid

Hole
Fermi liquid

$T(K)$

$\mu < 0$

$\mu > 0$

$\sim \sqrt{n}$

$\lambda \ln \Lambda \sqrt{n}$

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

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

M. Müller and S. Sachdev, PRB 78, 115419 (2008)
Impurity scattering dominates

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 $L = L_0$
Graphene

Predicted strange metal

$T(K)$

Quantum critical

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

Wiedemann-Franz obeyed !

Wiedemann-Franz violated !
Lorentz ratio $L = \kappa / (T \sigma)$

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

$Q \rightarrow$ electron density; $\mathcal{H} \rightarrow$ enthalpy density

$\sigma_Q \rightarrow$ quantum critical conductivity

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


J. Crossno et al., Science 351, 1058 (2016)
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. Bandurin¹, I. Torre²,³, R. Krishna Kumar¹,⁴, M. Ben Shalom¹,⁵, A. Tomadin⁶, A. Principi⁷, G. H. Auton⁵, E. Khestanova¹,⁵, K. S. Novoselov⁵, I. V. Grigorieva¹, L. A. Ponomarenko¹,⁴, A. K. Geim¹, M. Polini³,⁶

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 \( v \) (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.
Entangled quantum matter without quasiparticles

- No quasiparticle excitations

- Shortest possible “phase coherence” time, fastest possible local equilibration time, or fastest possible Lyapunov time towards quantum chaos, all of order $\frac{\hbar}{k_B T}$

- Realization in solvable SYK model, which saturates the lower bound on the Lyapunov time. Its properties have some similarities to non-rational, large central charge CFT2s.

- Remarkable match between SYK and quantum gravity of black holes with AdS$_2$ horizons, including a SL(2,R)-invariant Schwarzian effective action for thermal energy fluctuations.

- Experiments on graphene agree well with predictions of a theory of a nearly relativistic quantum liquid without quasiparticles.