Quantum matter without quasiparticles

Low Energy Challenges for High Energy Physicists

Perimeter Institute, Waterloo, Canada

May 29, 2014

Subir Sachdev

PERIMETER INSTITUTE FOR THEORETICAL PHYSICS

Talk online: sachdev.physics.harvard.edu

PHYSICS

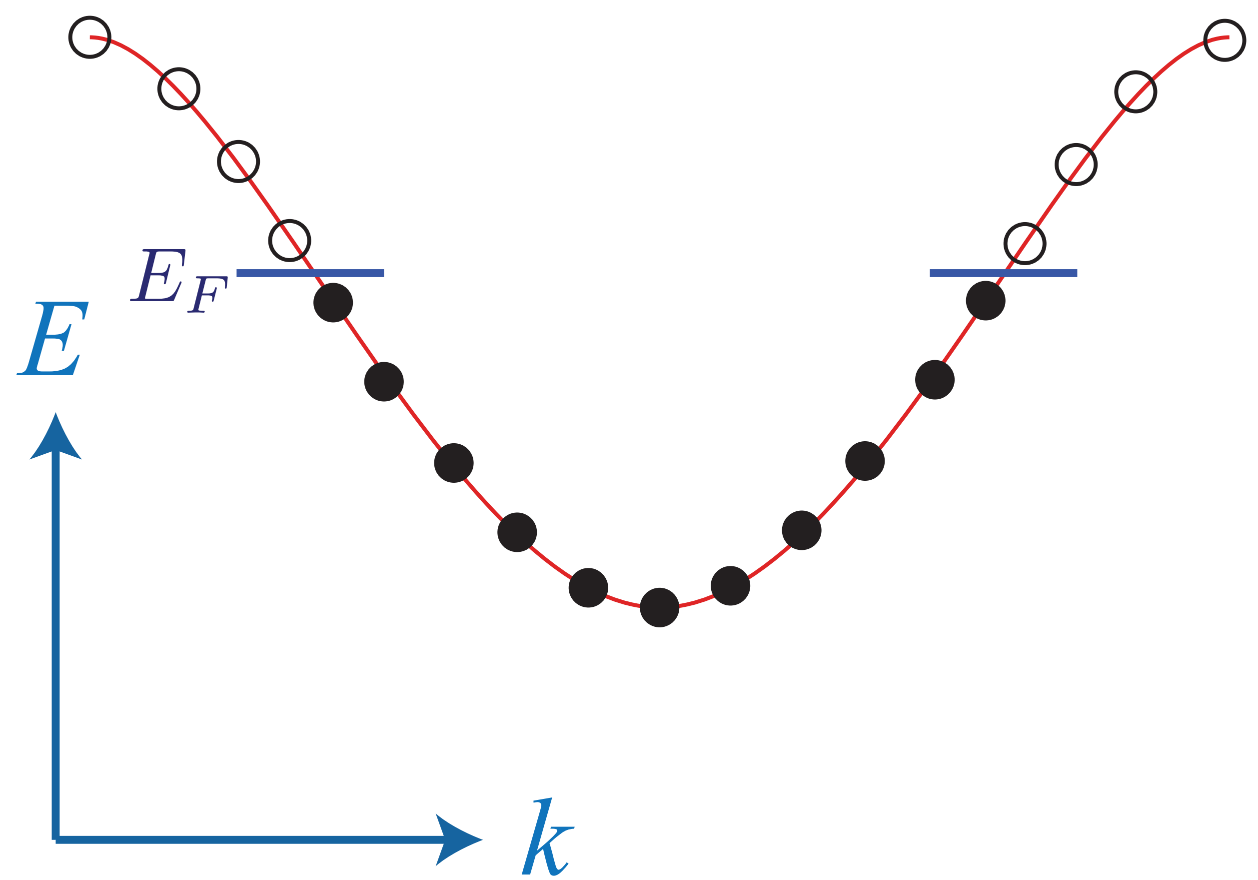
VE RI TAS

HARVARD

Foundations of quantum many body theory:

1. Ground states connected adiabatically to independent electron states

Metals

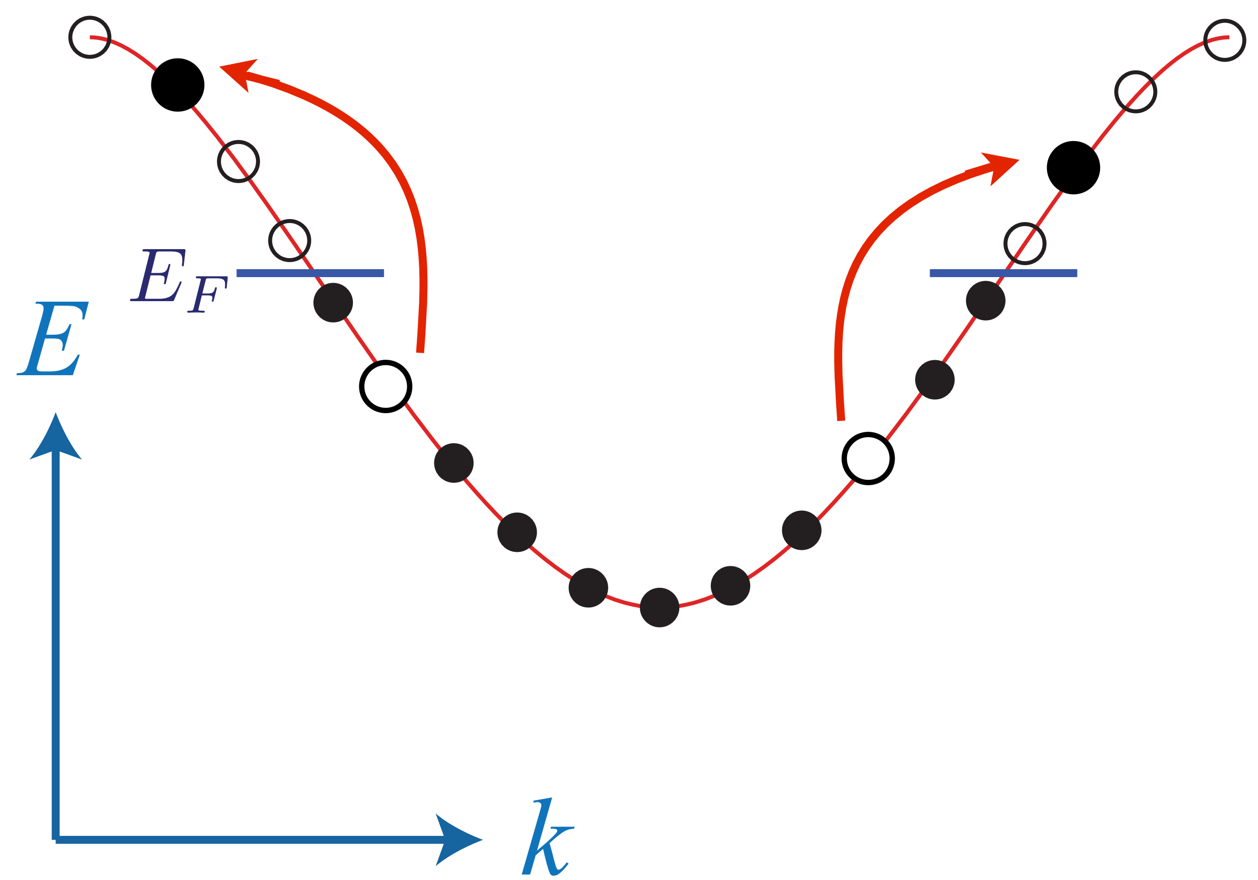


Foundations of quantum many body theory:

1. Ground states connected adiabatically to independent electron states

2. Boltzmann-Landau theory of quasiparticles

Metals



Modern phases of quantum matter:

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

2. Boltzmann-Landau theory of quasiparticles

Modern phases of quantum matter:

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

2. Boltzmann-Landau theory of quasiparticles

Famous examples:

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.

Modern phases of quantum matter:

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

2. Boltzmann-Landau theory of quasiparticles

Famous examples:

Electrons in one dimensional wires form the Luttinger liquid. The quanta of density oscillations ("phonons") are a *quasiparticle* basis of the low-energy Hilbert space. Similar comments apply to magnetic insulators in one dimension.

Modern phases of quantum matter:

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

2. *No quasiparticles*

Modern phases of quantum matter:

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

2. No quasiparticles

Only 2 examples:

1. Conformal field theories in spatial dimension $d > 1$

2. Quantum critical metals in dimension $d=2$

Outline

1. Conformal field theories in 2+1 dimensions

   Superfluid-insulator transition

   of ultracold bosonic atoms in an optical lattice

2. Theory of a non-Fermi liquid

   Non-quasiparticle transport at the

   Ising-nematic quantum critical point

Outline

1. Conformal field theories in 2+1 dimensions

   Superfluid-insulator transition

   of ultracold bosonic atoms in an optical lattice

2. Theory of a non-Fermi liquid

   Non-quasiparticle transport at the

   Ising-nematic quantum critical point

nature physics

ARTICLES

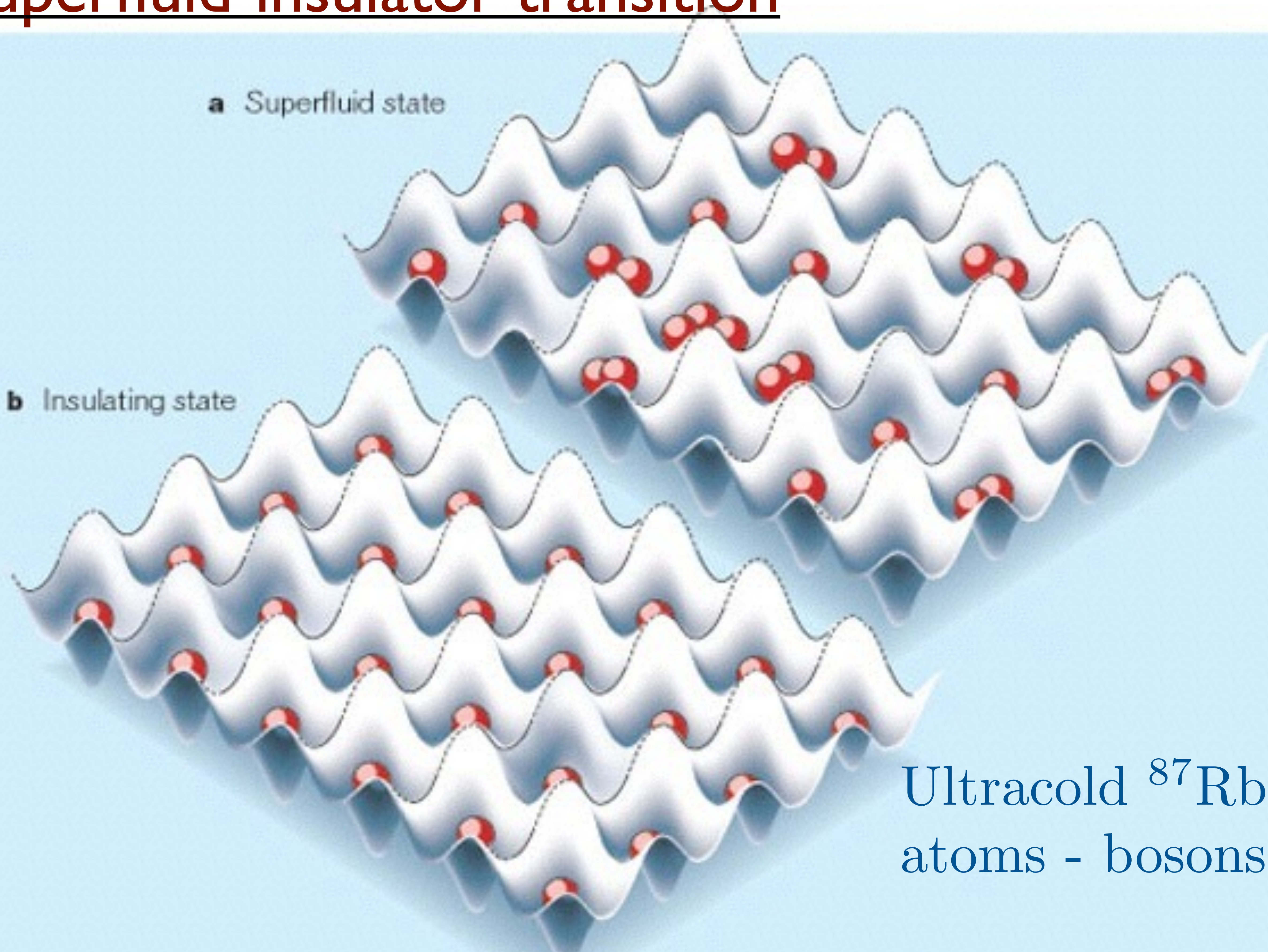
PUBLISHED ONLINE: 23 MARCH 2014 | DOI: 10.1038/NPHYS2913

The dynamics of quantum criticality revealed by quantum Monte Carlo and holography

William Witczak-Krempa
Perimeter

Erik Sorensen
McMaster

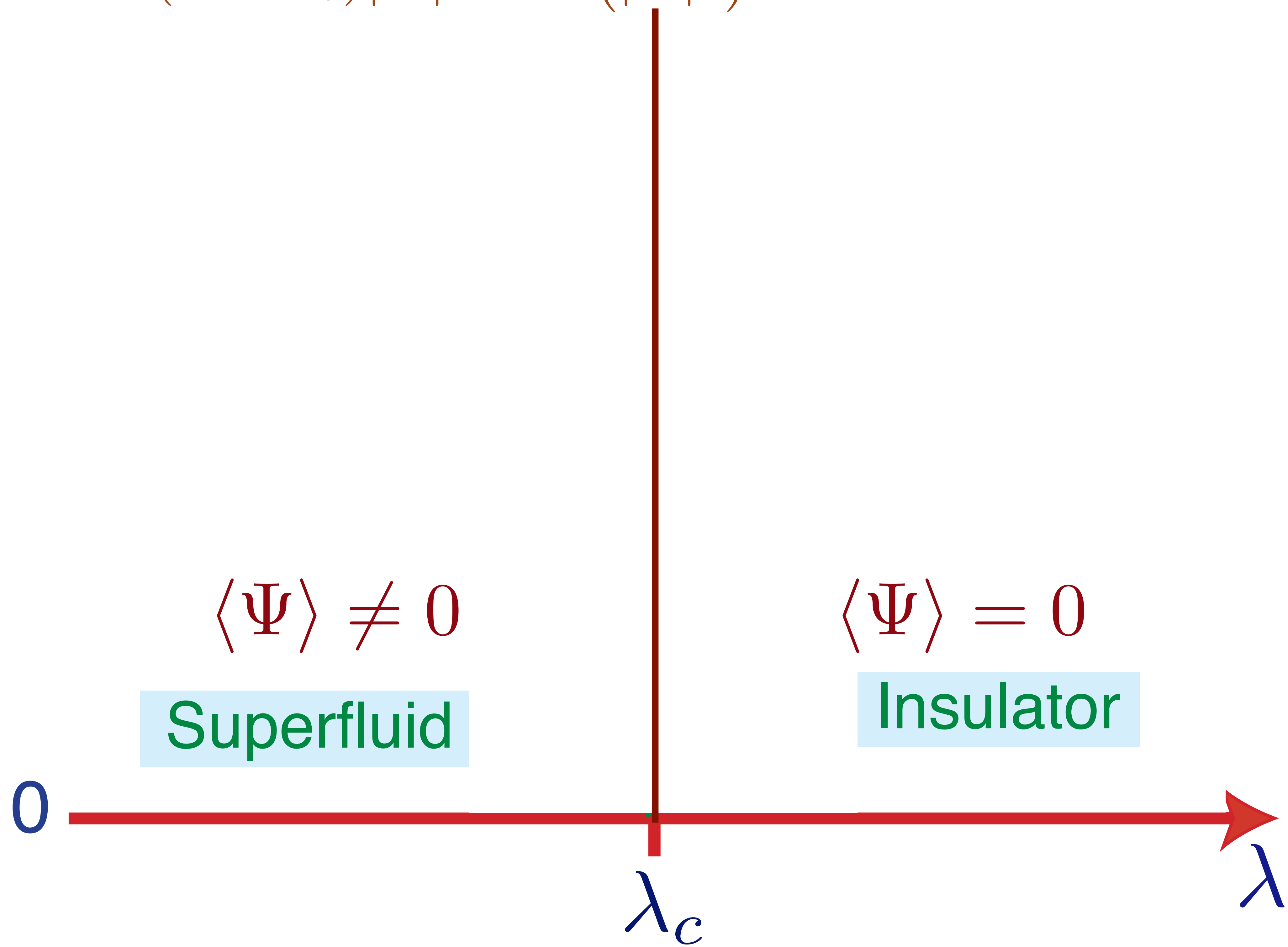
Superfluid-insulator transition



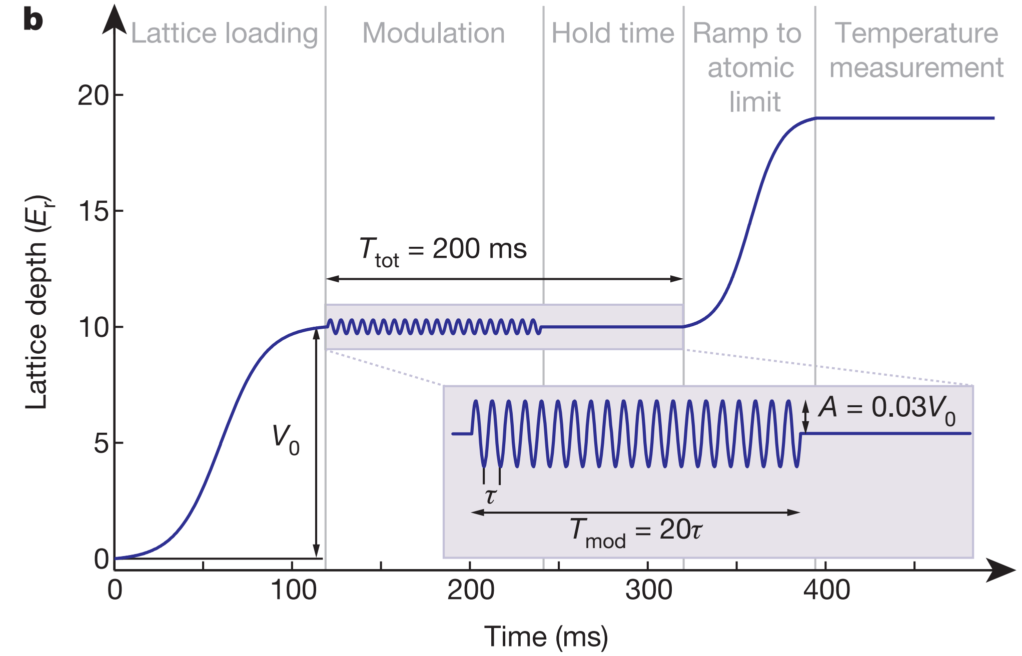
Ultracold ^{87}Rb atoms - bosons

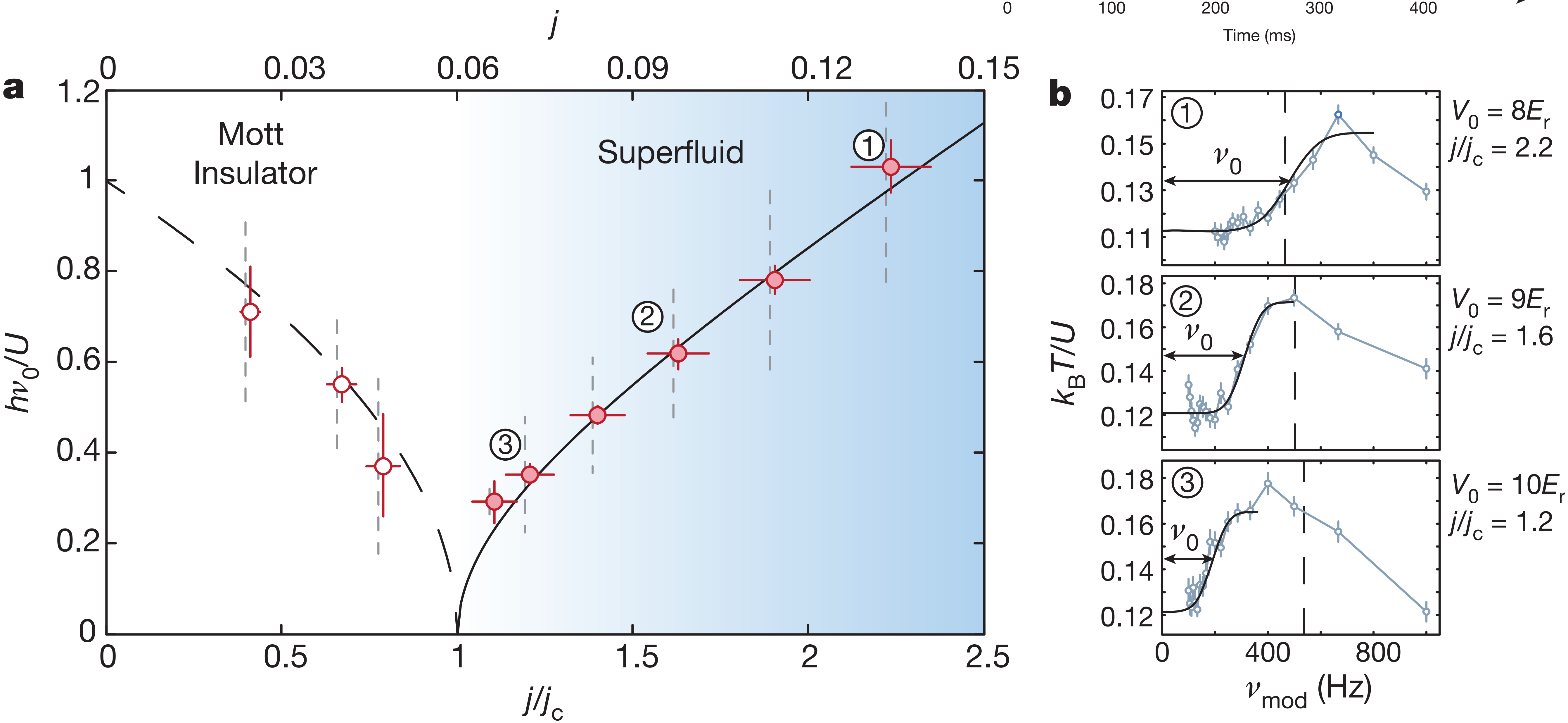
M. Greiner, O. Mandel, T. Esslinger, T. W. Hänsch, and I. Bloch, *Nature* **415**, 39 (2002).

$$
\begin{aligned}
\mathcal{S} &= \int d^2r dt \left[|\partial_t \Psi|^2 - c^2 |\nabla_r \Psi|^2 - V(\Psi) \right] \\
V(\Psi) &= (\lambda - \lambda_c)|\Psi|^2 + u \left(|\Psi|^2 \right)^2
\end{aligned}
$$



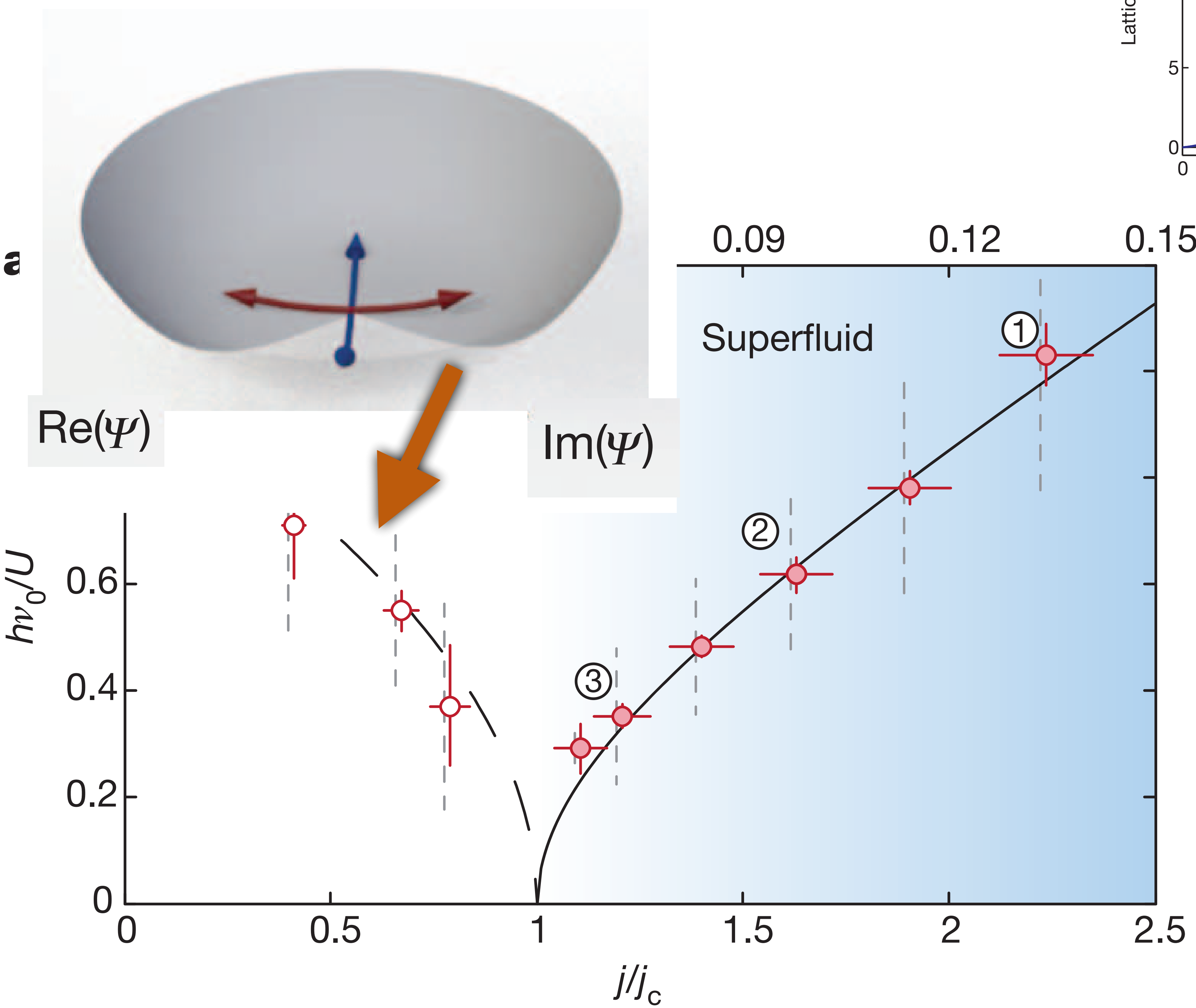
Excitations across the superfluid-insulator transition

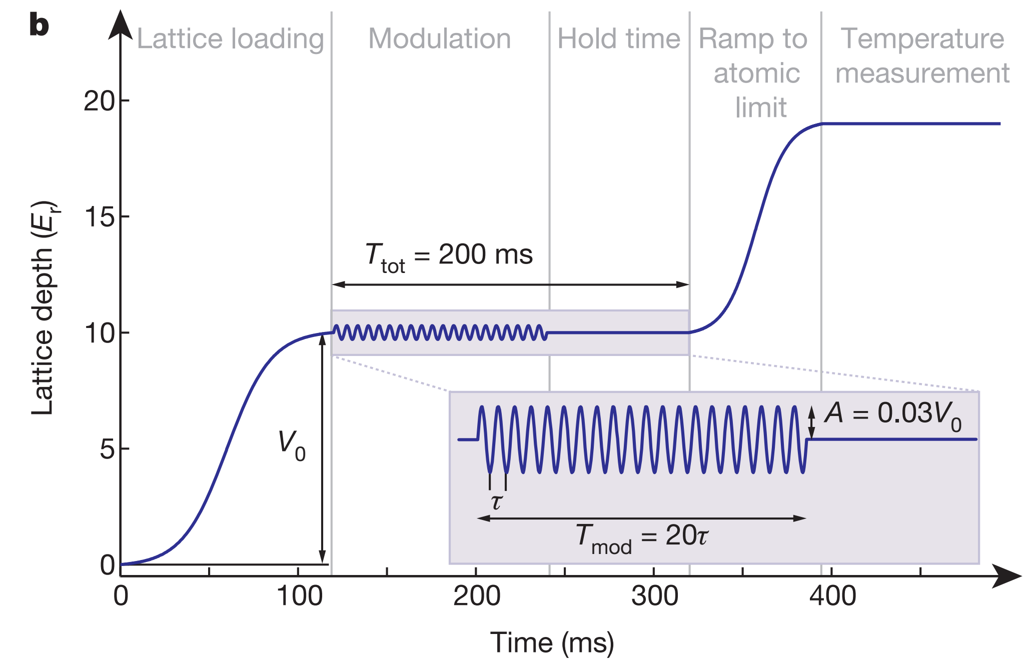


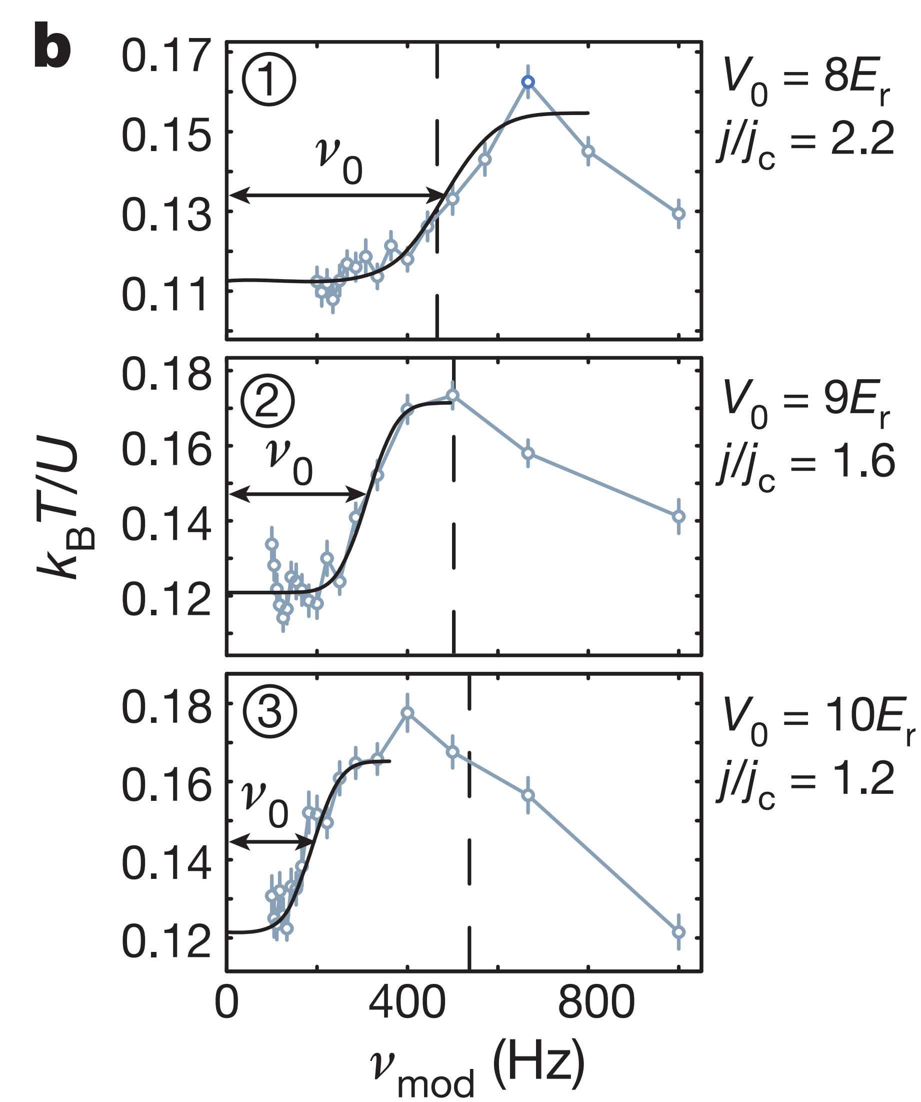


Manuel Endres, Takeshi Fukuhara, David Pekker, Marc Cheneau, Peter Schaub, Christian Gross, Eugene Demler, Stefan Kuhr, and Immanuel Bloch, *Nature* **487**, 454 (2012).

Excitations across the … transition

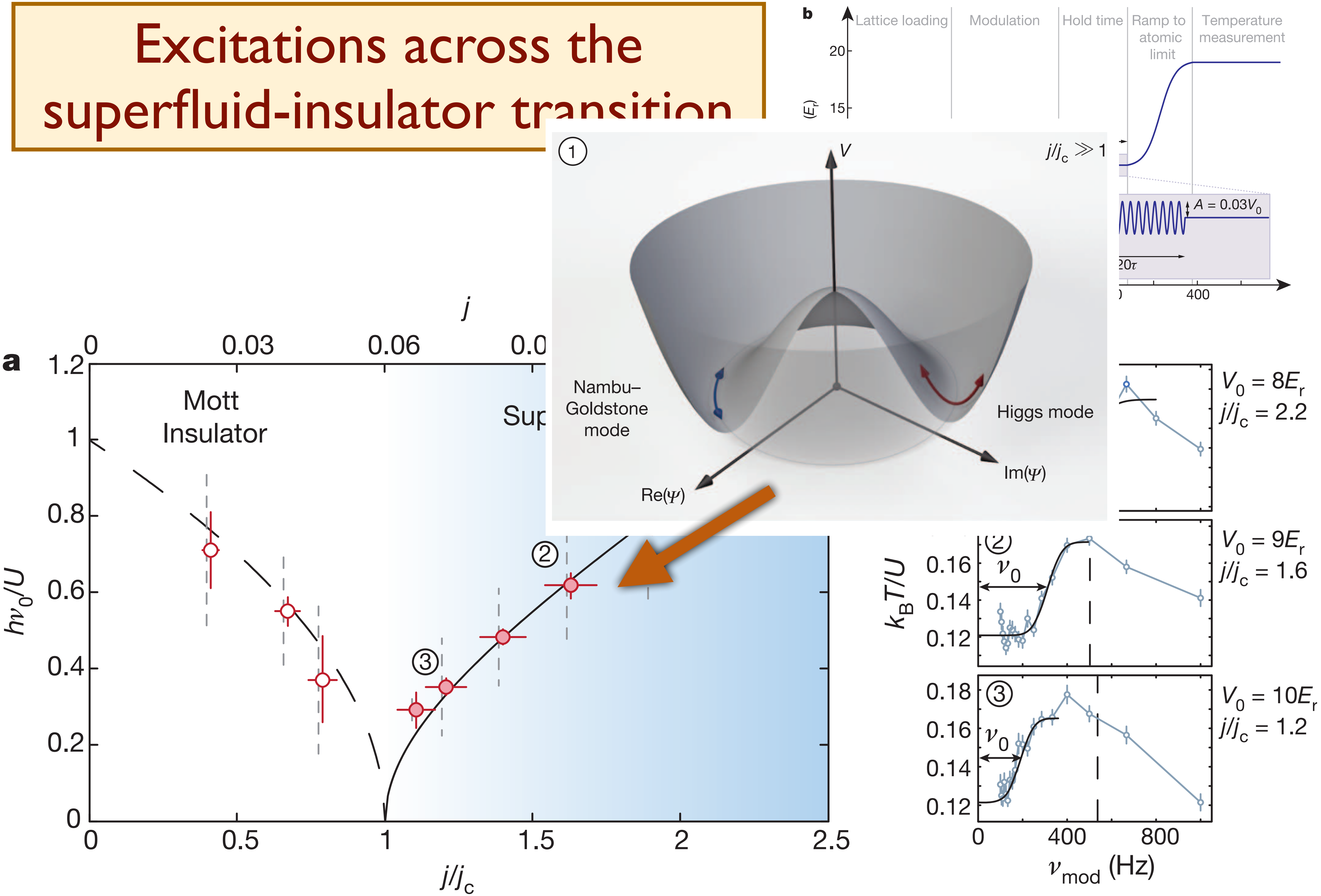






Manuel Endres, Takeshi Fukuhara, David Pekker, Marc Cheneau, Peter Schaub, Christian Gross, Eugene Demler, Stefan Kuhr, and Immanuel Bloch, *Nature* **487**, 454 (2012).

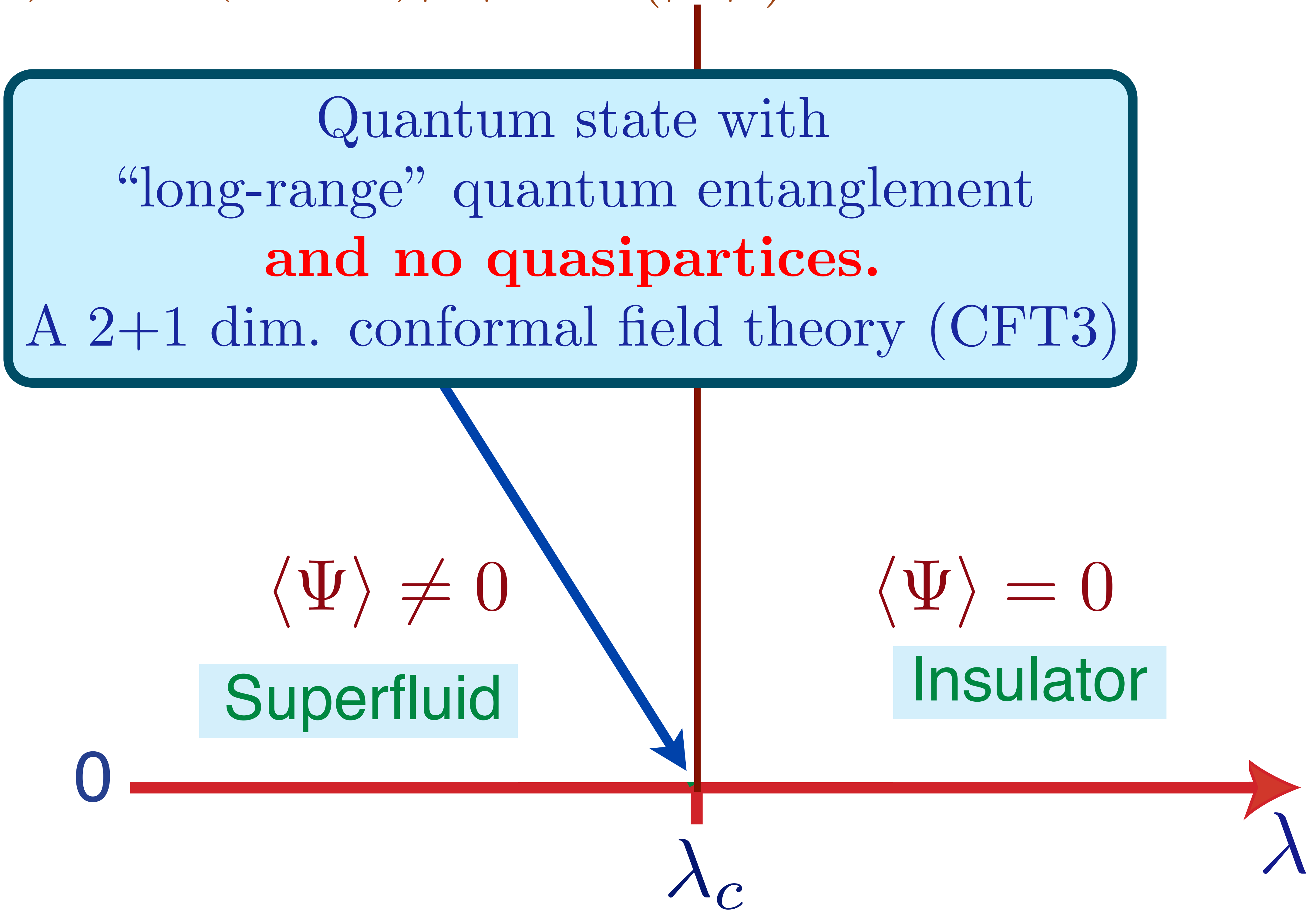
Excitations across the superfluid-insulator transition

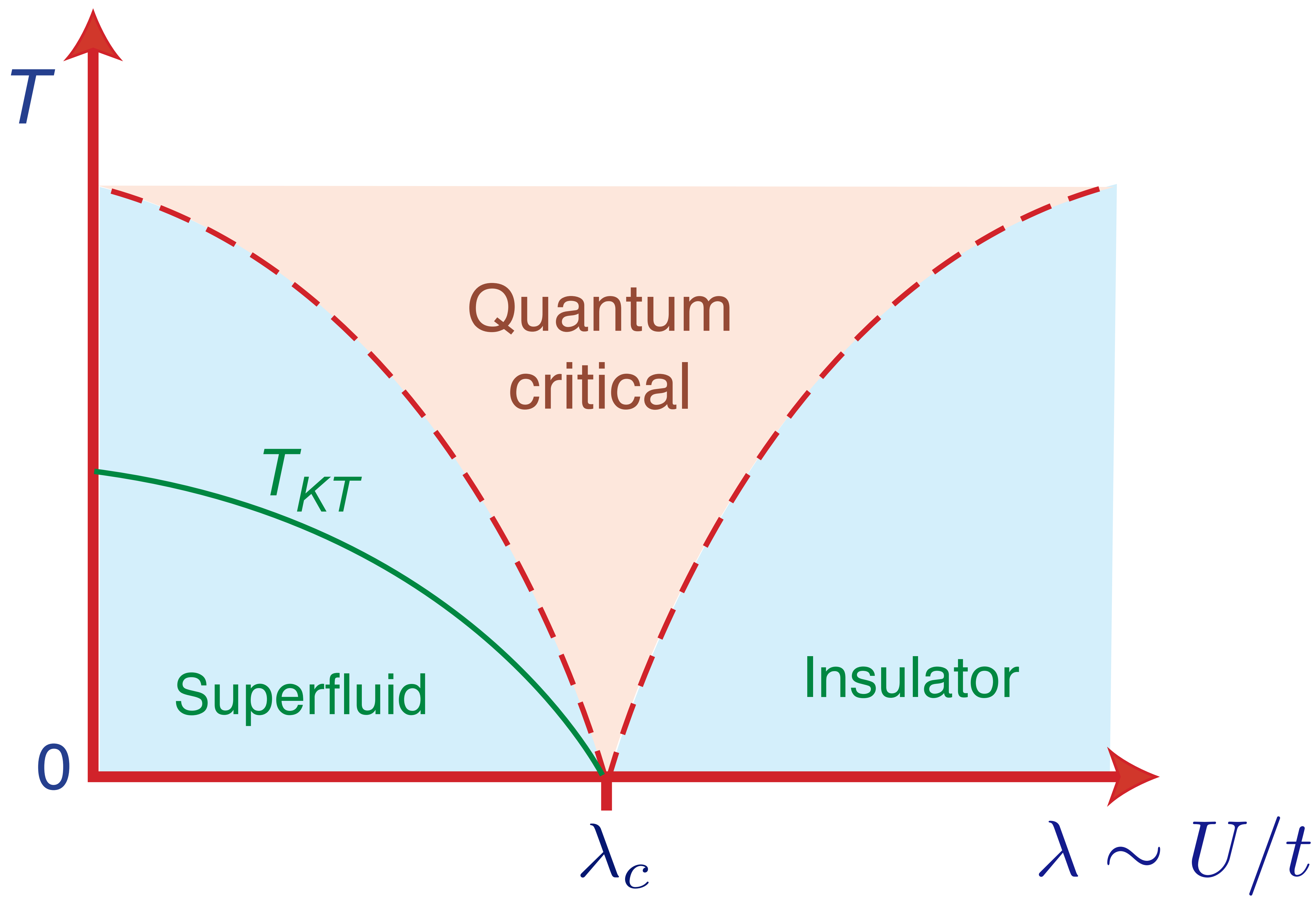


Manuel Endres, Takeshi Fukuhara, David Pekker, Marc Cheneau, Peter Schaub, Christian Gross, Eugene Demler, Stefan Kuhr, and Immanuel Bloch, *Nature* **487**, 454 (2012).

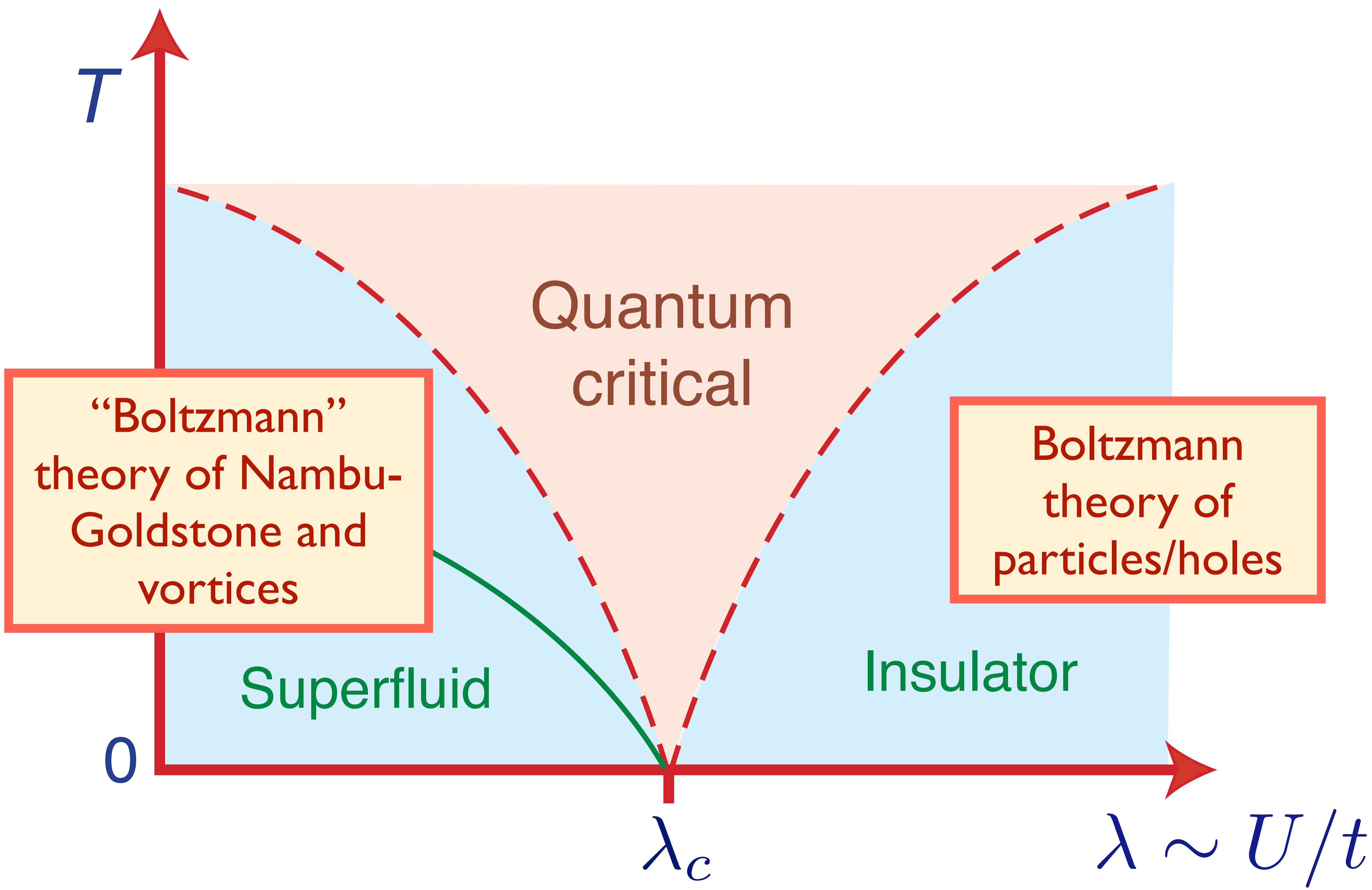
$$\mathcal{S} = \int d^2 r dt \left[|\partial_t \Psi|^2 - c^2 |\nabla_r \Psi|^2 - V(\Psi) \right]$$

$$V(\Psi) = (\lambda - \lambda_c)|\Psi|^2 + u\left(|\Psi|^2\right)^2$$

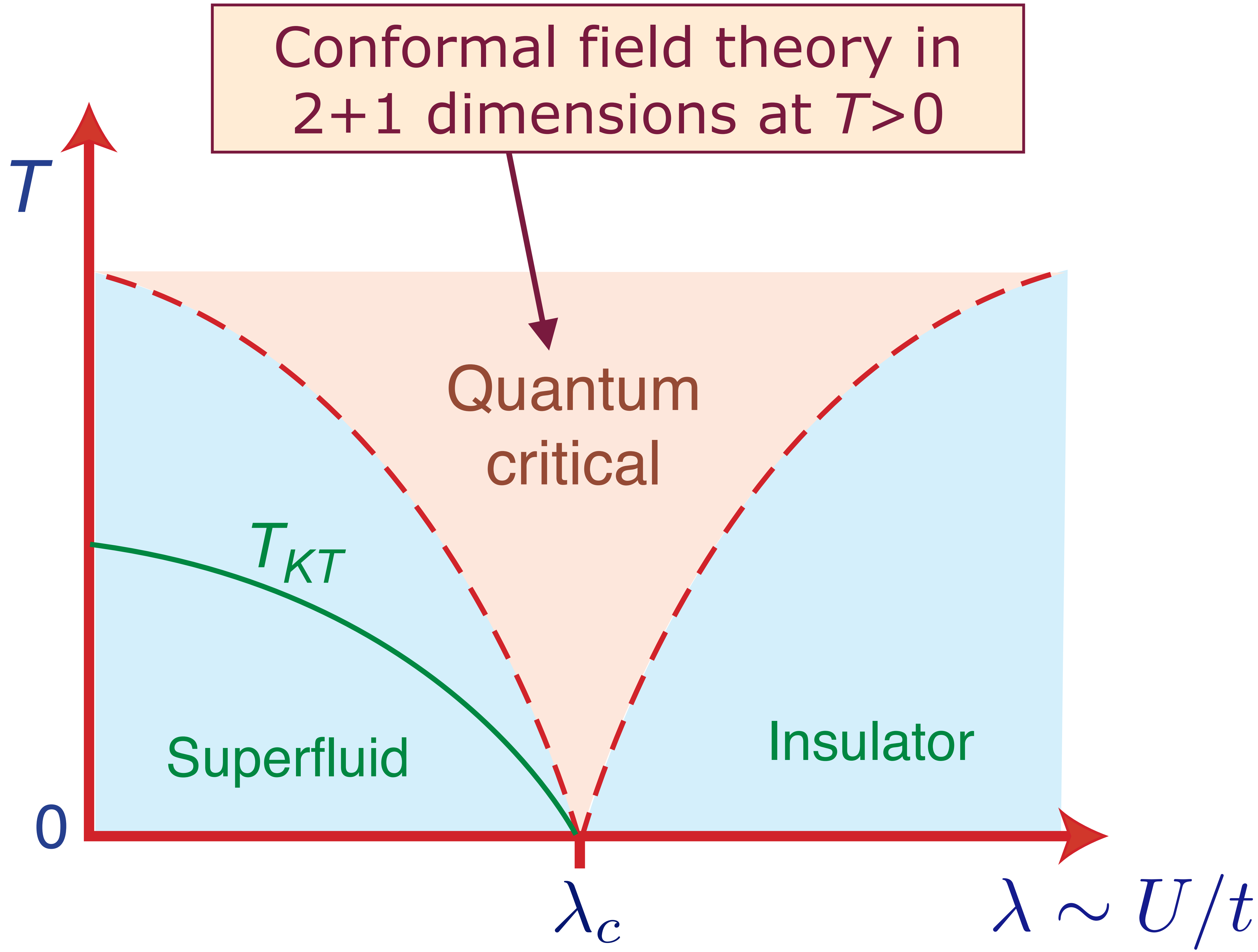


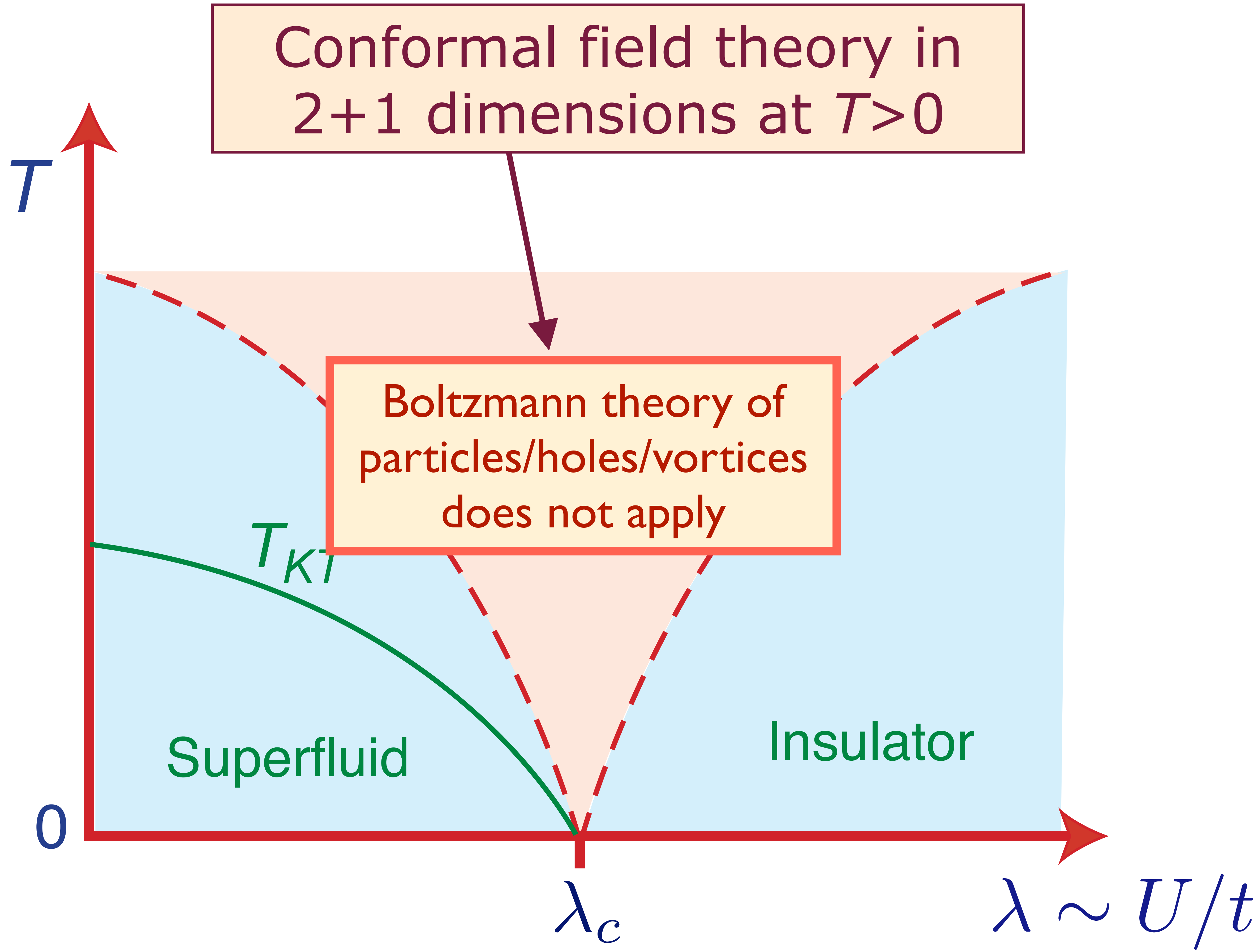


T
Quantum
critical
T_{KT}
Superfluid
Insulator
0
λ_c
$\lambda \sim U/t$

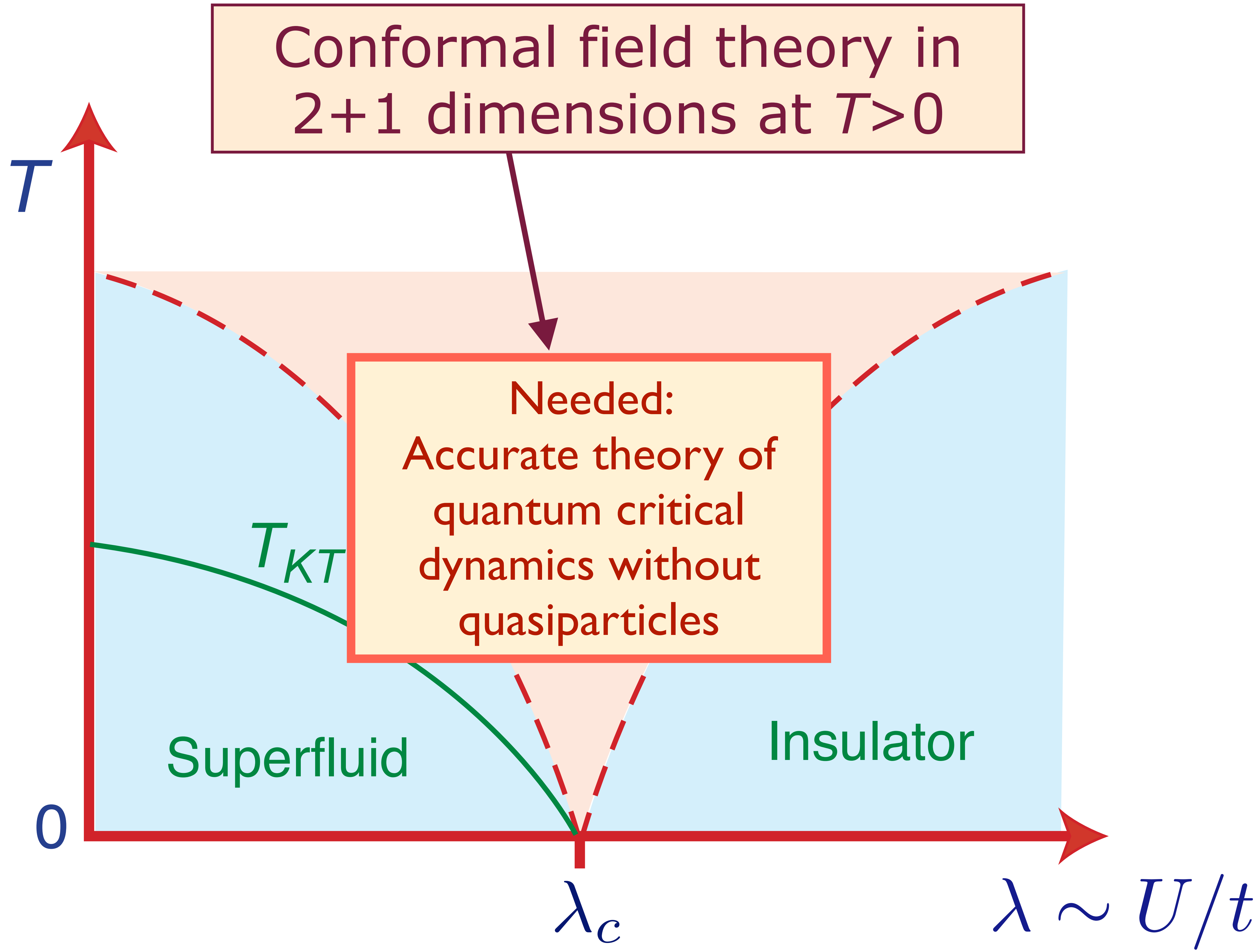


T
0
Quantum critical
“Boltzmann” theory of Nambu-Goldstone and vortices
Boltzmann theory of particles/holes
Superfluid
Insulator
λ_c
$\lambda \sim U/t$


Conformal field theory in 2+1 dimensions at $T>0$
T
Quantum critical
T_{KT}
Superfluid
Insulator
0
λ_c
$\lambda \sim U/t$



Conformal field theory in 2+1 dimensions at $T>0$
T
Boltzmann theory of particles/holes/vortices does not apply
T_{KT}
Superfluid
Insulator
0
λ_c
$\lambda \sim U/t$



Conformal field theory in 2+1 dimensions at T>0
T
Needed:
Accurate theory of quantum critical dynamics without quasiparticles
T_{KT}
Superfluid
Insulator
0
λ_c
$\lambda \sim U/t$

Traditional CMT

- Identify quasiparticles and their dispersions

- Compute scattering matrix elements of quasiparticles (or of collective modes)

Traditional CMT

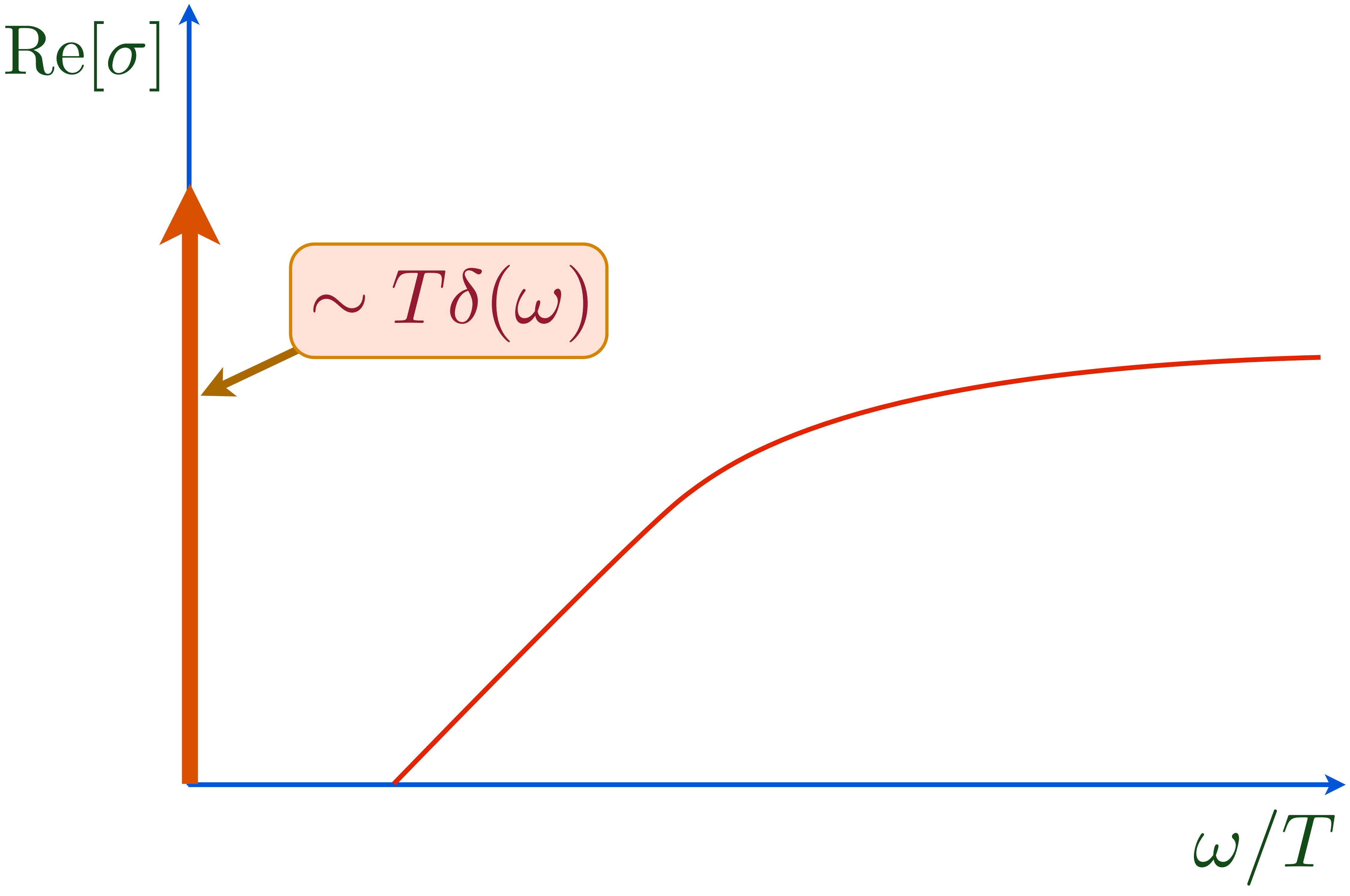
- Identify quasiparticles and their dispersions

- Compute scattering matrix elements of quasiparticles (or of collective modes)

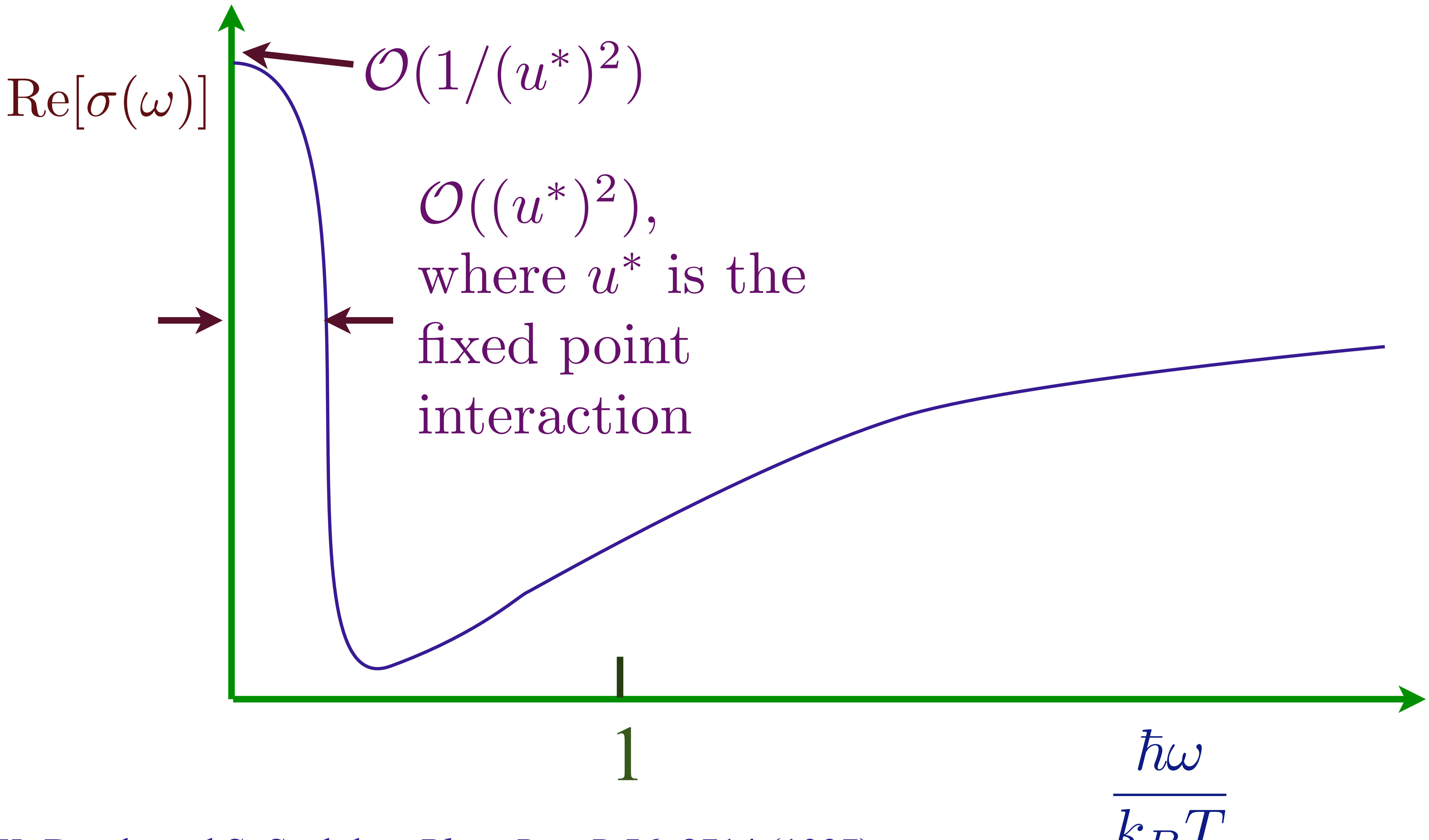
- These parameters are input into a quantum Boltzmann equation

- Deduce dissipative and dynamic properties at non-zero temperatures

Quasiparticle view of quantum criticality (Boltzmann equation): Electrical transport for a free CFT3



Quasiparticle view of quantum criticality (Boltzmann equation): Electrical transport for a (weakly) interacting CFT3



K. Damle and S. Sachdev, *Phys. Rev. B* **56**, 8714 (1997).

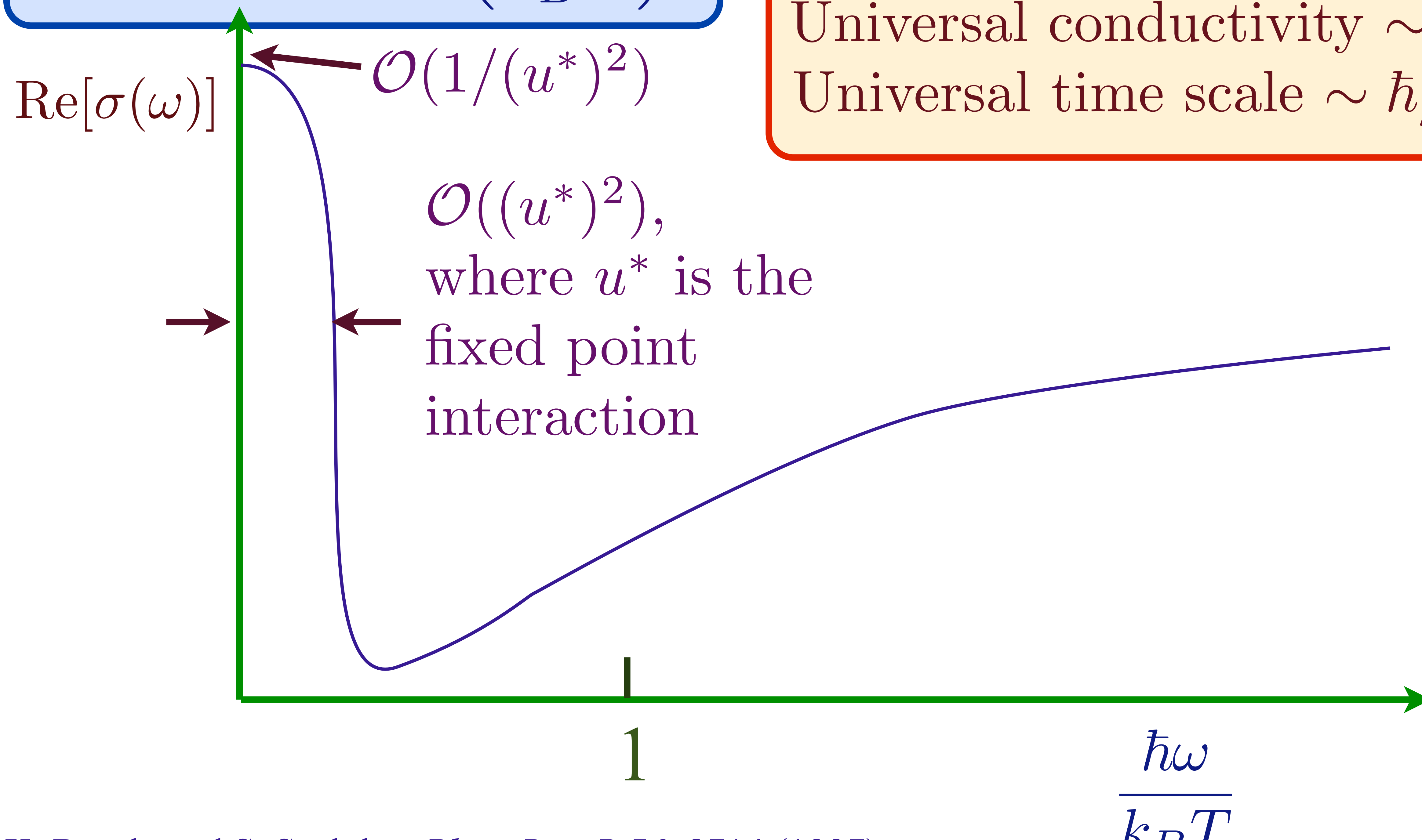
Quasiparticle view of quantum criticality (Boltzmann equation): Electrical transport for a (weakly) interacting CFT3

$$\sigma(\omega, T) = \frac{e^2}{h} \Sigma\left(\frac{\hbar\omega}{k_B T}\right) \quad ;$$

$\Sigma \rightarrow$ a universal function

Universal conductivity $\sim e^2/h$

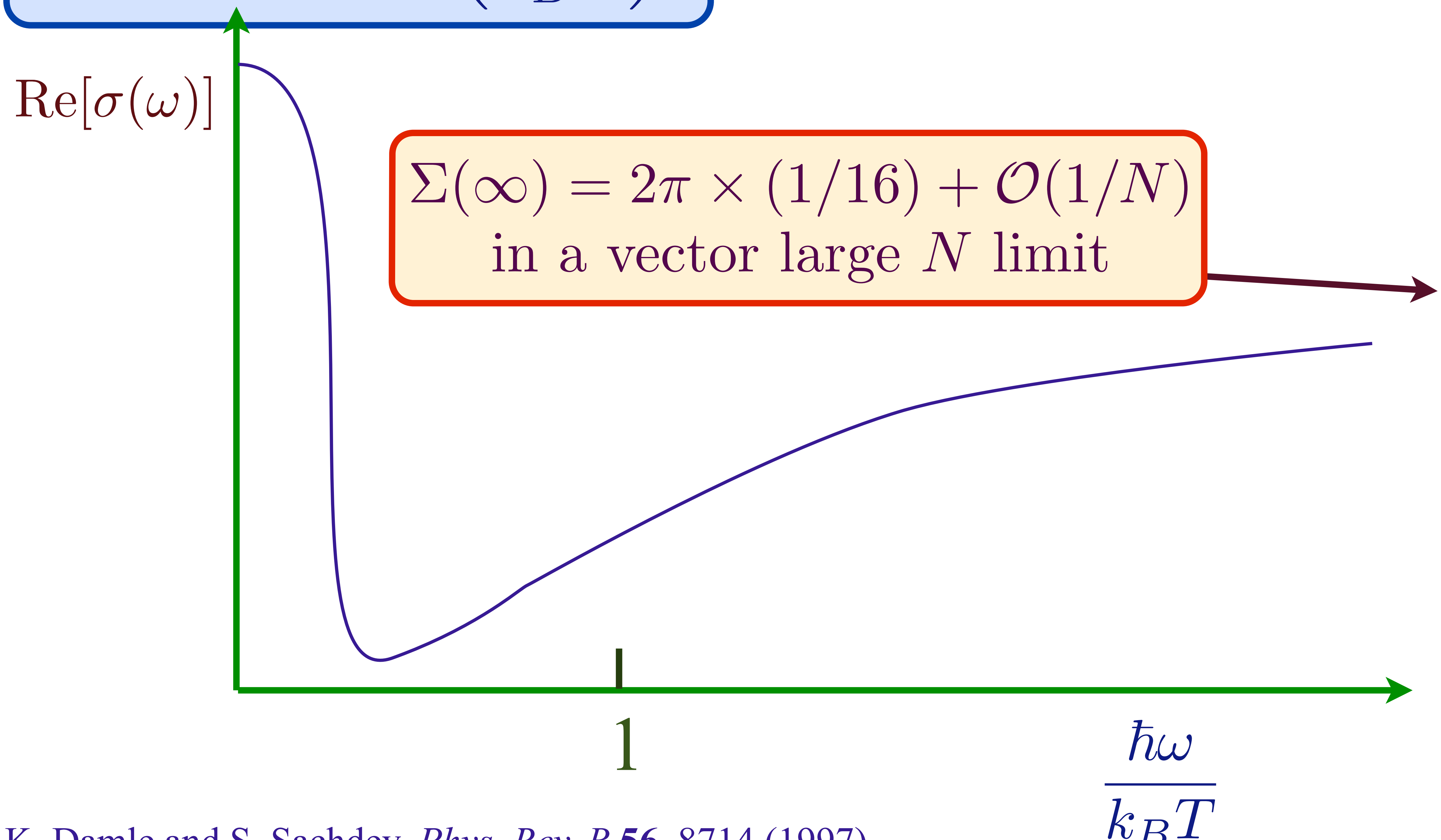
Universal time scale $\sim \hbar/k_B T$



K. Damle and S. Sachdev, *Phys. Rev. B* **56**, 8714 (1997).

Quasiparticle view of quantum criticality (Boltzmann equation): Electrical transport for a (weakly) interacting CFT3

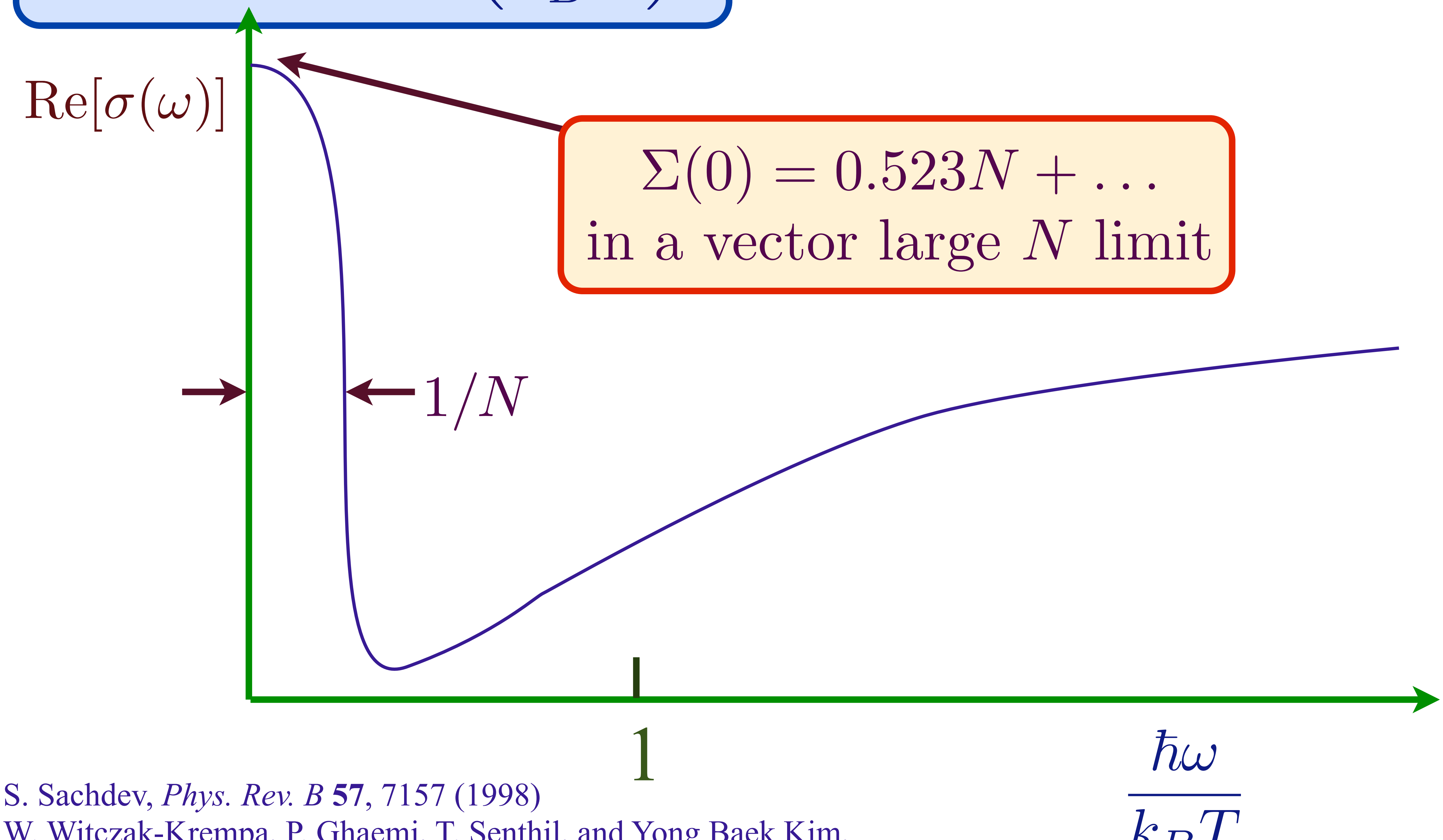
$$\sigma(\omega, T) = \frac{e^2}{h} \Sigma \left(\frac{\hbar\omega}{k_B T} \right)$$



K. Damle and S. Sachdev, *Phys. Rev. B* **56**, 8714 (1997).

Quasiparticle view of quantum criticality (Boltzmann equation): Electrical transport for a (weakly) interacting CFT3

$$\sigma(\omega, T) = \frac{e^2}{h}\Sigma\left(\frac{\hbar\omega}{k_B T}\right)$$

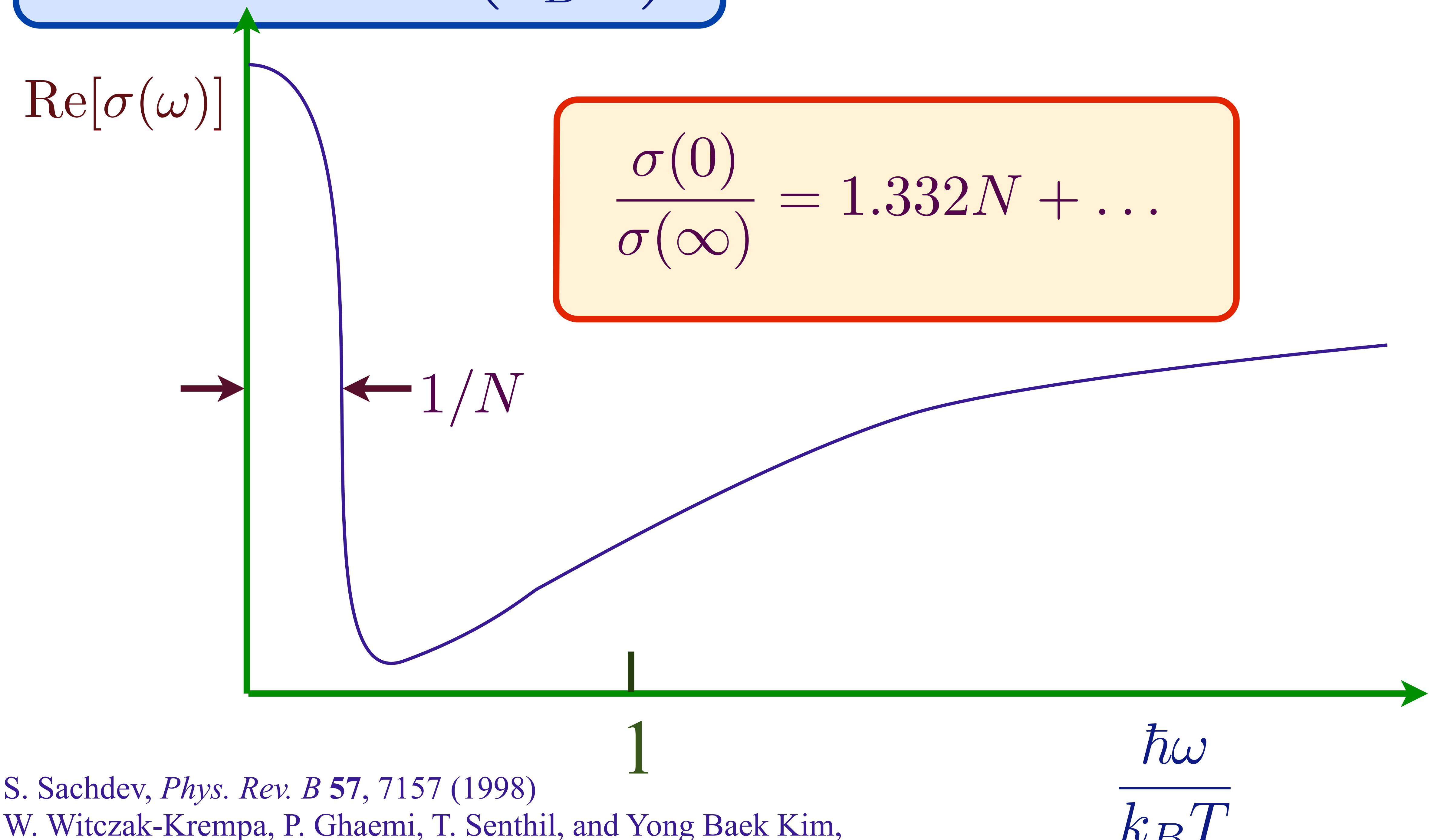


S. Sachdev, *Phys. Rev. B* **57**, 7157 (1998)
W. Witczak-Krempa, P. Ghaemi, T. Senthil, and Yong Baek Kim, *Phys. Rev. B* **86**, 24102 (2012)

Quasiparticle view of quantum criticality (Boltzmann equation): Electrical transport for a (weakly) interacting CFT3

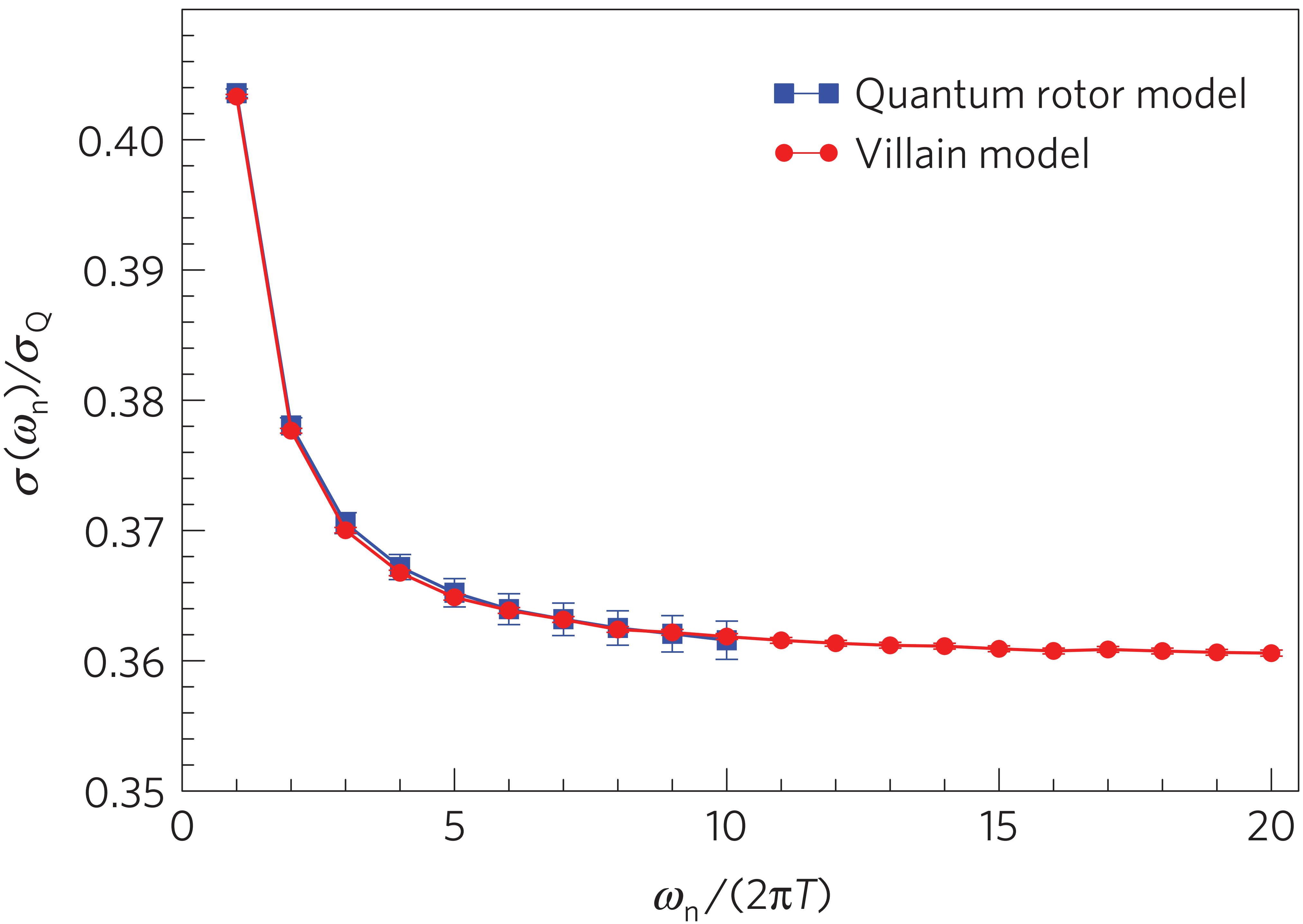
$$\sigma(\omega, T) = \frac{e^2}{h} \Sigma \left(\frac{\hbar\omega}{k_B T} \right)$$

$$\frac{\sigma(0)}{\sigma(\infty)} = 1.332N + \ldots$$



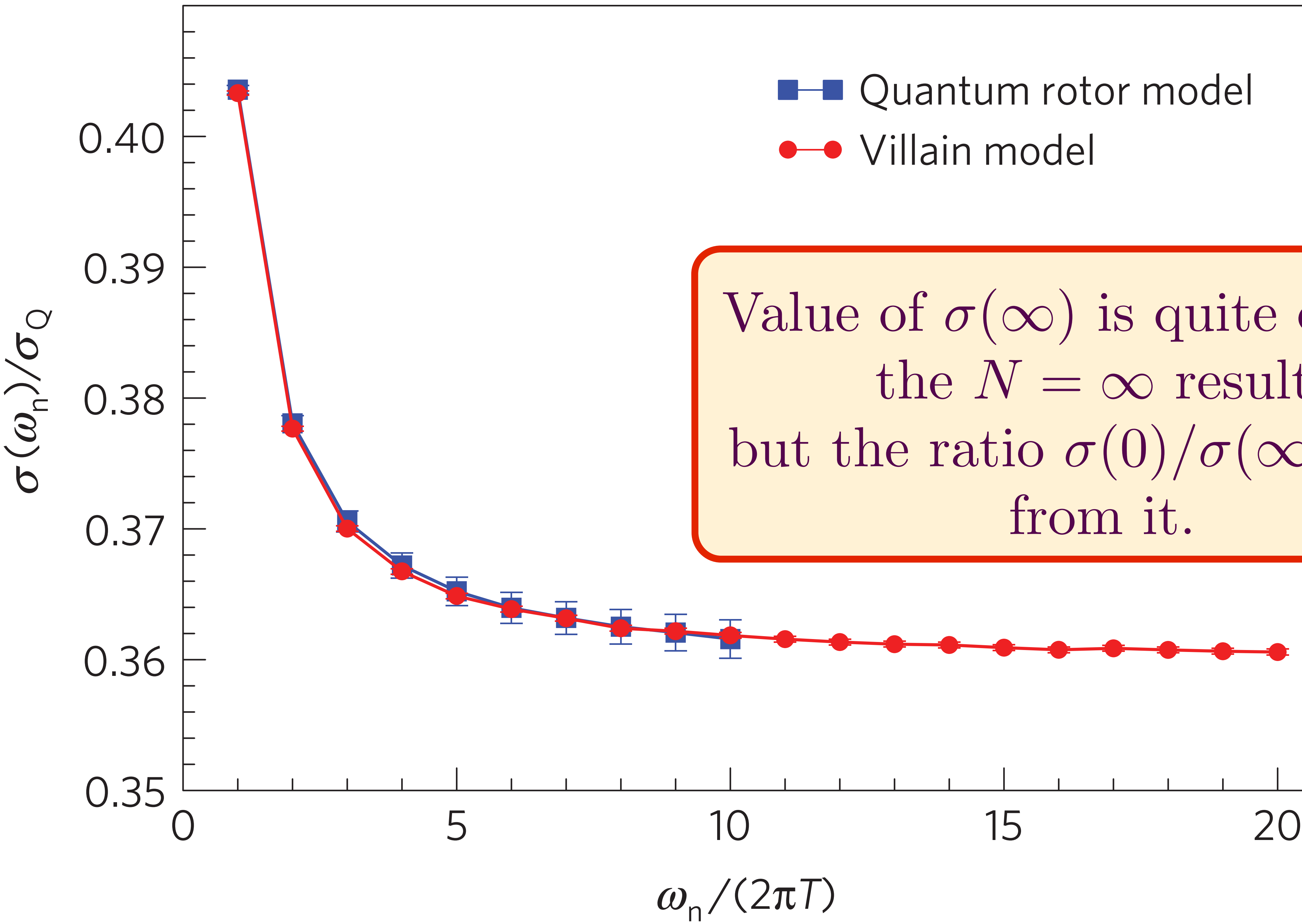
S. Sachdev, *Phys. Rev. B* **57**, 7157 (1998)
W. Witczak-Krempa, P. Ghaemi, T. Senthil, and Yong Baek Kim, *Phys. Rev. B* **86**, 24102 (2012)

Quantum Monte Carlo for lattice bosons



W. Witczak-Krempa, E. Sorensen, and S. Sachdev, Nature Physics **10**, 361 (2014)
See also K. Chen, L. Liu, Y. Deng, L. Pollet, and N. Prokof'ev, Phys. Rev. Lett. **112**, 030402 (2013)

Quantum Monte Carlo for lattice bosons



Value of $\sigma(\infty)$ is quite close to the $N = \infty$ result, but the ratio $\sigma(0)/\sigma(\infty)$ is far from it.

W. Witczak-Krempa, E. Sorensen, and S. Sachdev, Nature Physics **10**, 361 (2014)
See also K. Chen, L. Liu, Y. Deng, L. Pollet, and N. Prokof'ev, Phys. Rev. Lett. **112**, 030402 (2013)

Traditional CMT

- Identify quasiparticles and their dispersions

- Compute scattering matrix elements of quasiparticles (or of collective modes)

- These parameters are input into a quantum Boltzmann equation

- Deduce dissipative and dynamic properties at non-zero temperatures

Traditional CMT

- Identify quasiparticles and their dispersions
- Compute scattering matrix elements of quasiparticles (or of collective modes)
- These parameters are input into a quantum Boltzmann equation
- Deduce dissipative and dynamic properties at non-zero temperatures

Dynamics without quasiparticles

- Start with strongly interacting CFT without particle- or wave-like excitations
- Compute scaling dimensions and OPE co-efficients of operators of the CFT

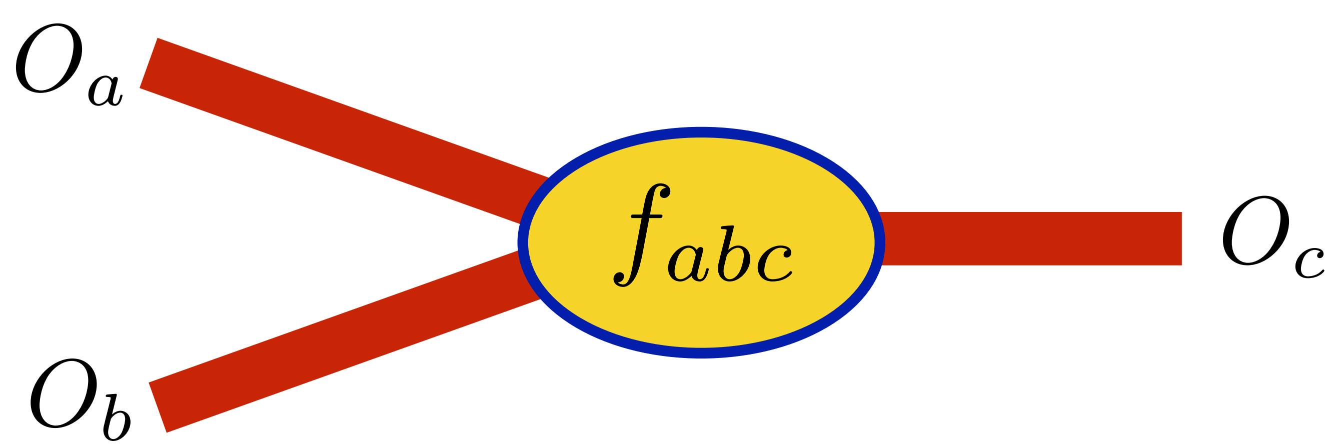
Basic characteristics of CFTs

Primary operators of CFT, $O_a(x)$, obey (at $T = 0$):

$$\boxed{\langle O_a(x) O_b(0) \rangle = \frac{\delta_{ab}}{|x|^{2\Delta_a}}}$$

where Δ_a is their scaling dimension. Their "interactions" are determined by the OPE (considering scalar operators only)

$$\boxed{\lim_{x' \to x} \langle O_a(x') O_b(x) O_c(0) \rangle = \frac{f_{abc}}{|x|^{\Delta_a + \Delta_b + \Delta_c}}}$$



The values of $\boxed{\{\Delta_a, f_{abc}\}}$ determine (in principle) all observable properties of the CFT, as constrained by conformal Ward identities. For the Wilson-Fisher CFT3, systematic methods exist to compute (in principle) all the $\{\Delta_a, f_{abc}\}$, and we will assume this data is *known*. This knowledge will be taken as an *input* to the computation of the finite T dynamics

Traditional CMT

- Identify quasiparticles and their dispersions

- Compute scattering matrix elements of quasiparticles (or of collective modes)

- These parameters are input into a quantum Boltzmann equation

- Deduce dissipative and dynamic properties at non-zero temperatures

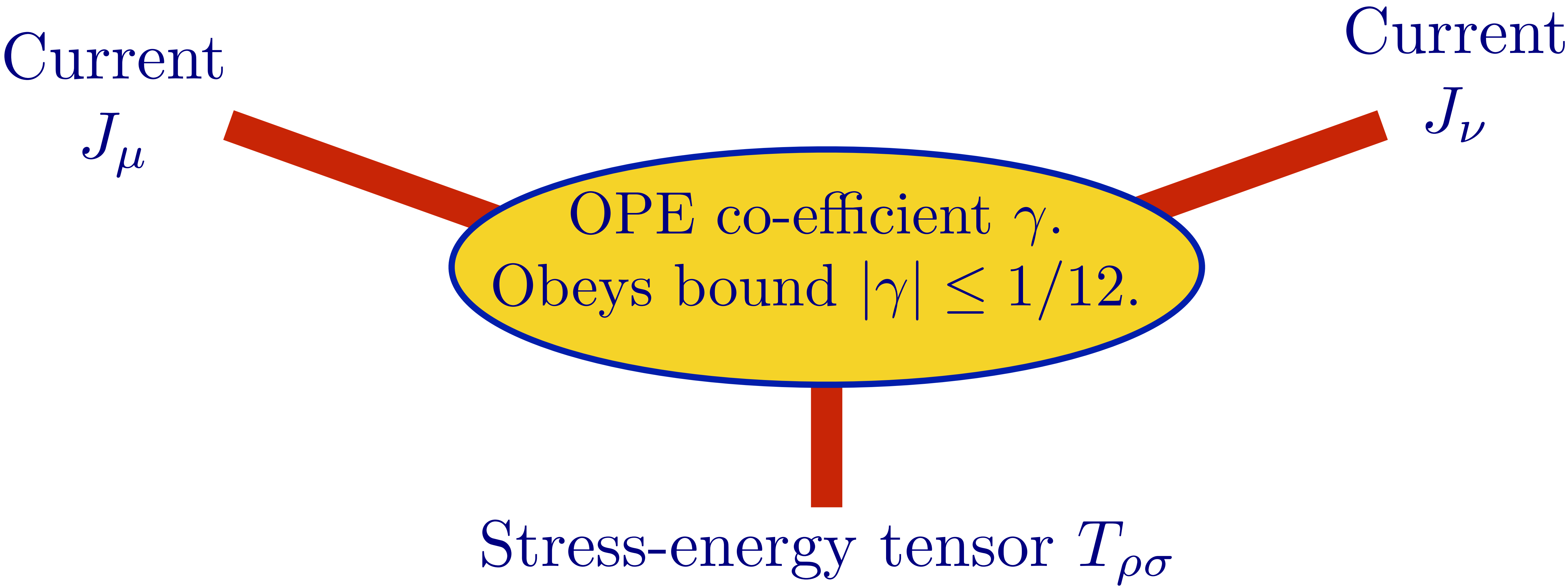
Dynamics without quasiparticles

- Start with strongly interacting CFT without particle- or wave-like excitations

- Compute scaling dimensions and OPE co-efficients of operators of the CFT

Traditional CMT	Dynamics without quasiparticles
Identify quasiparticles and their dispersions	Start with strongly interacting CFT without particle- or wave-like excitations
Compute scattering matrix elements of quasiparticles (or of collective modes)	Compute scaling dimensions and OPE co-efficients of operators of the CFT
These parameters are input into a quantum Boltzmann equation	Relate OPE co-efficients to couplings of an effective gravitational theory on AdS
Deduce dissipative and dynamic properties at non-zero temperatures	Non-zero *T* dynamics of CFT maps to dynamics of a "horizon" in (Einstein's) gravitational theory

Physical picture of electrical transport in a CFT3



Conductivity at $T > 0$ determined by "scattering" of current by thermal stress-energy tensor.

R. C. Myers, S. Sachdev, and A. Singh, *Physical Review D* **83**, 066017 (2011)
D. Chowdhury, S. Raju, S. Sachdev, A. Singh, and P. Strack, *Physical Review B* **87**, 085138 (2013).
D. M. Hofman and J. Maldacena, JHEP **0805** (2008) 012.

AdS$_4$ theory of quantum criticality

Most general effective holographic theory for linear charge transport with 4 spatial derivatives:

$$\boxed{\begin{aligned}\mathcal{S}_{\rm bulk} &= \frac{1}{g_M^2}\int d^4x\,\sqrt{g}\left[\frac{1}{4}F_{ab}F^{ab} + \gamma L^2 C_{abcd}F^{ab}F^{cd}\right] \\ &+ \int d^4x\sqrt{g}\left[-\frac{1}{2\kappa^2}\left(R + \frac{6}{L^2}\right)\right],\end{aligned}}$$

This action is characterized by 3 dimensionless parameters, which can be linked to data of the CFT (OPE coefficients): 2-point correlators of the conserved current J_μ and the stress energy tensor $T_{\mu\nu}$, and a 3-point T, J, J correlator. Constraints from both the CFT and the gravitational theory bound $|\gamma| \leq 1/12 = 0.0833$..

R. C. Myers, S. Sachdev, and A. Singh, *Phys. Rev. D* **83**, 066017 (2011)
D. Chowdhury, S. Raju, S. Sachdev, A. Singh, and P. Strack, *Phys. Rev. B* **87**, 085138 (2013)

AdS$_4$ theory of quantum criticality

- The AdS$_4$ solutions satisfy two sum rules which are expected to be satisfied by all CFT3s:

$$
\begin{aligned}
\int_0^\infty d\omega \left(\mathrm{Re}\left[\Sigma(\omega)\right] - \Sigma(\infty)\right) &= 0 \\
\int_0^\infty d\omega \left(\mathrm{Re}\left[\frac{1}{\Sigma(\omega)}\right] - \frac{1}{\Sigma(\infty)}\right) &= 0
\end{aligned}
$$

The second sum rule relies on the existence of a S-dual CFT3.

W. Witczak-Krempa and S. Sachdev, *Phys. Rev. B* **86**, 235115 (2012)

AdS_4 theory of quantum criticality

- The AdS_4 solutions satisfy two sum rules which are expected to be satisfied by all CFT3s:

$$\int_0^\infty d\omega \left(\mathrm{Re}\left[\Sigma(\omega)\right] - \Sigma(\infty) \right) = 0$$

$$\int_0^\infty d\omega \left(\mathrm{Re}\left[\frac{1}{\Sigma(\omega)}\right] - \frac{1}{\Sigma(\infty)} \right) = 0$$

  The second sum rule relies on the existence of a S-dual CFT3.

- The poles *and zeros* of the complex function $\Sigma(\omega)$ are expected to be in the lower-half plane.

W. Witczak-Krempa and S. Sachdev, *Phys. Rev. B* **86**, 235115 (2012)

AdS$_4$ theory of quantum criticality

- The AdS$_4$ solutions satisfy two sum rules which are expected to be satisfied by all CFT3s:

$$\begin{aligned}\int_0^\infty d\omega \left(\mathrm{Re}\left[\Sigma(\omega)\right] - \Sigma(\infty)\right) &= 0 \\ \int_0^\infty d\omega \left(\mathrm{Re}\left[\frac{1}{\Sigma(\omega)}\right] - \frac{1}{\Sigma(\infty)}\right) &= 0\end{aligned}$$

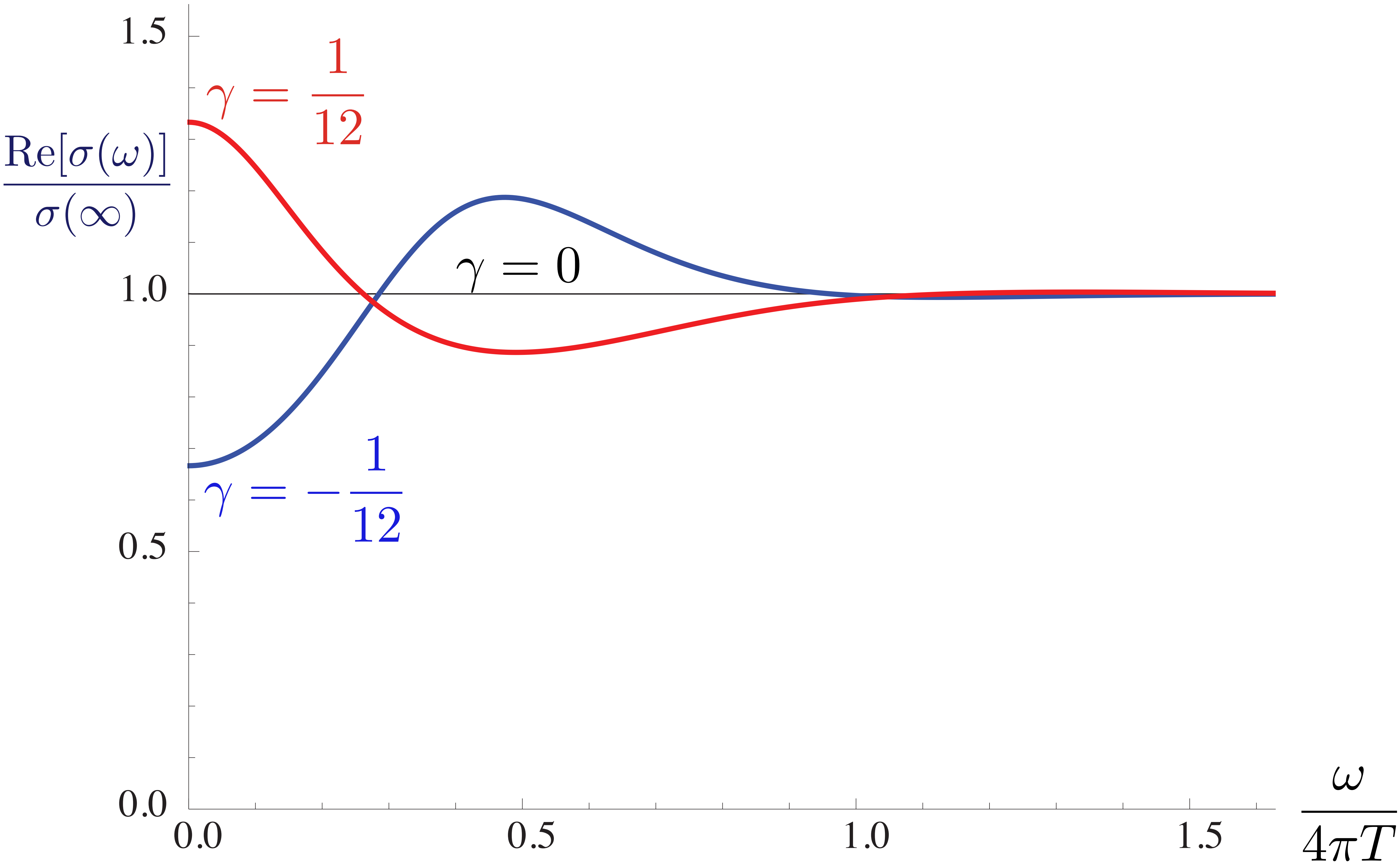
  The second sum rule relies on the existence of a S-dual CFT3.

- The poles *and zeros* of the complex function $\Sigma(\omega)$ are expected to be in the lower-half plane.

- The Boltzmann theory does *not* obey the above exact properties.

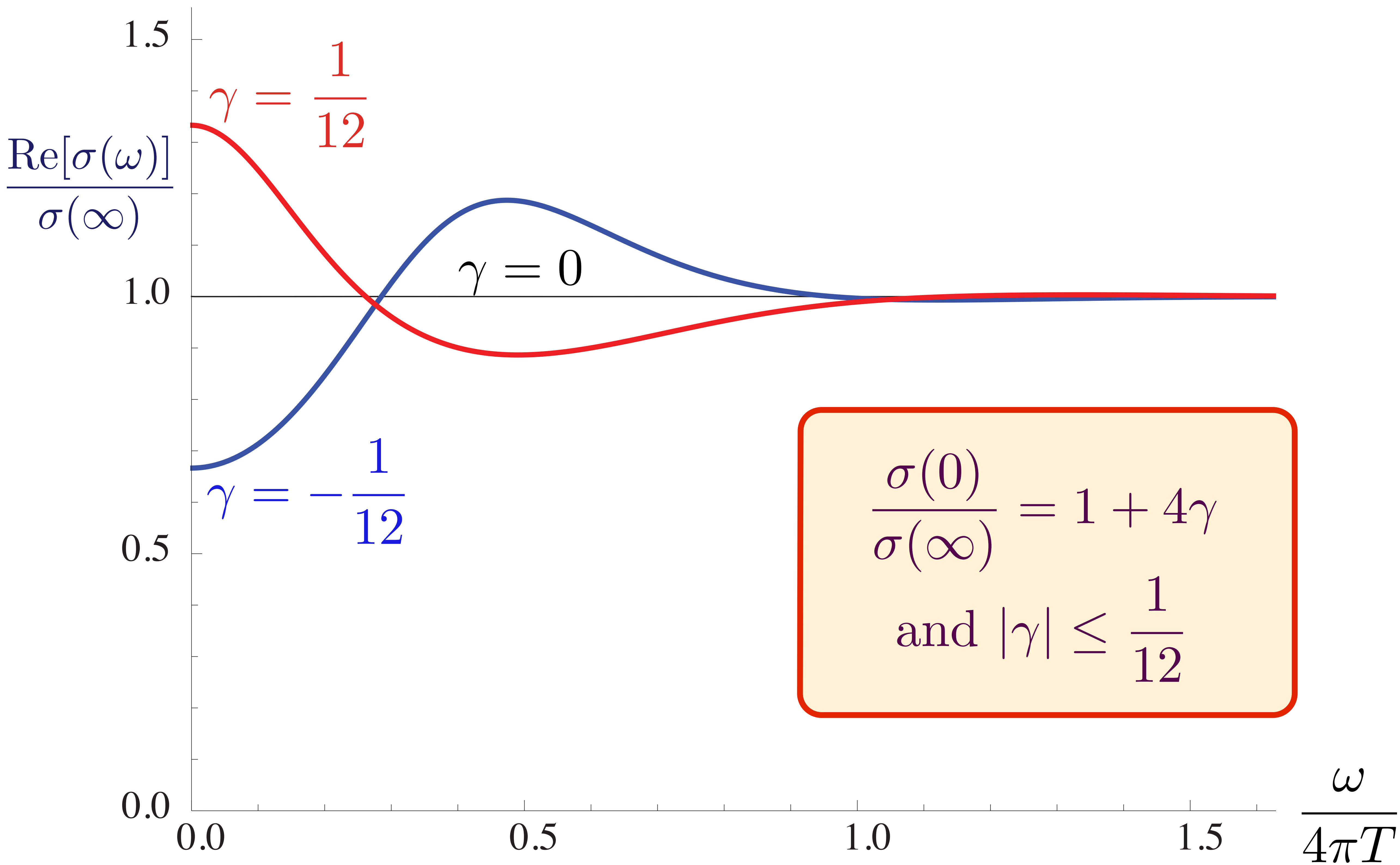
W. Witczak-Krempa and S. Sachdev, *Phys. Rev. B* **86**, 235115 (2012)

AdS$_4$ theory of quantum criticality



R. C. Myers, S. Sachdev, and A. Singh, *Physical Review D* **83**, 066017 (2011)

AdS$_4$ theory of quantum criticality

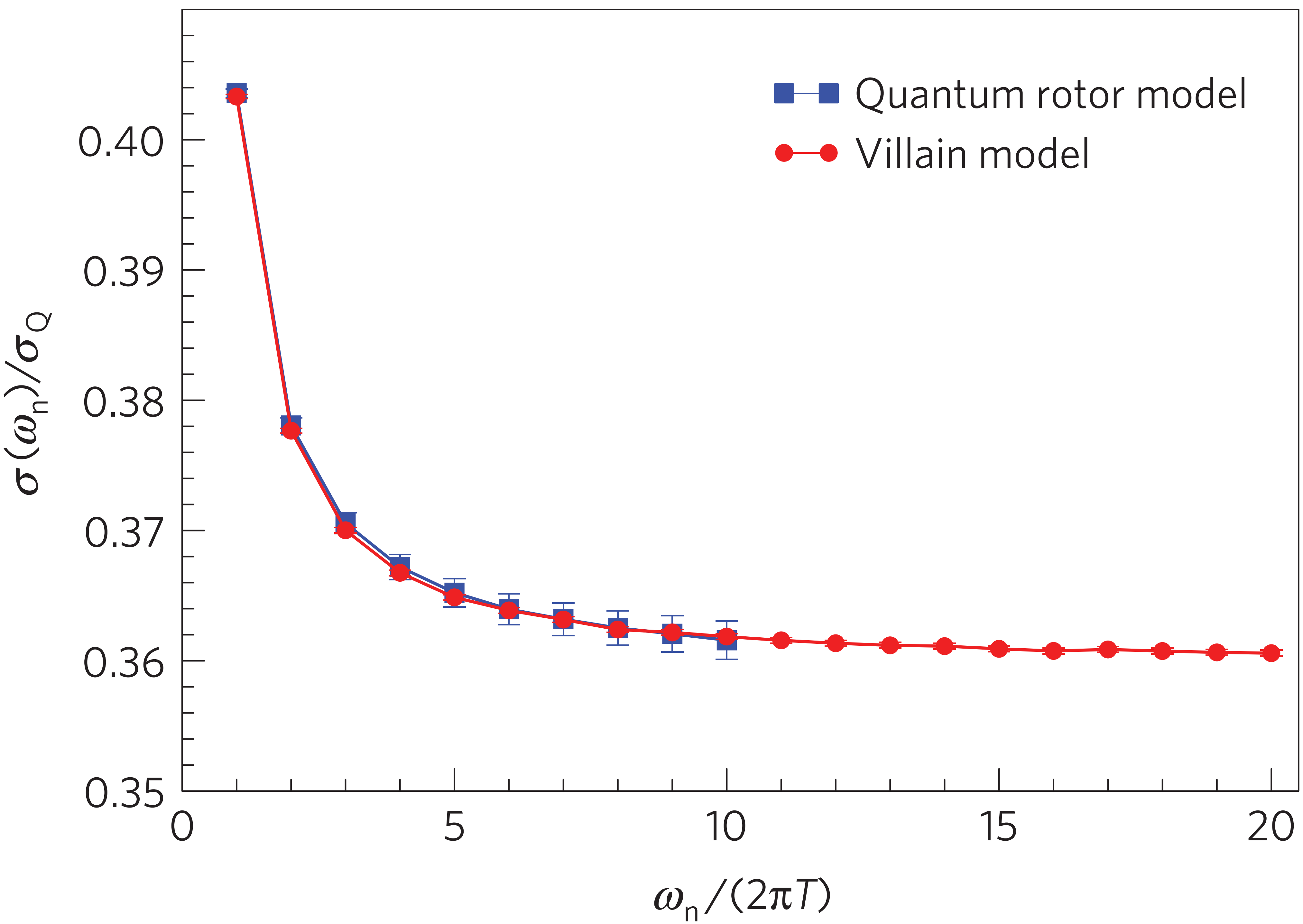


$$\frac{\sigma(0)}{\sigma(\infty)} = 1 + 4\gamma$$

and $|\gamma| \leq \frac{1}{12}$

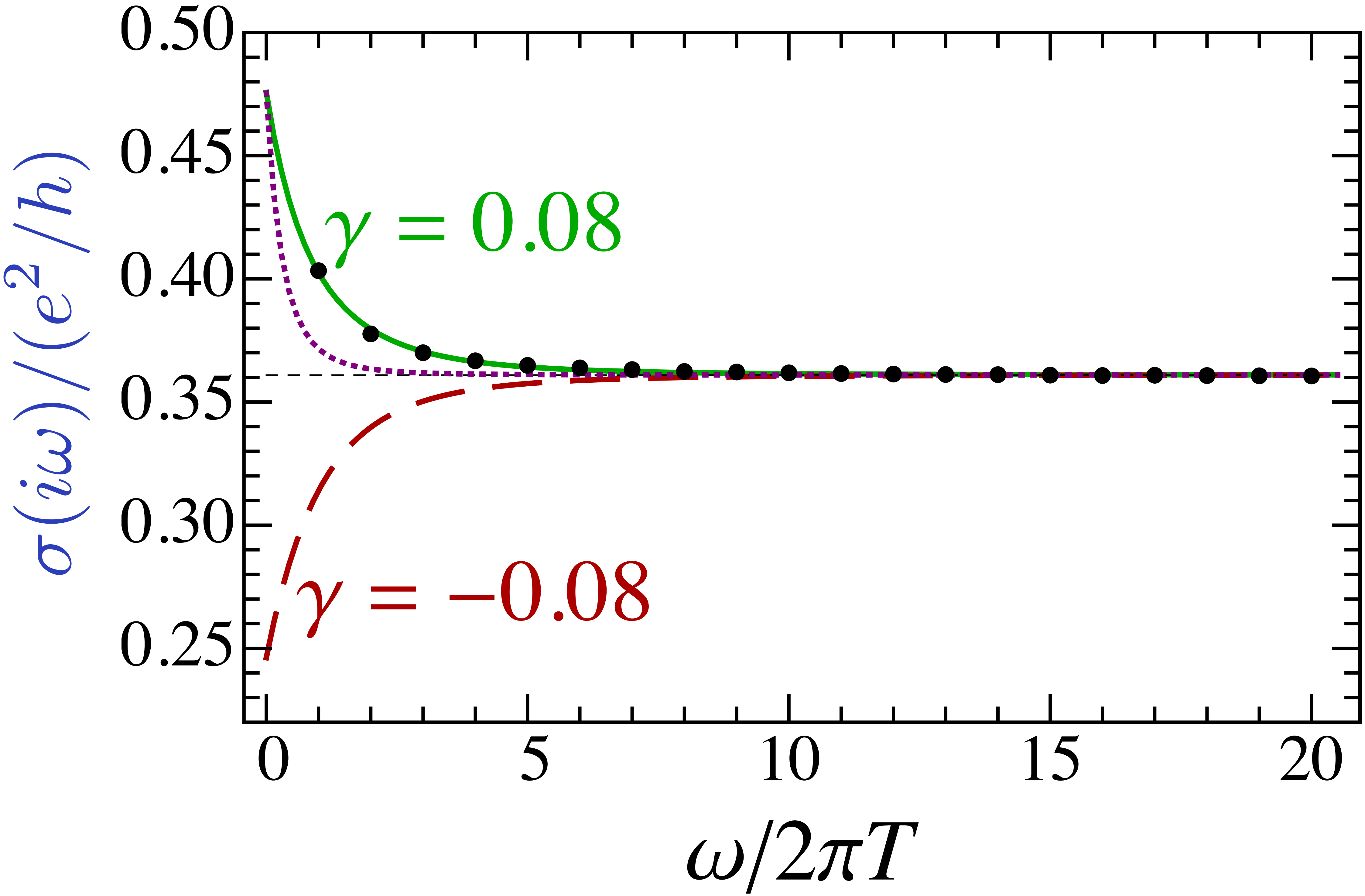
R. C. Myers, S. Sachdev, and A. Singh, *Physical Review D* **83**, 066017 (2011)

Quantum Monte Carlo for lattice bosons



W. Witczak-Krempa, E. Sorensen, and S. Sachdev, Nature Physics **10**, 361 (2014)
See also K. Chen, L. Liu, Y. Deng, L. Pollet, and N. Prokof'ev, Phys. Rev. Lett. **112**, 030402 (2013)

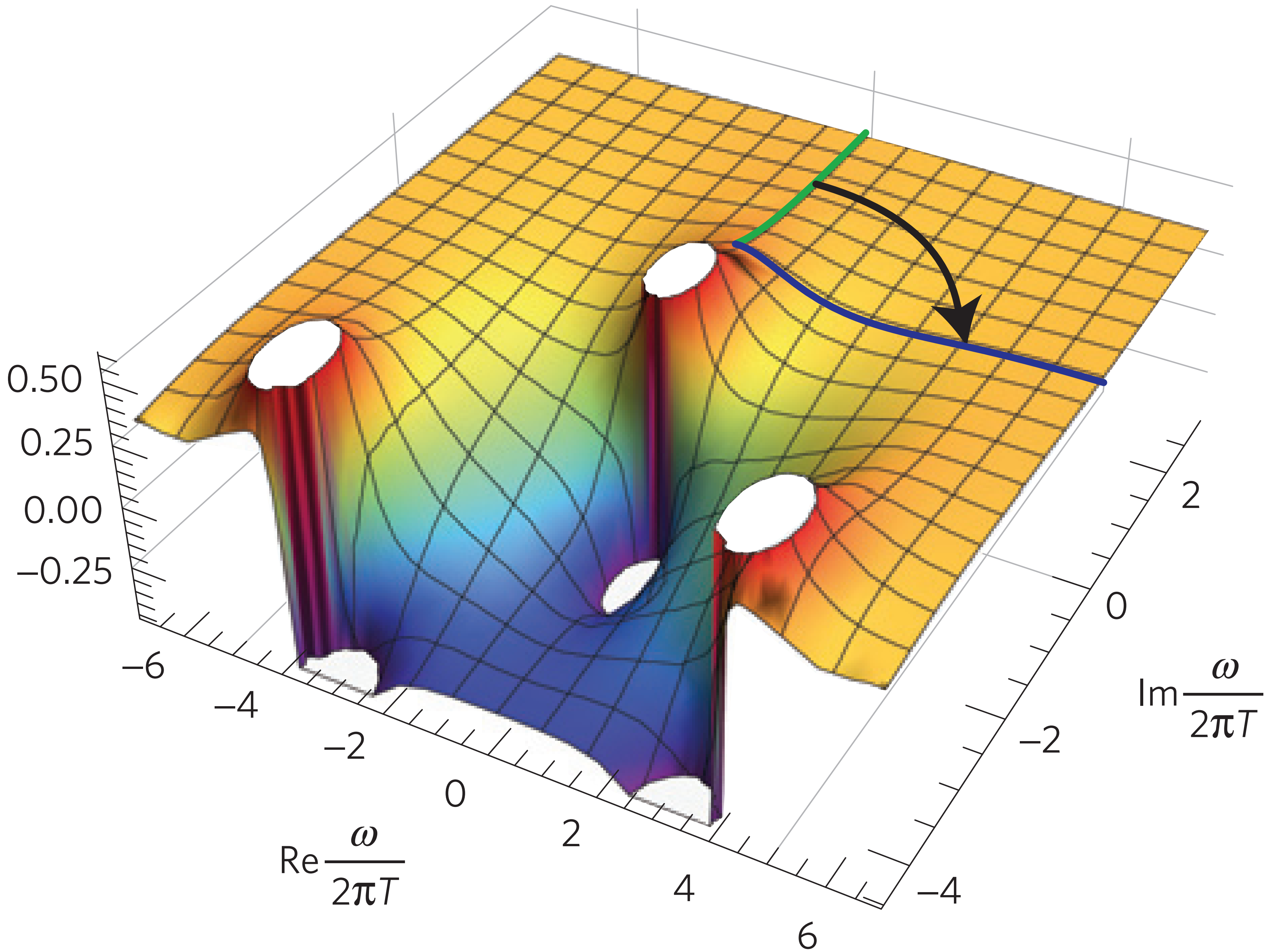
Transport revealed by QMC and holography



Fit of holographic theory to Monte Carlo,
after rescaling frequency axis

W. Witczak-Krempa, E. Sorensen, and S. Sachdev, Nature Physics **10**, 361 (2014)

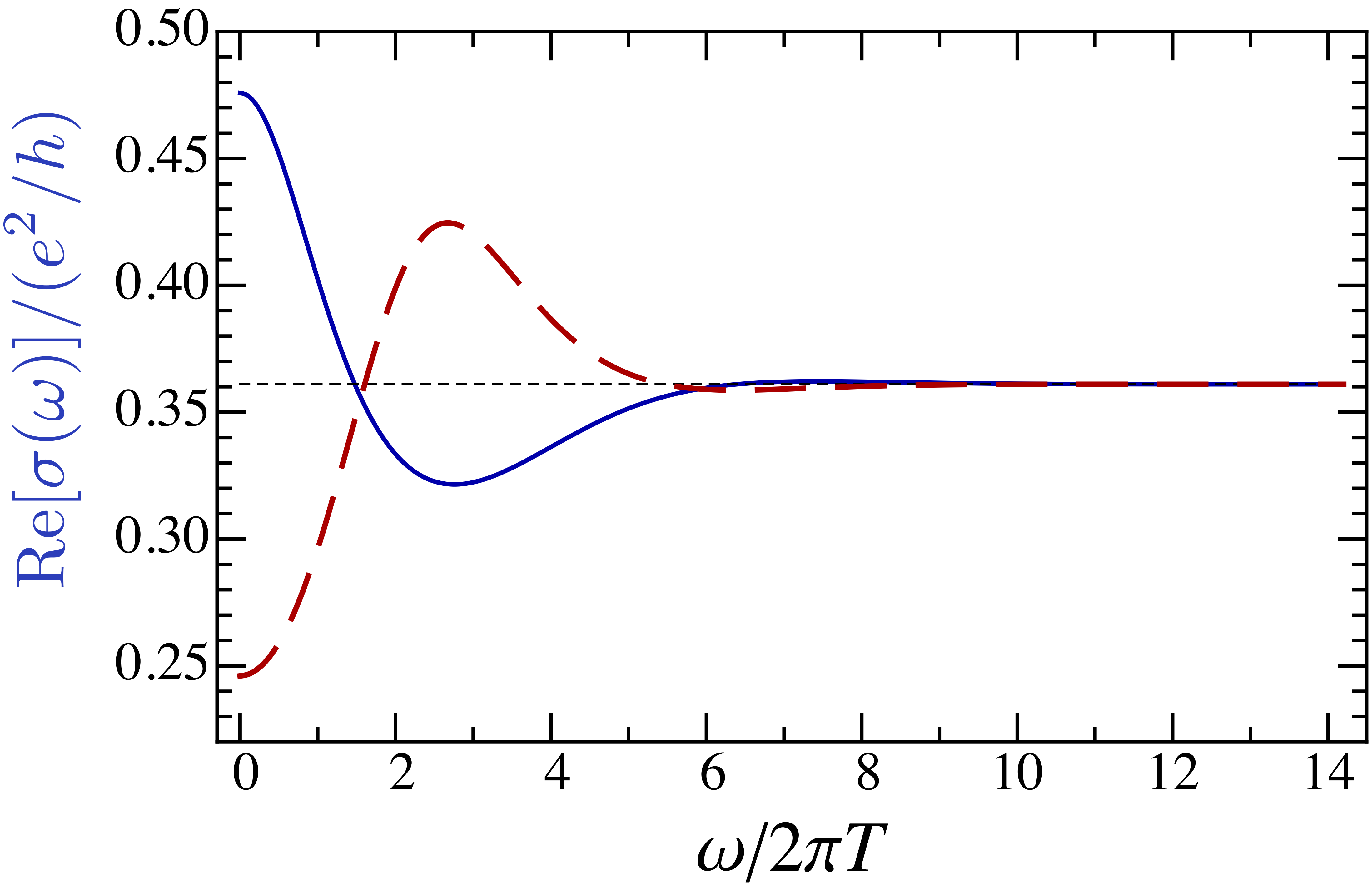
Analytic continuation by a holographic model



Obeys 2 non-trivial sum rules, and all poles and zeros are in the lower-half plane

W. Witczak-Krempa, E. Sorensen, and S. Sachdev, Nature Physics **10**, 361 (2014)

Transport revealed by QMC and holography



Predictions of holographic theory,
after analytic continuation to real frequencies

W. Witczak-Krempa, E. Sorensen, and S. Sachdev, Nature Physics **10**, 361 (2014)

Outline

1. Conformal field theories in 2+1 dimensions

   Superfluid-insulator transition

   of ultracold bosonic atoms in an optical lattice

2. Theory of a non-Fermi liquid

   Non-quasiparticle transport at the

   Ising-nematic quantum critical point

Outline

1. Conformal field theories in 2+1 dimensions

Superfluid-insulator transition

of ultracold bosonic atoms in an optical lattice

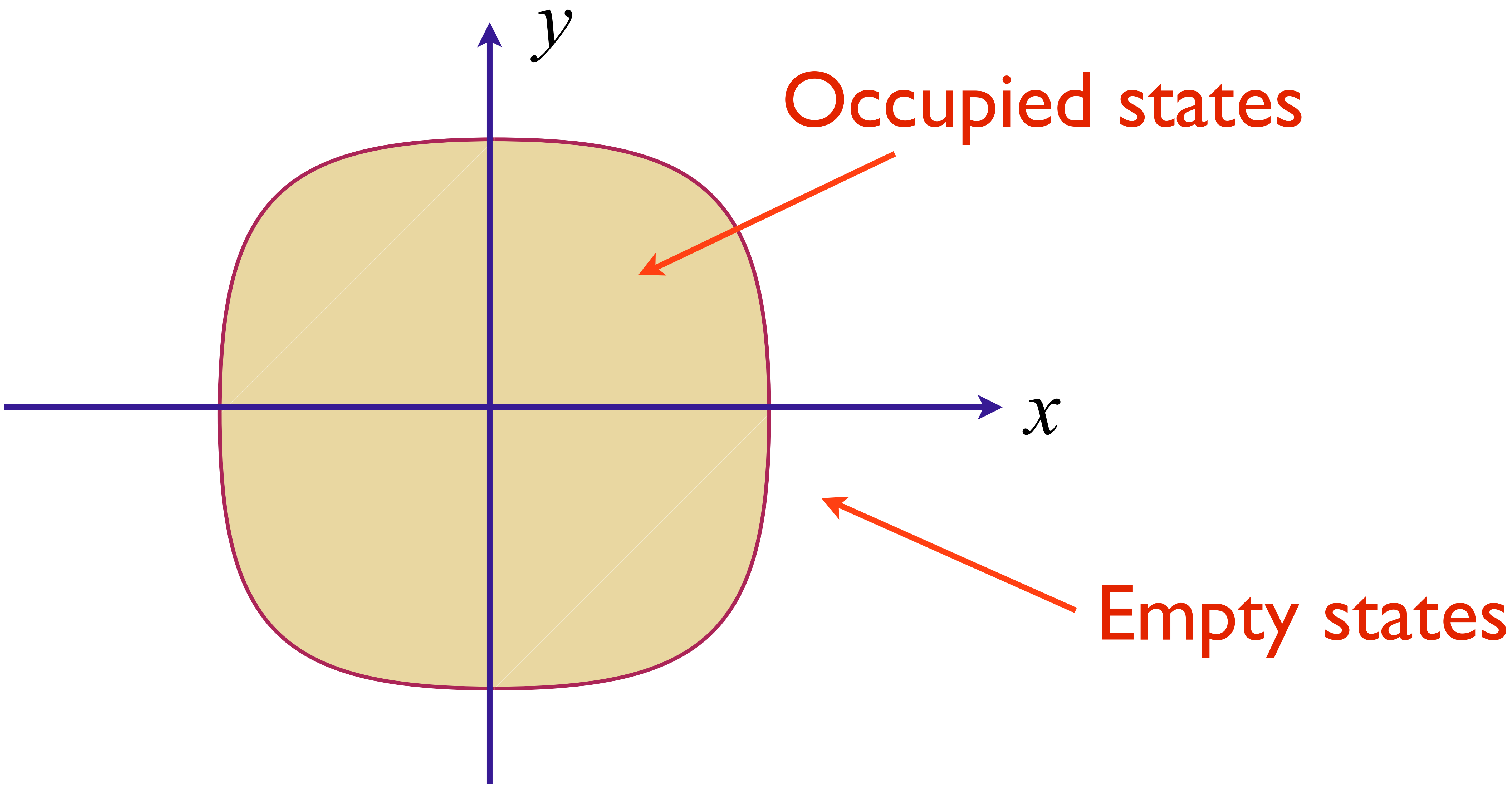
2. Theory of a non-Fermi liquid

Non-quasiparticle transport at the

Ising-nematic quantum critical point

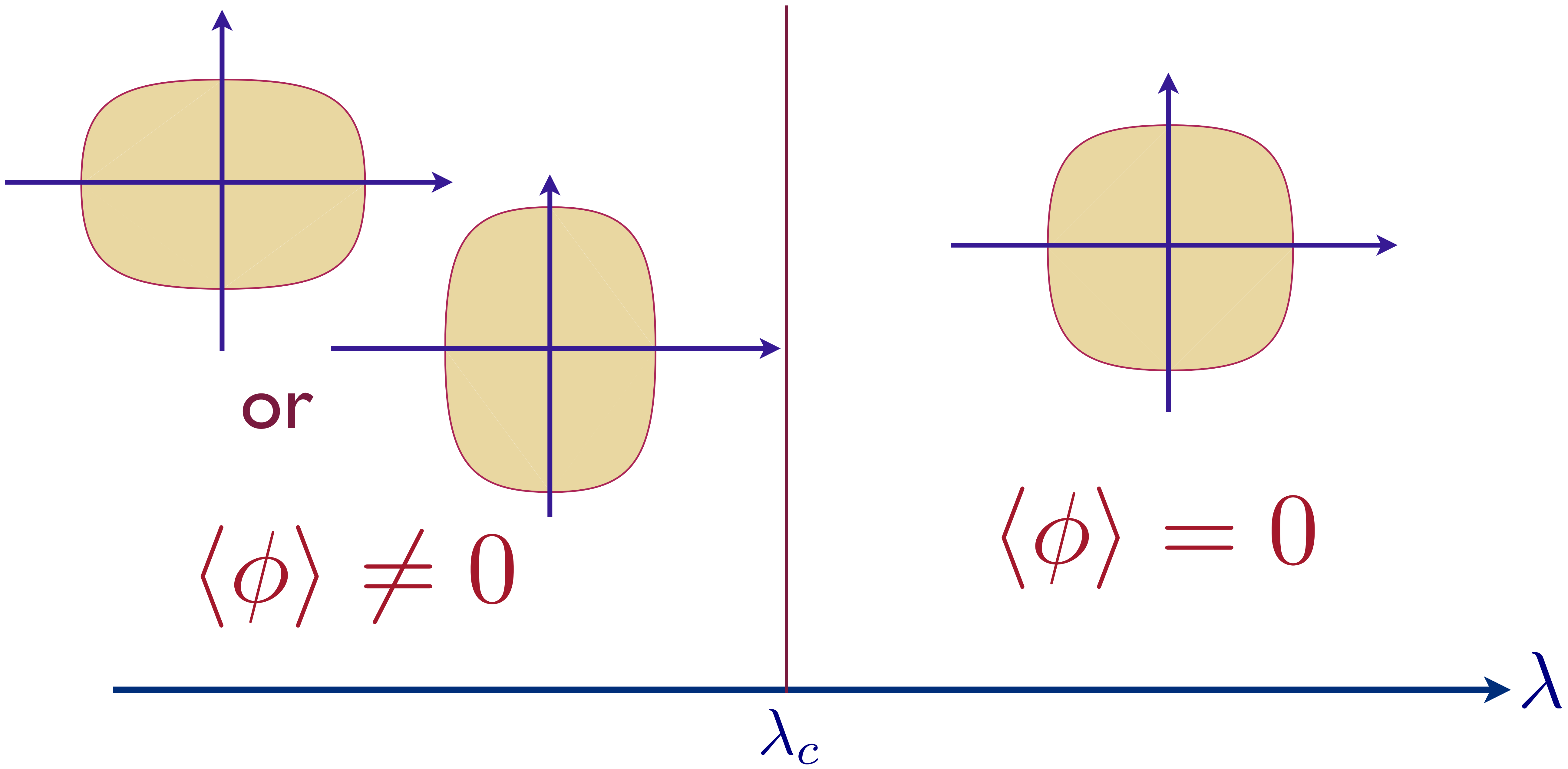
Max Metlitski

Quantum criticality of Ising-nematic ordering in a metal



A metal with a Fermi surface
with full square lattice symmetry

Quantum criticality of Ising-nematic ordering in a metal



Pomeranchuk instability as a function of coupling λ

Quantum criticality of Ising-nematic ordering in a metal

The "standard model":

$$\mathcal{S}_\phi = \int d^2 r d\tau \left[(\partial_\tau \phi)^2 + c^2 (\nabla \phi)^2 + (\lambda - \lambda_c) \phi^2 + u \phi^4 \right]$$

$$\mathcal{S}_c = \sum_{\alpha=1}^{N_f} \sum_{\mathbf{k}} \int d\tau c_{\mathbf{k}\alpha}^\dagger \left(\partial_\tau + \varepsilon_{\mathbf{k}} \right) c_{\mathbf{k}\alpha}$$

$$\mathcal{S}_{\phi c} = -g \int d\tau \sum_{\alpha=1}^{N_f} \sum_{\mathbf{k},\mathbf{q}} \phi_{\mathbf{q}} \left(\cos k_x - \cos k_y \right) c_{\mathbf{k}+\mathbf{q}/2,\alpha}^\dagger c_{\mathbf{k}-\mathbf{q}/2,\alpha}$$

Quantum criticality of Ising-nematic ordering in a metal

The "standard model":

$$\mathcal{S}_\phi = \int d^2r d\tau \left[(\partial_\tau \phi)^2 + c^2 (\nabla \phi)^2 + (\lambda - \lambda_c)\phi^2 + u\phi^4 \right]$$

$$\mathcal{S}_c = \sum_{\alpha=1}^{N_f} \sum_{\mathbf{k}} \int d\tau c^\dagger_{\mathbf{k}\alpha} \left(\partial_\tau + \varepsilon_{\mathbf{k}} \right) c_{\mathbf{k}\alpha}$$

$$\mathcal{S}_{\phi c} = -g \int d\tau \sum_{\alpha=1}^{N_f} \sum_{\mathbf{k},\mathbf{q}} \phi_{\mathbf{q}} \left(\cos k_x - \cos k \right.$$

Field theory of bosonic order parameter

Quantum criticality of Ising-nematic ordering in a metal

The "standard model":

$$\mathcal{S}_\phi = \int d^2r d\tau \left[(\partial_\tau \phi)^2 + c^2 (\nabla \phi)^2 + (\lambda - \lambda_c)\phi^2 + u\phi^4 \right]$$

$$\mathcal{S}_c = \sum_{\alpha=1}^{N_f} \sum_{\mathbf{k}} \int d\tau c^\dagger_{\mathbf{k}\alpha} \left(\partial_\tau + \varepsilon_{\mathbf{k}} \right) c_{\mathbf{k}\alpha}$$

Electrons with a Fermi surface: $\varepsilon_{\mathbf{k}} = -2t(\cos k_x + \cos k_y) - \mu \ldots$

$$\mathcal{S}_{\phi c} = -g \int d\tau \sum_{\alpha=1}^{N_f} \sum_{\mathbf{k},\mathbf{q}} \phi_{\mathbf{q}} \left(\cos k_x - \cos k_y \right) c^\dagger_{\mathbf{k}+\mathbf{q}/2,\alpha} c_{\mathbf{k}-\mathbf{q}/2,\alpha}$$

Quantum criticality of Ising-nematic ordering in a metal

The "standard model":

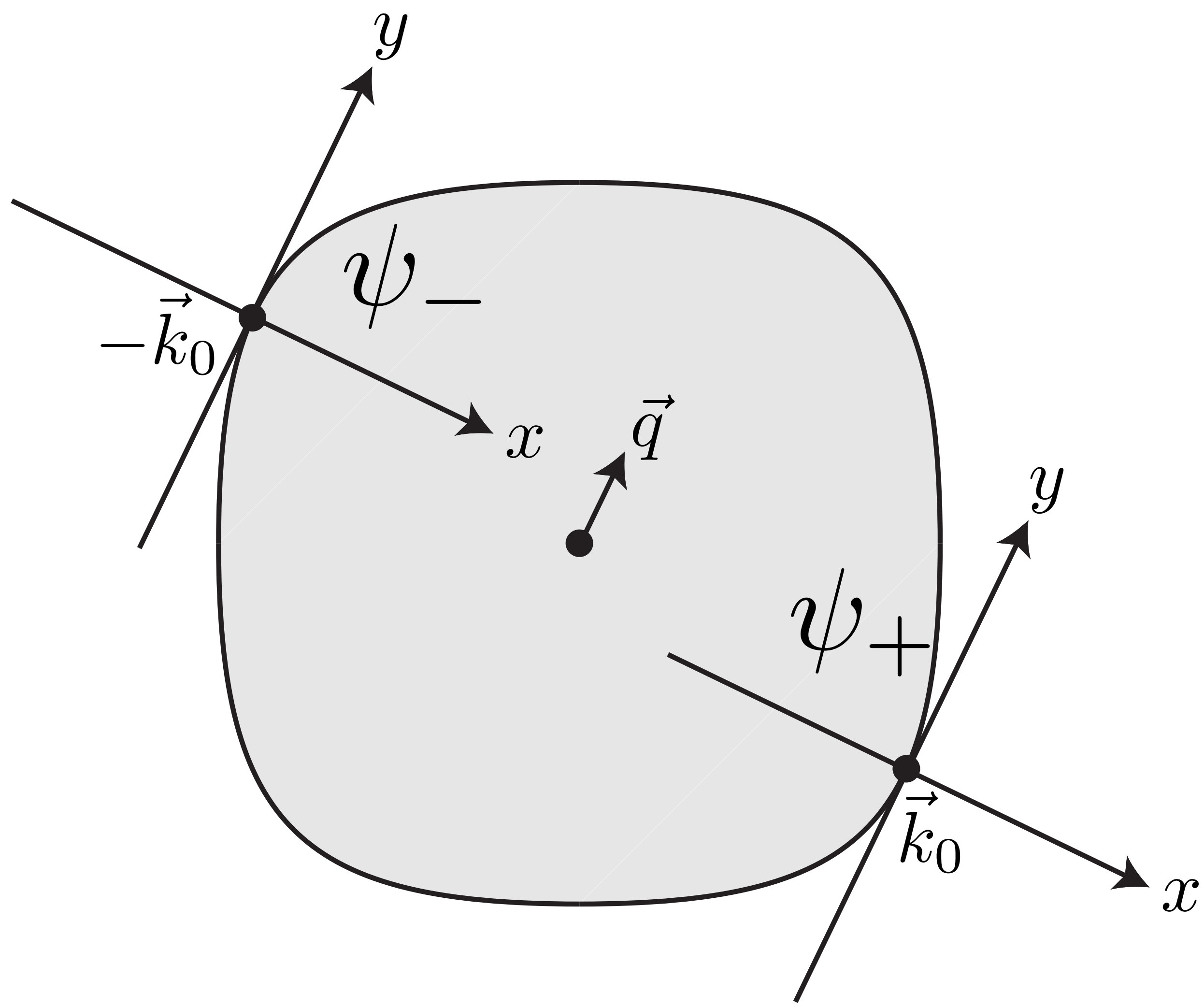
$$\mathcal{S}_\phi = \int d^2 r d\tau \left[(\partial_\tau \phi)^2 + c^2 (\nabla \phi)^2 + (\lambda - \right.$$

$$\mathcal{S}_c = \sum_{\alpha=1}^{N_f} \sum_{\mathbf{k}} \int d\tau c_{\mathbf{k}\alpha}^\dagger \left(\partial_\tau + \varepsilon_{\mathbf{k}} \right) c_{\mathbf{k}\alpha}$$

$$\mathcal{S}_{\phi c} = -g \int d\tau \sum_{\alpha=1}^{N_f} \sum_{\mathbf{k},\mathbf{q}} \phi_{\mathbf{q}} \left(\cos k_x - \cos k_y \right) c_{\mathbf{k}+\mathbf{q}/2,\alpha}^\dagger c_{\mathbf{k}-\mathbf{q}/2,\alpha}$$

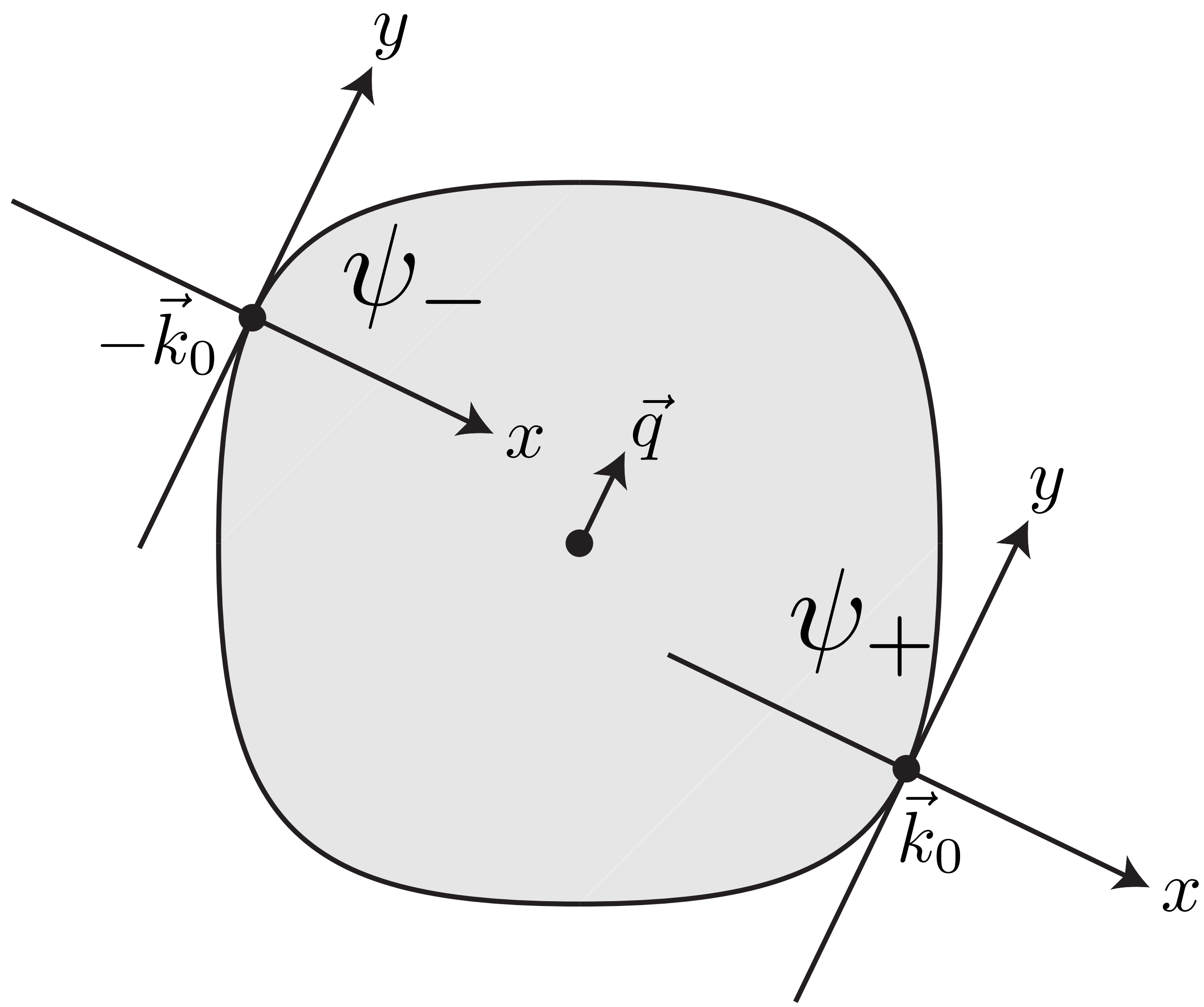
"Yukawa" coupling between bosons and fermions

Quantum criticality of Ising-nematic ordering in a metal



- ϕ fluctuation at wavevector $\vec{q}$ couples most efficiently to fermions near $\pm\vec{k}_0$.

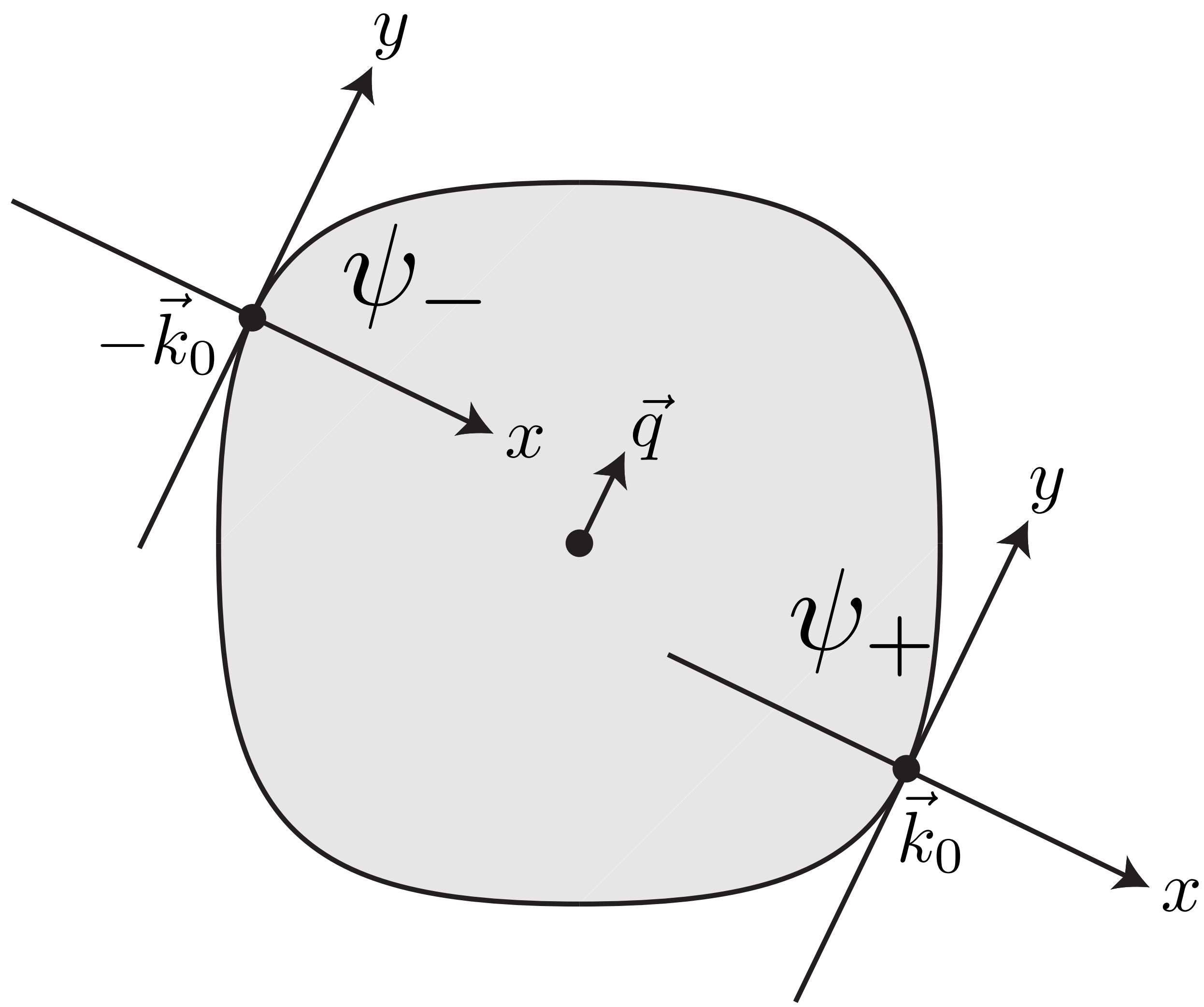
Quantum criticality of Ising-nematic ordering in a metal



- ϕ fluctuation at wavevector $\vec{q}$ couples most efficiently to fermions near $\pm\vec{k}_0$.

- Expand fermion kinetic energy at wavevectors about $\pm\vec{k}_0$ and boson (ϕ) kinetic energy about $\vec{q} = 0$.

Quantum criticality of Ising-nematic ordering in a metal



$$\mathcal{L}[\psi_\pm, \phi] =$$

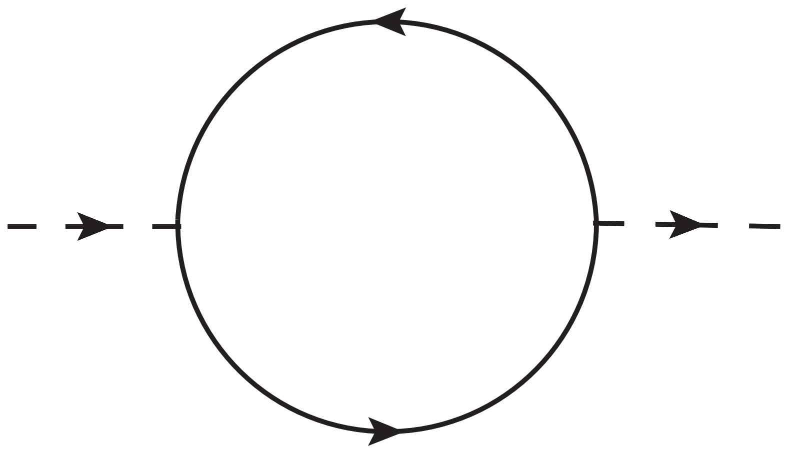
$$\psi_+^\dagger \left(\partial_\tau - i\partial_x - \partial_y^2\right) \psi_+ + \psi_-^\dagger \left(\partial_\tau + i\partial_x - \partial_y^2\right) \psi_-$$

$$-\phi \left(\psi_+^\dagger \psi_+ + \psi_-^\dagger \psi_-\right) + \frac{1}{2g^2} \left(\partial_y \phi\right)^2$$

M. A. Metlitski and S. Sachdev, Phys. Rev. B **82**, 075127 (2010)

Quantum criticality of Ising-nematic ordering in a metal

$$
\begin{aligned}
\mathcal{L} &= \psi_+^\dagger \left(\partial_\tau - i\partial_x - \partial_y^2\right)\psi_+ + \psi_-^\dagger \left(\partial_\tau + i\partial_x - \partial_y^2\right)\psi_- \\
&\quad - \phi\left(\psi_+^\dagger \psi_+ + \psi_-^\dagger \psi_-\right) + \frac{1}{2g^2}\left(\partial_y \phi\right)^2
\end{aligned}
$$



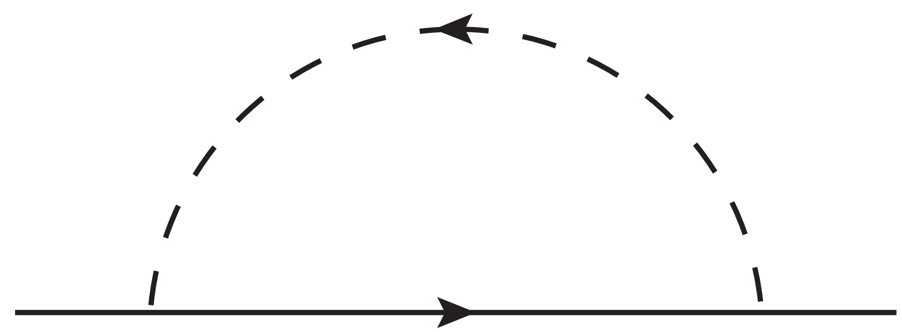
One loop ϕ self-energy with N_f fermion flavors:

$$
\begin{aligned}
\Sigma_\phi(\vec{q},\omega) &= N_f \int \frac{d^2k}{4\pi^2}\frac{d\Omega}{2\pi}\frac{1}{\left[-i(\Omega+\omega)+k_x+q_x+(k_y+q_y)^2\right]\left[-i\Omega-k_x+k_y^2\right]} \\
&= \frac{N_f}{4\pi}\frac{|\omega|}{|q_y|}
\end{aligned}
$$

Landau-damping

Quantum criticality of Ising-nematic ordering in a metal

$$\mathcal{L} = \psi_+^\dagger \left(\partial_\tau - i\partial_x - \partial_y^2\right) \psi_+ + \psi_-^\dagger \left(\partial_\tau + i\partial_x - \partial_y^2\right) \psi_- - \phi \left(\psi_+^\dagger \psi_+ + \psi_-^\dagger \psi_-\right) + \frac{1}{2g^2} \left(\partial_y \phi\right)^2$$



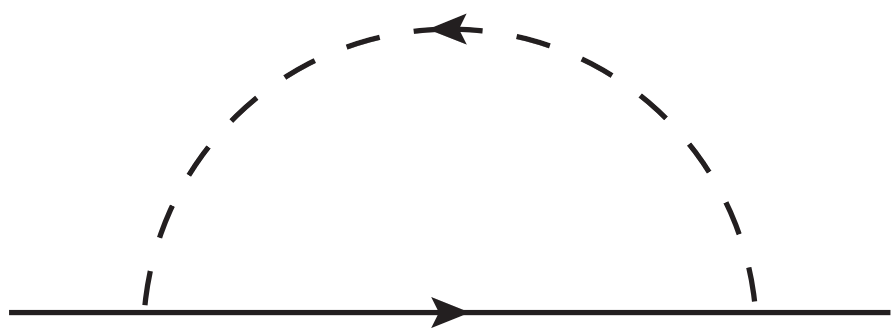
Electron self-energy at order $1/N_f$:

$$\Sigma(\vec{k}, \Omega) = -\frac{1}{N_f} \int \frac{d^2 q}{4\pi^2} \frac{d\omega}{2\pi} \frac{1}{\left[-i(\omega + \Omega) + k_x + q_x + (k_y + q_y)^2\right] \left[\frac{q_y^2}{g^2} + \frac{|\omega|}{|q_y|}\right]}$$

$$= -i \frac{2}{\sqrt{3} N_f} \left(\frac{g^2}{4\pi}\right)^{2/3} \mathrm{sgn}(\Omega) |\Omega|^{2/3}$$

Quantum criticality of Ising-nematic ordering in a metal

$$\mathcal{L} = \psi_+^\dagger \left(\partial_\tau - i\partial_x - \partial_y^2\right)\psi_+ + \psi_-^\dagger \left(\partial_\tau + i\partial_x - \partial_y^2\right)\psi_- - \phi\left(\psi_+^\dagger\psi_+ + \psi_-^\dagger\psi_-\right) + \frac{1}{2g^2}\left(\partial_y\phi\right)^2$$



Electron self-energy at order $1/N_f$:

$$\Sigma(\vec{k},\Omega) = -\frac{1}{N_f}\int\frac{d^2q}{4\pi^2}\frac{d\omega}{2\pi}\frac{1}{\left[-i(\omega+\Omega)+k_x+q_x+(k_y+q_y)^2\right]\left[\frac{q_y^2}{g^2}+\frac{|\omega|}{|q_y|}\right]}$$

$$= -i\frac{2}{\sqrt{3}N_f}\left(\frac{g^2}{4\pi}\right)^{2/3}\mathrm{sgn}(\Omega)|\Omega|^{2/3}$$

$\sim |\Omega|^{d/3}$ in dimension d.

Quantum criticality of Ising-nematic ordering in a metal

$$\mathcal{L} = \psi_+^\dagger \left(\partial_\tau - i\partial_x - \partial_y^2\right) \psi_+ + \psi_-^\dagger \left(\partial_\tau + i\partial_x - \partial_y^2\right) \psi_- - \phi \left(\psi_+^\dagger \psi_+ + \psi_-^\dagger \psi_-\right) + \frac{1}{2g^2} \left(\partial_y \phi\right)^2$$

Schematic form of ϕ and fermion Green's functions in d dimensions

$$D(\vec{q},\omega) = \frac{1/N_f}{q_\perp^2 + \dfrac{|\omega|}{|q_\perp|}} \quad , \quad G_f(\vec{q},\omega) = \frac{1}{q_x + q_\perp^2 - i\mathrm{sgn}(\omega)|\omega|^{d/3}/N_f}$$

In the boson case, $q_\perp^2 \sim \omega^{1/z_b}$ with $z_b = 3/2$.
In the fermion case, $q_x \sim q_\perp^2 \sim \omega^{1/z_f}$ with $z_f = 3/d$.

Note $z_f < z_b$ for $d > 2$ $\Rightarrow$ Fermions have *higher* energy than bosons, and perturbation theory in g is OK.
Strongly-coupled theory in $d = 2$.

Quantum criticality of Ising-nematic ordering in a metal

$$
\begin{aligned}
\mathcal{L} &= \psi_+^\dagger \left(\partial_\tau - i\partial_x - \partial_y^2\right) \psi_+ + \psi_-^\dagger \left(\partial_\tau + i\partial_x - \partial_y^2\right) \psi_- \\
&\quad - \phi \left(\psi_+^\dagger \psi_+ + \psi_-^\dagger \psi_-\right) + \frac{1}{2g^2} \left(\partial_y \phi\right)^2
\end{aligned}
$$

Schematic form of ϕ and fermion Green's functions in $d = 2$

$$
D(\vec{q}, \omega) = \frac{1/N_f}{q_y^2 + \dfrac{|\omega|}{|q_y|}} \quad , \quad G_f(\vec{q}, \omega) = \frac{1}{q_x + q_y^2 - i \operatorname{sgn}(\omega) |\omega|^{2/3} / N_f}
$$

In *both* cases $q_x \sim q_y^2 \sim \omega^{1/z}$, with $z = 3/2$. Note that the bare term $\sim \omega$ in G_f^{-1} is irrelevant.

Strongly-coupled theory without quasiparticles.

Quantum criticality of Ising-nematic ordering in a metal

$$\begin{aligned}\mathcal{L} &= \psi_+^\dagger \left(\partial_\tau - i\partial_x - \partial_y^2\right) \psi_+ + \psi_-^\dagger \left(\partial_\tau + i\partial_x - \partial_y^2\right) \psi_- \\ &\quad - \phi \left(\psi_+^\dagger \psi_+ + \psi_-^\dagger \psi_-\right) + \frac{1}{2g^2} \left(\partial_y \phi\right)^2\end{aligned}$$

Simple scaling argument for $z = 3/2$.

Quantum criticality of Ising-nematic ordering in a metal

$$\begin{aligned}\mathcal{L} &= \psi_+^\dagger \left(\cancel{\partial_\tau} - i\partial_x - \partial_y^2\right)\psi_+ + \psi_-^\dagger \left(\cancel{\partial_\tau} + i\partial_x - \partial_y^2\right)\psi_- \\ &\quad - \phi\left(\psi_+^\dagger \psi_+ + \psi_-^\dagger \psi_-\right) + \frac{1}{2g^2}\left(\partial_y \phi\right)^2\end{aligned}$$

Simple scaling argument for $z = 3/2$.

Quantum criticality of Ising-nematic ordering in a metal

$$
\begin{aligned}
\mathcal{L} &= \psi_+^\dagger \left(\cancel{\partial_\tau} - i\partial_x - \partial_y^2\right) \psi_+ + \psi_-^\dagger \left(\cancel{\partial_\tau} + i\partial_x - \partial_y^2\right) \psi_- \\
&\quad - \phi \left(\psi_+^\dagger \psi_+ + \psi_-^\dagger \psi_-\right) + \frac{1}{2g^2} \left(\partial_y \phi\right)^2
\end{aligned}
$$

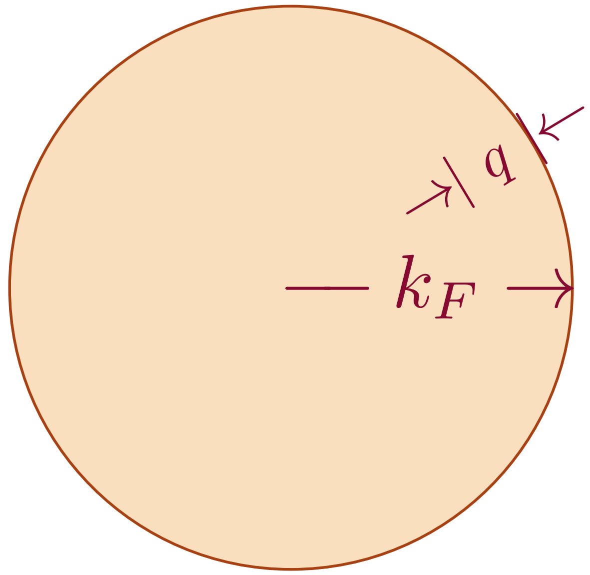
Simple scaling argument for $z = 3/2$.

Under the rescaling $x \to x/s$, $y \to y/s^{1/2}$, and $\tau \to \tau/s^z$, we find invariance provided

$$
\begin{aligned}
\phi &\to \phi\, s \\
\psi &\to \psi\, s^{(2z+1)/4} \\
g &\to g\, s^{(3-2z)/4}
\end{aligned}
$$

So the action is invariant provided $z = 3/2$.

<u>FL</u>
<u>Fermi liquid</u>



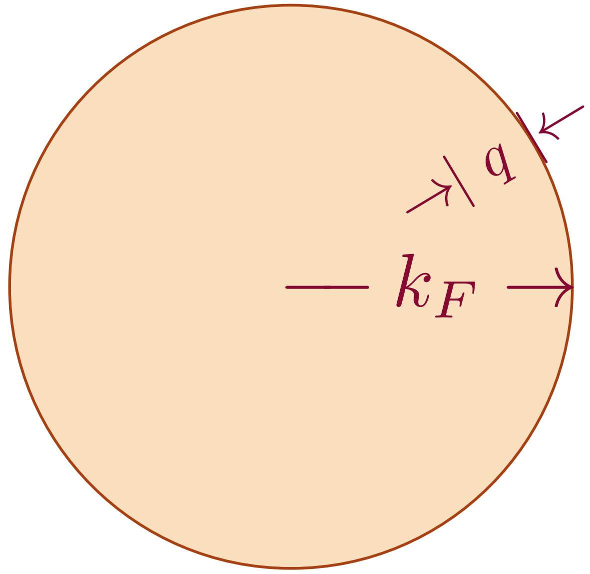
- $k_F^d \sim \mathcal{Q}$, the fermion density

- Sharp fermionic excitations near Fermi surface with $\omega \sim |q|^z$, and $z = 1$.

- Entropy density $S \sim T^{(d-\theta)/z}$ with violation of hyperscaling exponent $\theta = d - 1$.

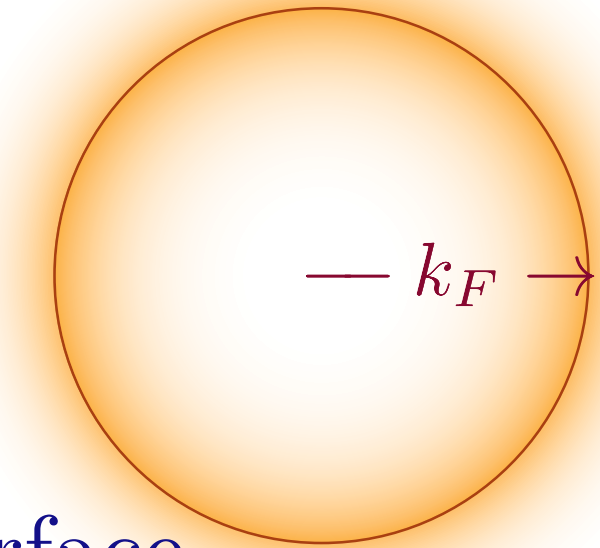
- Entanglement entropy $S_E \sim k_F^{d-1} P \ln P$.

FL Fermi liquid



- $k_F^d \sim \mathcal{Q}$, the fermion density

NFL Nematic QCP

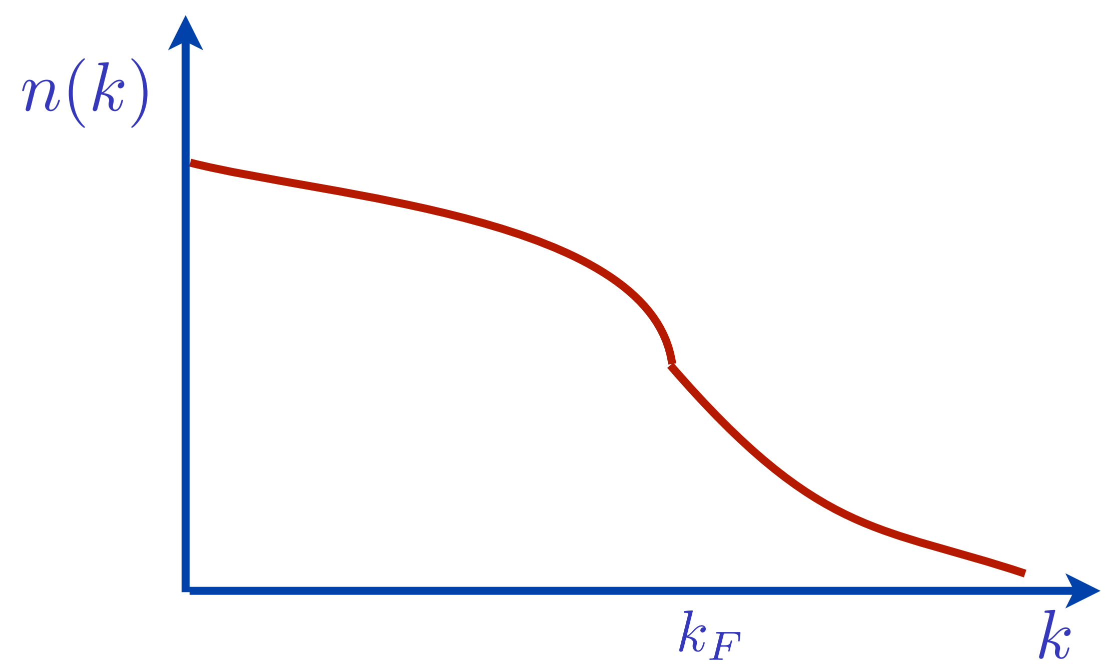


- Fermi surface with $k_F^d \sim \mathcal{Q}$.

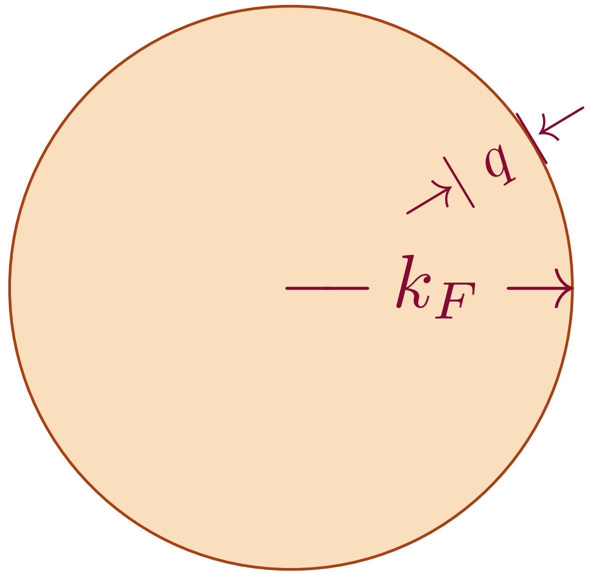
- Sharp fermionic excitations near Fermi surface with $\omega \sim |q|^z$, and $z = 1$.

- Entropy density $S \sim T^{(d-\theta)/z}$ with violation of hyperscaling exponent $\theta = d - 1$.

- Entanglement entropy $S_E \sim k_F^{d-1} P \ln P$.

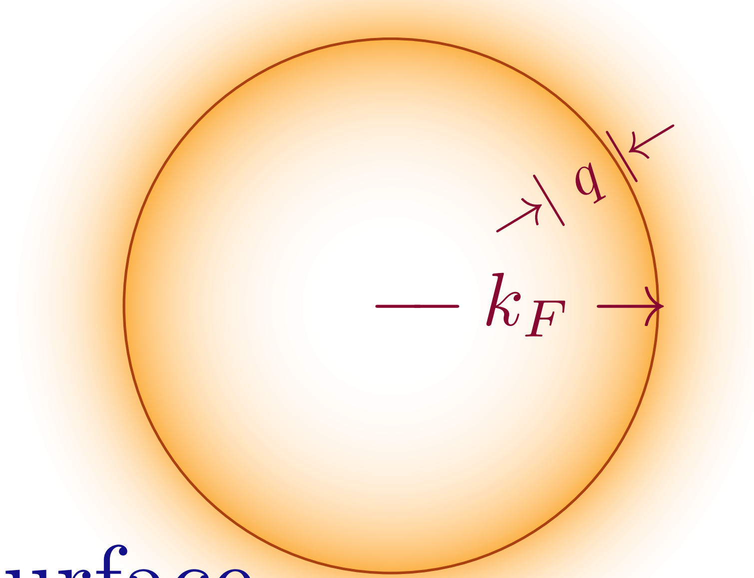


FL Fermi liquid



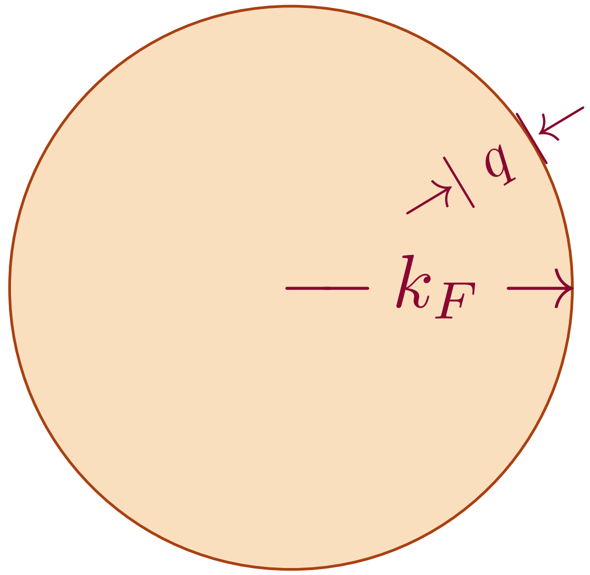
- $k_F^d \sim \mathcal{Q}$, the fermion density
- Sharp fermionic excitations near Fermi surface with $\omega \sim |q|^z$, and $z = 1$.
- Entropy density $S \sim T^{(d-\theta)/z}$ with violation of hyperscaling exponent $\theta = d - 1$.
- Entanglement entropy $S_E \sim k_F^{d-1} P \ln P$.

NFL Nematic QCP



- Fermi surface with $k_F^d \sim \mathcal{Q}$.
- Diffuse fermionic excitations with $z = 3/2$ to three loops.

FL Fermi liquid



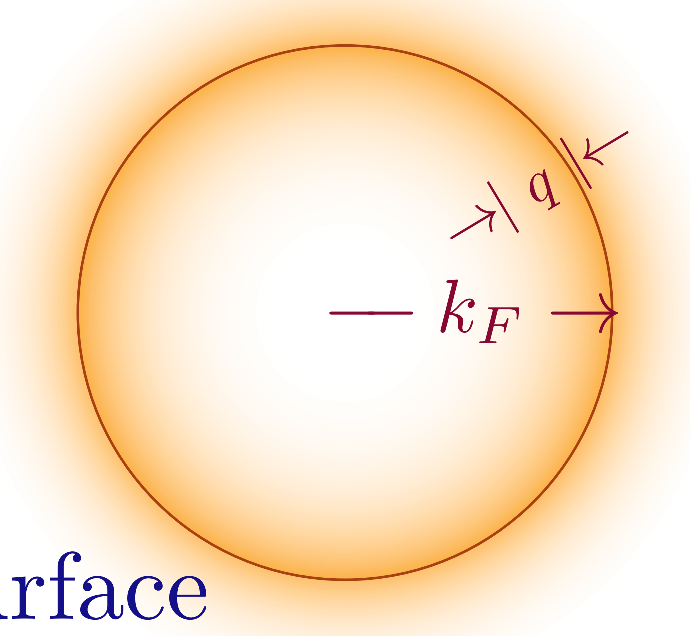
- $k_F^d \sim \mathcal{Q}$, the fermion density

- Sharp fermionic excitations near Fermi surface with $\omega \sim |q|^z$, and $z = 1$.

- Entropy density $S \sim T^{(d-\theta)/z}$ with violation of hyperscaling exponent $\theta = d - 1$.

- Entanglement entropy $S_E \sim k_F^{d-1} P \ln P$.

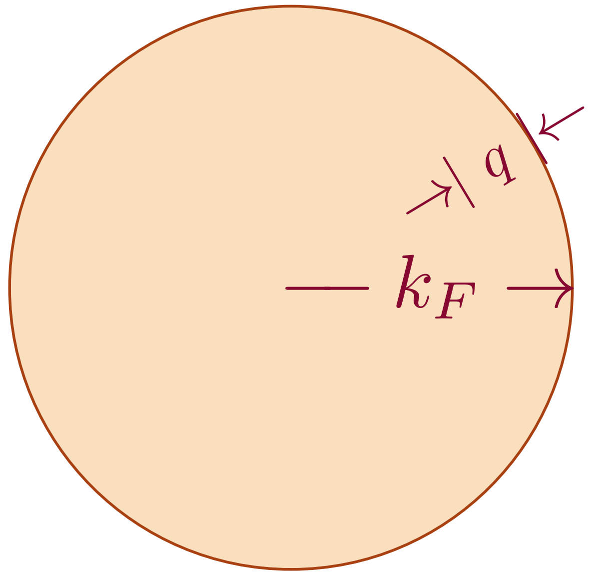
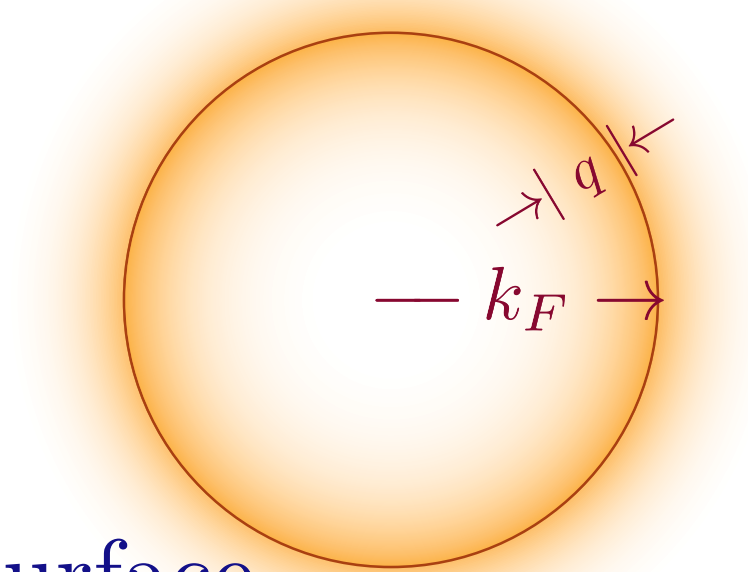
NFL Nematic QCP



- Fermi surface with $k_F^d \sim \mathcal{Q}$.

- Diffuse fermionic excitations with $z = 3/2$ to three loops.

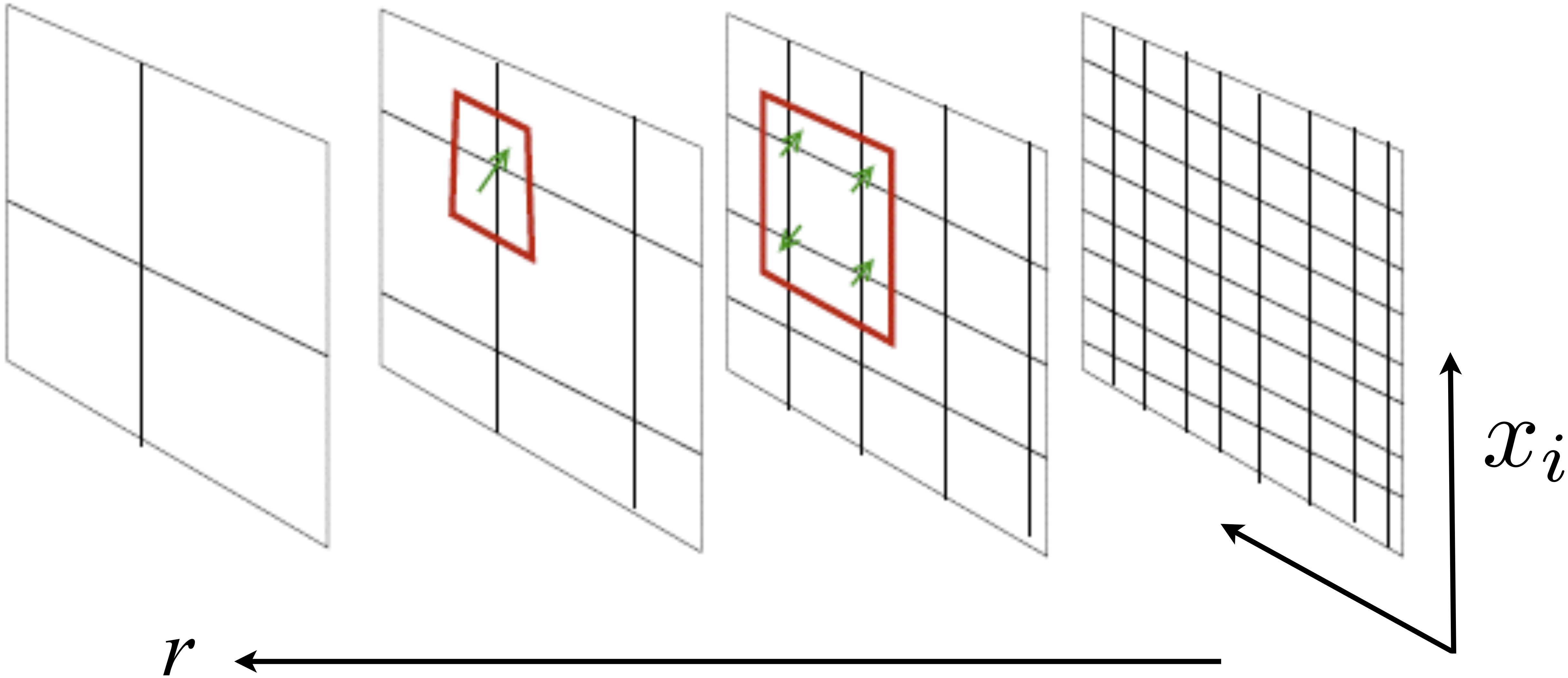
- $S \sim T^{(d-\theta)/z}$ with $\theta = d - 1$.

FL Fermi liquid	**NFL Nematic QCP**
	
• $k_F^d \sim \mathcal{Q}$, the fermion density	• Fermi surface with $k_F^d \sim \mathcal{Q}$.
• Sharp fermionic excitations near Fermi surface with $\omega \sim \|q\|^z$, and $z = 1$.	• Diffuse fermionic excitations with $z = 3/2$ to three loops.
• Entropy density $S \sim T^{(d-\theta)/z}$ with violation of hyperscaling exponent $\theta = d - 1$.	• $S \sim T^{(d-\theta)/z}$ with $\theta = d - 1$.
• Entanglement entropy $S_E \sim k_F^{d-1} P \ln P$.	• $S_E \sim k_F^{d-1} P \ln P$.

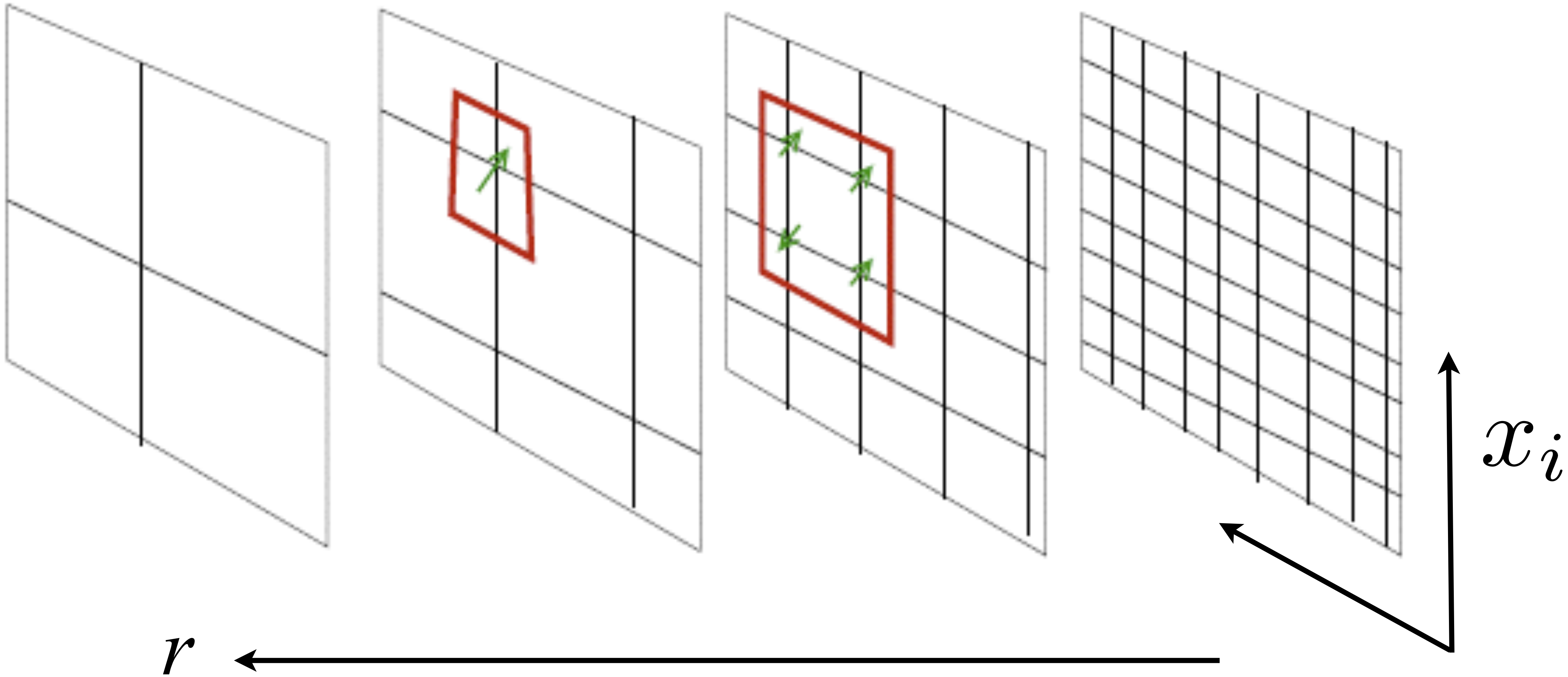
Liza Huijse

Brian Swingle

Generalized holography



Generalized holography

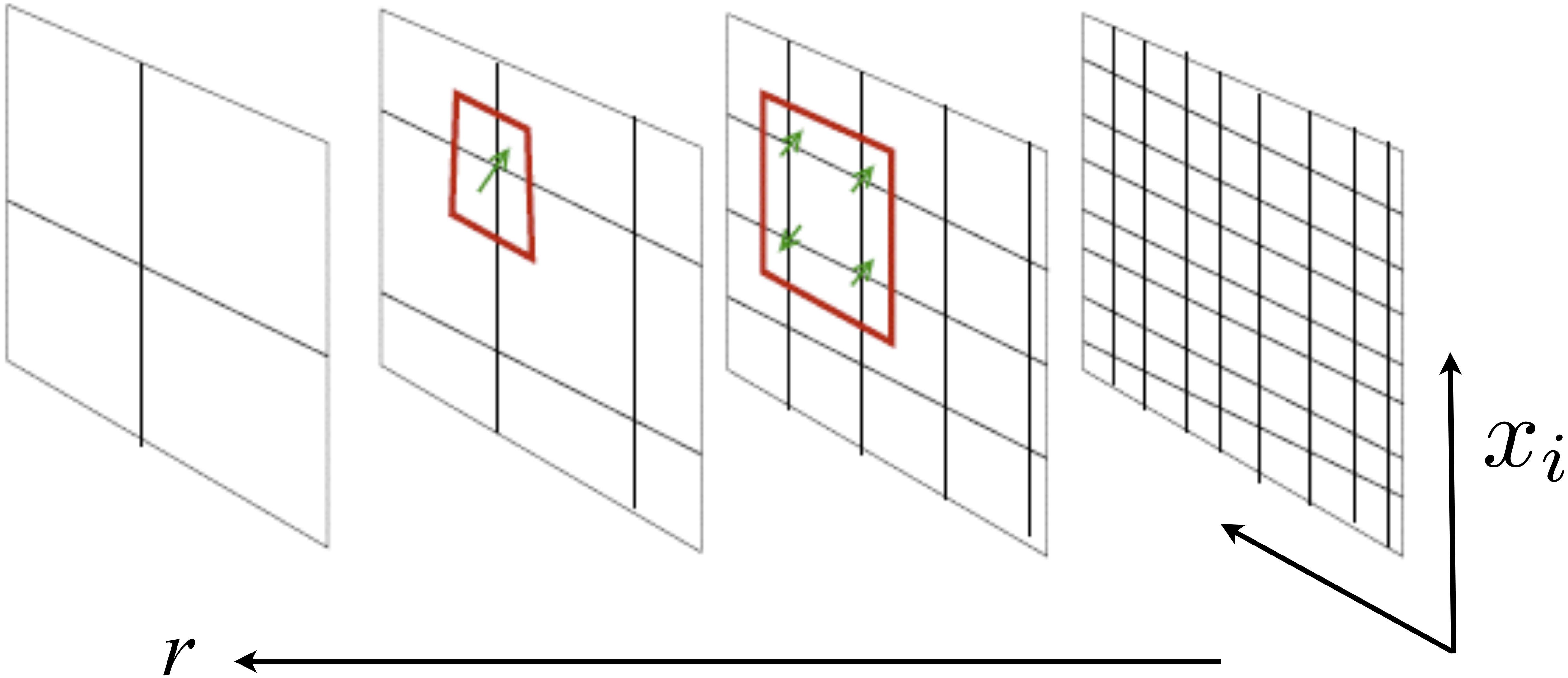


Consider a metric which transforms under rescaling as

$$x_i \to \zeta\, x_i, \quad t \to \zeta^z\, t, \quad ds \to \zeta^{\theta/d}\, ds.$$

Recall: conformal matter has $\theta = 0$, $z = 1$, and the metric is anti-de Sitter

Generalized holography



The most general such metric is

$$ds^2 = \frac{1}{r^2}\left(-\frac{dt^2}{r^{2d(z-1)/(d-\theta)}} + r^{2\theta/(d-\theta)}dr^2 + dx_i^2\right)$$

Generalized holography

$$ds^2 = \frac{1}{r^2}\left(-\frac{dt^2}{r^{2d(z-1)/(d-\theta)}} + r^{2\theta/(d-\theta)}dr^2 + dx_i^2\right)$$

This is the most general metric which is invariant under the scale transformation

$$\begin{aligned} x_i &\rightarrow \zeta\, x_i \\ t &\rightarrow \zeta^z\, t \\ ds &\rightarrow \zeta^{\theta/d}\, ds. \end{aligned}$$

This identifies z as the dynamic critical exponent ($z = 1$ for "relativistic" quantum critical points). We will see shortly that θ is the violation of hyperscaling exponent.
We have used reparametrization invariance in r to define it so that it scales as

$$r \rightarrow \zeta^{(d-\theta)/d} r\,.$$

L. Huijse, S. Sachdev, B. Swingle, Physical Review B **85**, 035121 (2012)

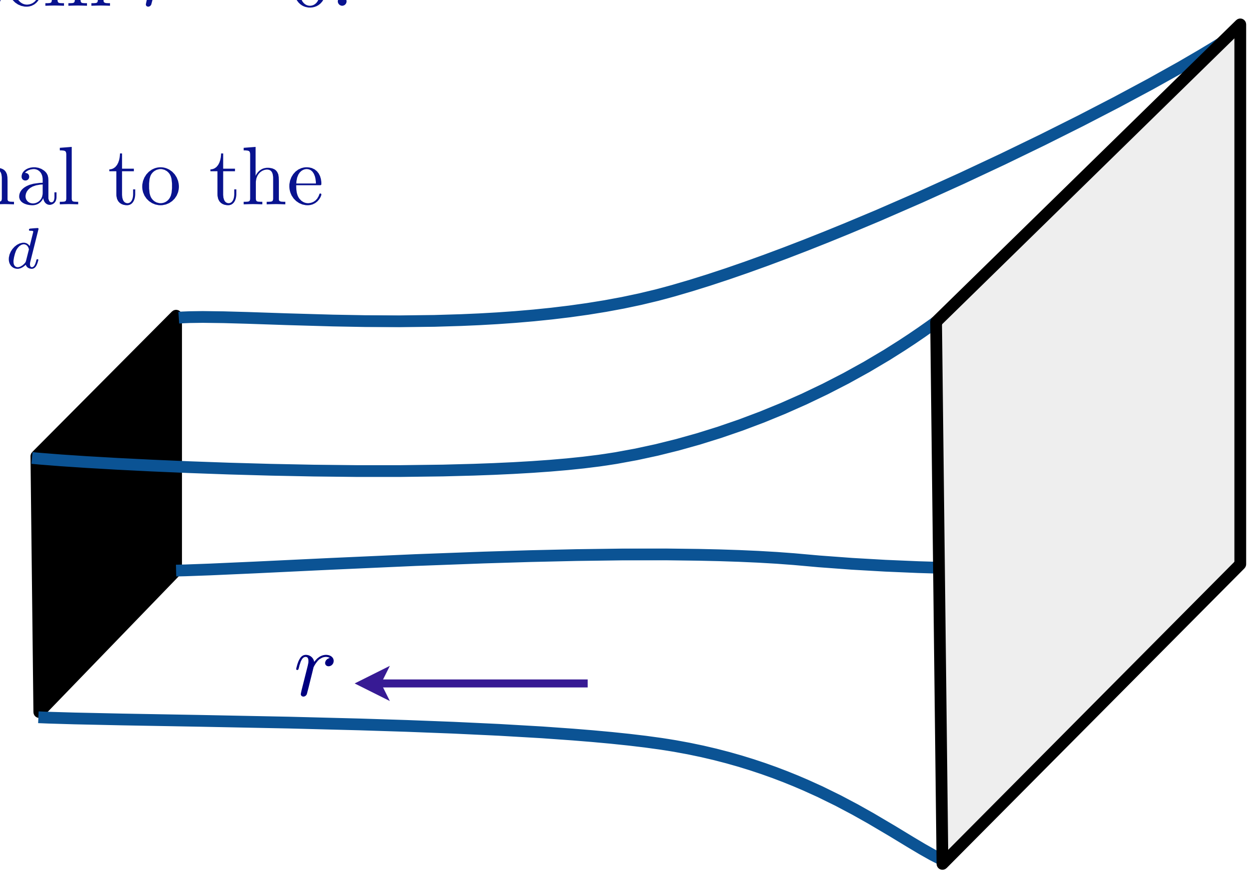
Generalized holography

$$ds^2 = \frac{1}{r^2}\left(-\frac{dt^2}{r^{2d(z-1)/(d-\theta)}} + r^{2\theta/(d-\theta)}dr^2 + dx_i^2\right)$$

At $T > 0$, there is a "black-brane" at $r = r_h$.

The Beckenstein-Hawking entropy of the black-brane is the thermal entropy of the quantum system $r = 0$.

The entropy density, S, is proportional to the "area" of the horizon, and so $S \sim r_h^{-d}$



L. Huijse, S. Sachdev, B. Swingle, Physical Review B **85**, 035121 (2012)

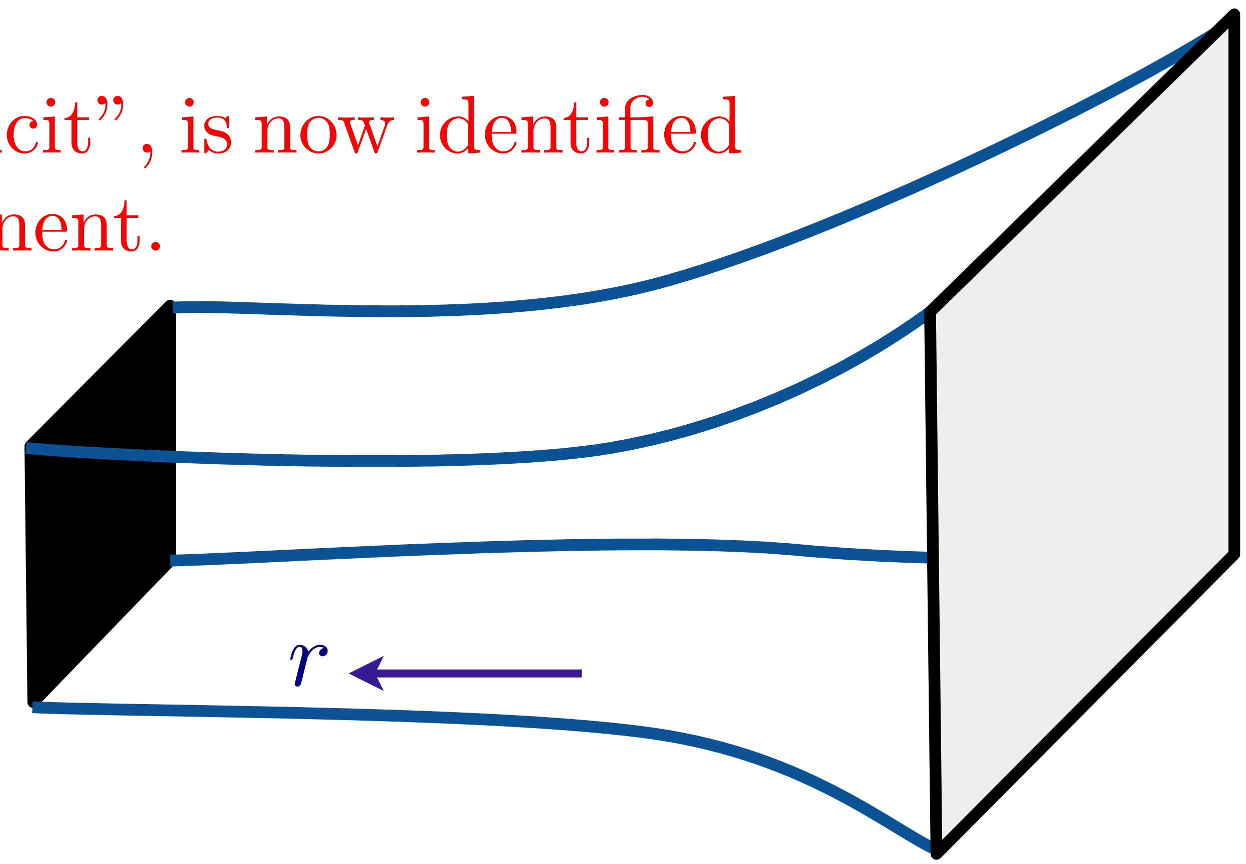
Generalized holography

$$\boxed{ds^2 = \frac{1}{r^2}\left(-\frac{dt^2}{r^{2d(z-1)/(d-\theta)}} + r^{2\theta/(d-\theta)}dr^2 + dx_i^2\right)}$$

Under rescaling $r \to \zeta^{(d-\theta)/d} r$, and the temperature $T \sim t^{-1}$, and so

$$S \sim T^{(d-\theta)/z} = T^{d_{\rm eff}/z}$$

where $\theta = d - d_{\rm eff}$, the "dimension deficit", is now identified as the violation of hyperscaling exponent.



L. Huijse, S. Sachdev, B. Swingle, Physical Review B **85**, 035121 (2012)

Generalized holography

$$ds^2 = \frac{1}{r^2}\left(-\frac{dt^2}{r^{2d(z-1)/(d-\theta)}} + r^{2\theta/(d-\theta)}dr^2 + dx_i^2\right)$$

The null energy condition (stability condition for gravity) yields a new inequality

$$z \geq 1 + \frac{\theta}{d}$$

N. Ogawa, T. Takayanagi, and T. Ugajin, JHEP **1201**, 125 (2012).
L. Huijse, S. Sachdev, B. Swingle, Physical Review B **85**, 035121 (2012)

Generalized holography

$$ds^2 = \frac{1}{r^2}\left(-\frac{dt^2}{r^{2d(z-1)/(d-\theta)}} + r^{2\theta/(d-\theta)}dr^2 + dx_i^2\right)$$

The null energy condition (stability condition for gravity) yields a new inequality

$$z \geq 1 + \frac{\theta}{d}$$

The Fermi liquid has $\theta = d - 1$ and $z = 1$: so the Fermi liquid does *not* have such a gravity dual.

N. Ogawa, T. Takayanagi, and T. Ugajin, JHEP **1201**, 125 (2012).
L. Huijse, S. Sachdev, B. Swingle, Physical Review B **85**, 035121 (2012)

Generalized holography

$$ds^2 = \frac{1}{r^2}\left(-\frac{dt^2}{r^{2d(z-1)/(d-\theta)}} + r^{2\theta/(d-\theta)}dr^2 + dx_i^2\right)$$

The null energy condition (stability condition for gravity) yields a new inequality

$$z \geq 1 + \frac{\theta}{d}$$

The Fermi liquid has $\theta = d - 1$ and $z = 1$: so the Fermi liquid does *not* have such a gravity dual.
The non-Fermi liquid in $d = 2$ has $\theta = d - 1$, and this implies $z \geq 3/2$. So the lower bound is precisely the value obtained for the non-Fermi liquid!

N. Ogawa, T. Takayanagi, and T. Ugajin, JHEP **1201**, 125 (2012).
L. Huijse, S. Sachdev, B. Swingle, Physical Review B **85**, 035121 (2012)

Generalized holography

$$ds^2 = \frac{1}{r^2}\left(-\frac{dt^2}{r^{2d(z-1)/(d-\theta)}} + r^{2\theta/(d-\theta)} dr^2 + dx_i^2 \right)$$

Application of the Ryu-Takayanagi minimal area formula to a dual Einstein-Maxwell-dilaton theory yields

$$S_E \sim \begin{cases} P & , \quad \text{for } \theta < d-1 \\ P \ln P & , \quad \text{for } \theta = d-1 \\ P^{\theta/(d-1)} & , \quad \text{for } \theta > d-1 \end{cases} .$$

N. Ogawa, T. Takayanagi, and T. Ugajin, JHEP **1201**, 125 (2012).
L. Huijse, S. Sachdev, B. Swingle, Physical Review B **85**, 035121 (2012)

Generalized holography

$$ds^2 = \frac{1}{r^2}\left(-\frac{dt^2}{r^{2d(z-1)/(d-\theta)}} + r^{2\theta/(d-\theta)}dr^2 + dx_i^2\right)$$

Application of the Ryu-Takayanagi minimal area formula to a dual Einstein-Maxwell-dilaton theory yields

$$S_E \sim \begin{cases} P & , \quad \text{for } \theta < d-1 \\ P \ln P & , \quad \text{for } \theta = d-1 \\ P^{\theta/(d-1)} & , \quad \text{for } \theta > d-1 \end{cases} .$$

The non-Fermi liquid has log-violation of "area law", and this appears precisely at the correct value $\theta = d - 1$!

N. Ogawa, T. Takayanagi, and T. Ugajin, JHEP **1201**, 125 (2012).
L. Huijse, S. Sachdev, B. Swingle, Physical Review B **85**, 035121 (2012)

Generalized holography

$$ds^2 = \frac{1}{r^2}\left(-\frac{dt^2}{r^{2d(z-1)/(d-\theta)}} + r^{2\theta/(d-\theta)}dr^2 + dx_i^2\right)$$

Application of the Ryu-Takayanagi minimal area formula to a dual Einstein-Maxwell-dilaton theory yields

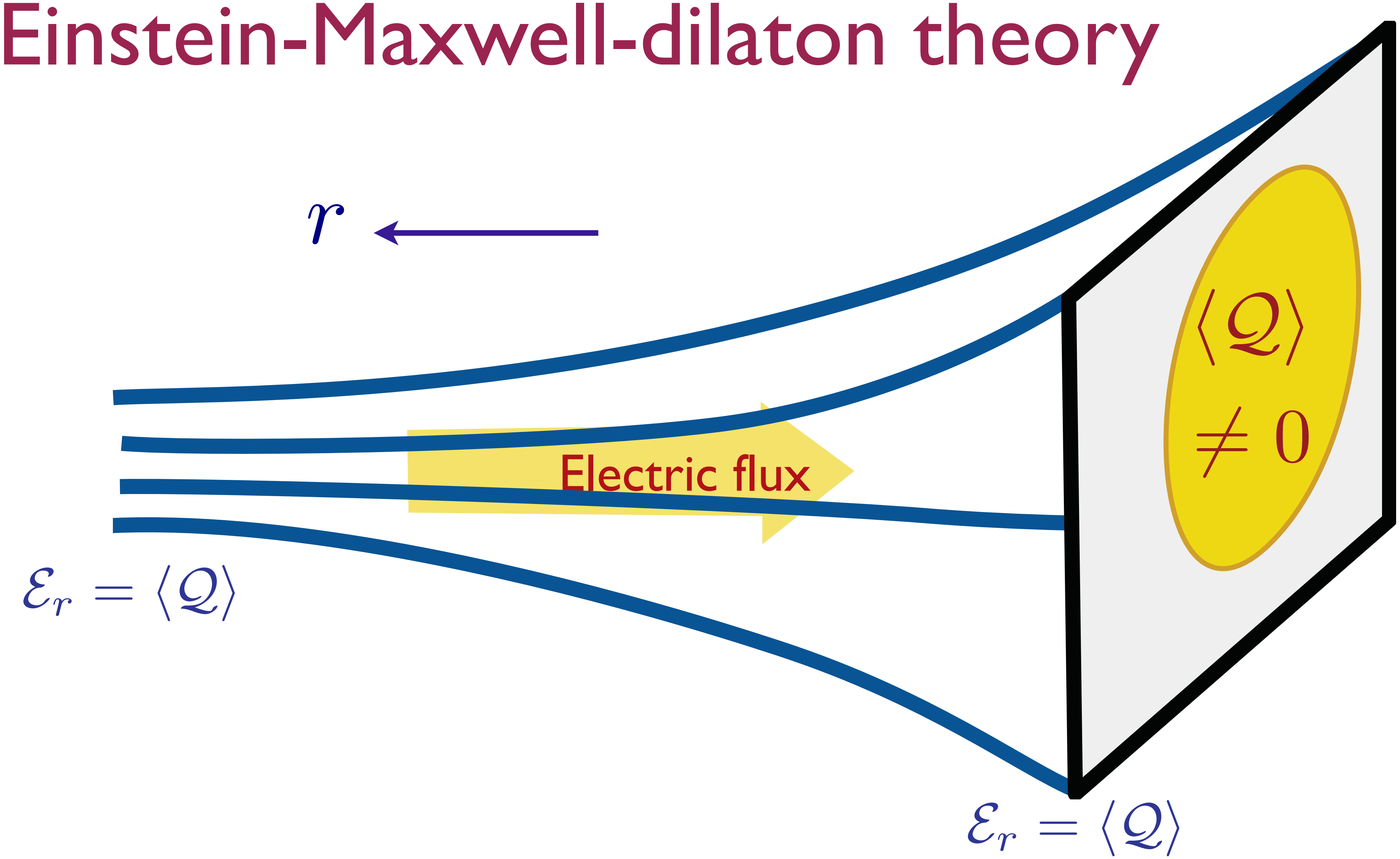
$$S_E \sim \begin{cases} P & , \quad \text{for } \theta < d-1 \\ P \ln P & , \quad \text{for } \theta = d-1 \\ P^{\theta/(d-1)} & , \quad \text{for } \theta > d-1 \end{cases}.$$

The non-Fermi liquid has log-violation of "area law", and this appears precisely at the correct value $\theta = d-1$!
Moreover, the co-efficient of $P \ln P$ computed holographically is independent of the shape of the entangling region just as expected for a circular Fermi surface!!

N. Ogawa, T. Takayanagi, and T. Ugajin, JHEP **1201**, 125 (2012).
L. Huijse, S. Sachdev, B. Swingle, Physical Review B **85**, 035121 (2012)

Holography of a non-Fermi liquid

Einstein-Maxwell-dilaton theory



$$\mathcal{S} = \int d^{d+2}x \sqrt{-g} \left[\frac{1}{2\kappa^2} \left(R - 2(\nabla\Phi)^2 - \frac{V(\Phi)}{L^2} \right) - \frac{Z(\Phi)}{4e^2} F_{ab} F^{ab} \right]$$

with $Z(\Phi) = Z_0 e^{\alpha\Phi}$, $V(\Phi) = -V_0 e^{-\beta\Phi}$, as $\Phi \to \infty$.

C. Charmousis, B. Gouteraux, B. S. Kim, E. Kiritsis and R. Meyer, JHEP **1011**, 151 (2010).
S. S. Gubser and F. D. Rocha, Phys. Rev. D **81**, 046001 (2010).
N. Iizuka, N. Kundu, P. Narayan and S. P. Trivedi, arXiv:1105.1162 [hep-th].

Holography of a non-Fermi liquid

$$\boxed{ds^2 = \frac{1}{r^2}\left(-\frac{dt^2}{r^{2d(z-1)/(d-\theta)}} + r^{2\theta/(d-\theta)}dr^2 + dx_i^2\right)}$$

The $r \to \infty$ limit of the metric of the Einstein-Maxwell-dilaton (EMD) theory has the most general form with

$$\begin{aligned}
\theta &= \frac{d^2\beta}{\alpha + (d-1)\beta} \\
z &= 1 + \frac{\theta}{d} + \frac{8(d(d-\theta)+\theta)^2}{d^2(d-\theta)\alpha^2}.
\end{aligned}$$

Holography of a non-Fermi liquid

$$ds^2 = \frac{1}{r^2}\left(-\frac{dt^2}{r^{2d(z-1)/(d-\theta)}} + r^{2\theta/(d-\theta)}dr^2 + dx_i^2\right)$$

Computation of the entanglement entropy in the EMD theory via the Ryu-Takayanagi formula for $\theta = d - 1$ yields

$$\mathcal{S}_E = \mathcal{C}_E \mathcal{Q}^{(d-1)/d} P \ln P$$

where $\mathcal{C}_E$ is independent of UV details.

L. Huijse, S. Sachdev, B. Swingle, Physical Review B **85**, 035121 (2012)

Holography of a non-Fermi liquid

$$ds^2 = \frac{1}{r^2}\left(-\frac{dt^2}{r^{2d(z-1)/(d-\theta)}} + r^{2\theta/(d-\theta)}dr^2 + dx_i^2\right)$$

Computation of the entanglement entropy in the EMD theory via the Ryu-Takayanagi formula for $\theta = d - 1$ yields

$$\mathcal{S}_E = \mathcal{C}_E \mathcal{Q}^{(d-1)/d} P \ln P$$

where $\mathcal{C}_E$ is independent of UV details.
This is precisely as expected for a Fermi surface, when combined with the Luttinger theorem!

L. Huijse, S. Sachdev, B. Swingle, Physical Review B **85**, 035121 (2012)

Holography of a non-Fermi liquid

$$ds^2 = \frac{1}{r^2}\left(-\frac{dt^2}{r^{2d(z-1)/(d-\theta)}} + r^{2\theta/(d-\theta)}dr^2 + dx_i^2\right)$$

Computation of the entanglement entropy in the EMD

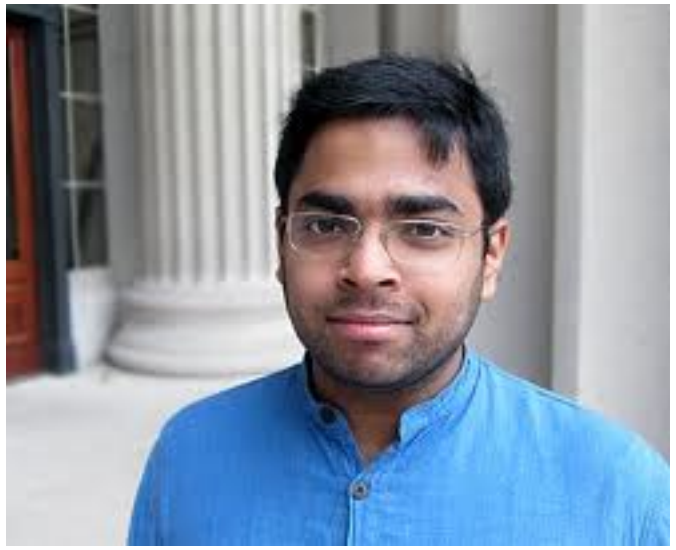
Corrections to the Ryu-Takayangi formula from the bulk entanglement entropy also correctly yield the entanglement entropy and Fermi surface volume of probe fermions.

L. Huijse, S. Sachdev, and B. Swingle, arXiv:1308.3234

combined with the Luttinger theorem!

L. Huijse, S. Sachdev, B. Swingle, Physical Review B **85**, 035121 (2012)

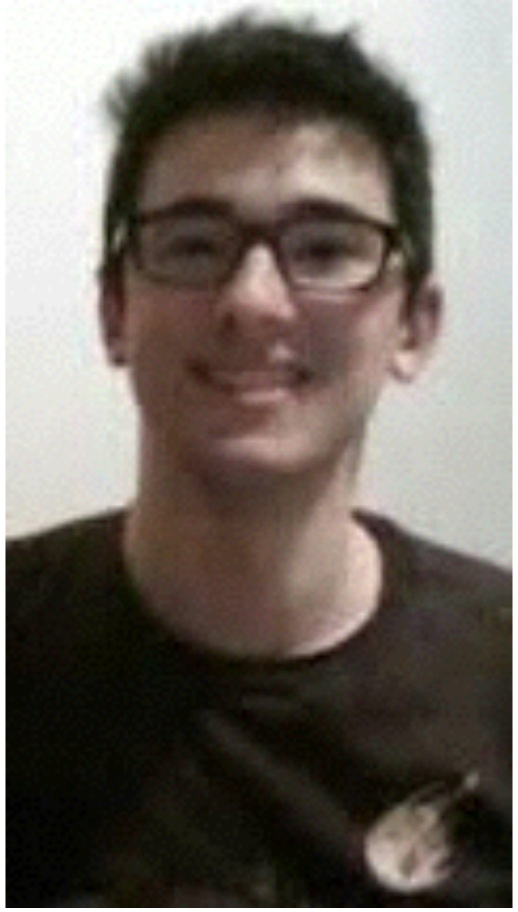
Sean Hartnoll
Stanford



Raghu Mahajan
Stanford

Matthias Punk
Innsbruck

Koenraad Schalm
Leiden



Andrew Lucas
Harvard

Quantum criticality of Ising-nematic ordering in a metal

Boltzmann view of electrical transport:

- Identify charge carriers: electrons near the Fermi surface. Compute the scattering rate of these charged excitations off the bosonic ϕ fluctuations.

Quantum criticality of Ising-nematic ordering in a metal

Boltzmann view of electrical transport:

- Identify charge carriers: electrons near the Fermi surface. Compute the scattering rate of these charged excitations off the bosonic ϕ fluctuations.

- Analogous to electron-phonon scattering in metals, where we have "Bloch's law": a resistivity $\rho(T) \sim T^5$.

Quantum criticality of Ising-nematic ordering in a metal

Boltzmann view of electrical transport:

- Identify charge carriers: electrons near the Fermi surface. Compute the scattering rate of these charged excitations off the bosonic ϕ fluctuations.

- Analogous to electron-phonon scattering in metals, where we have "Bloch's law": a resistivity $\rho(T) \sim T^5$.

- "Bloch's law" for the Ising-nematic critical point yields $\rho(T) \sim T^{4/3}$.

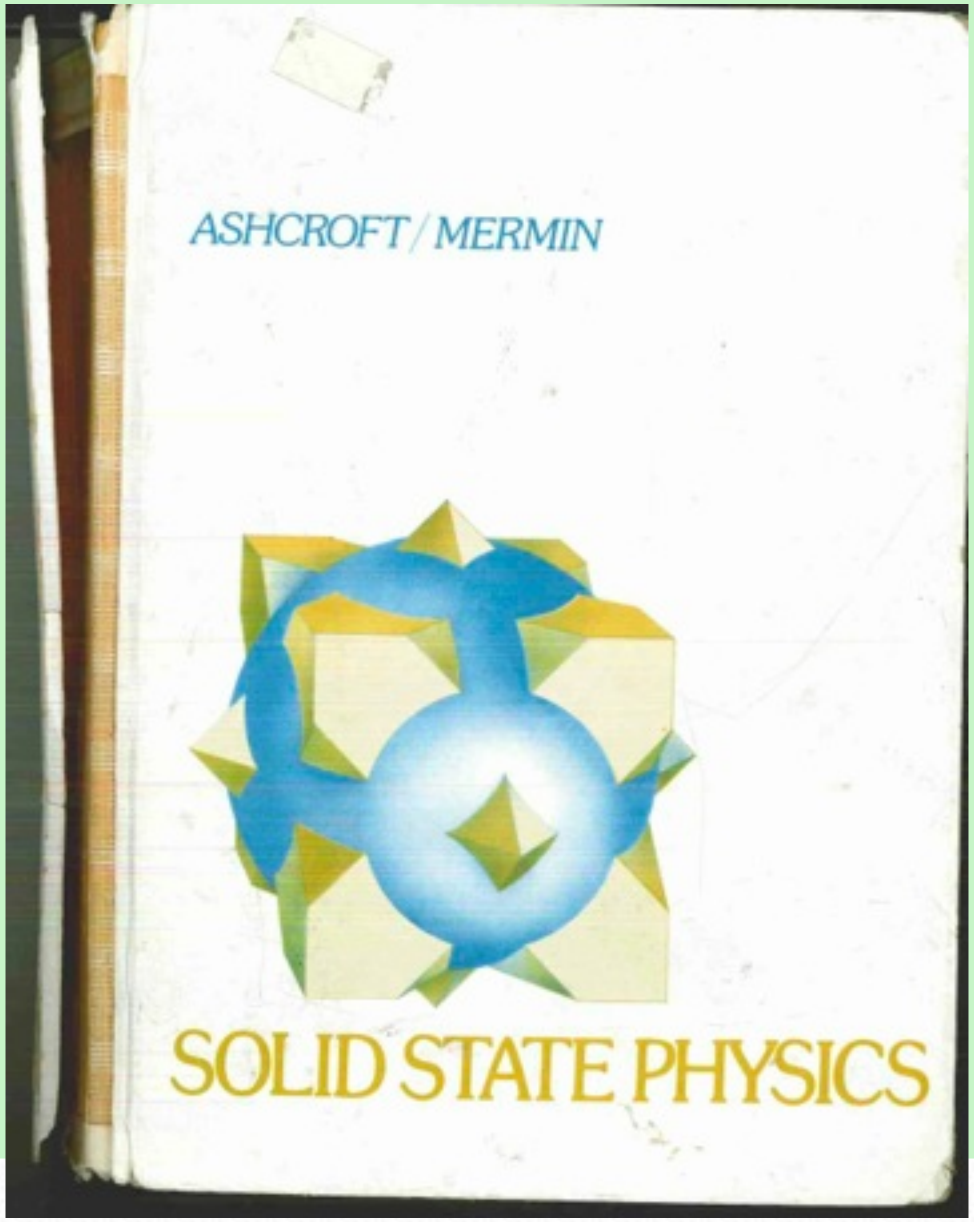
Quantum criticality of Ising-nematic ordering in a metal

Boltzmann view of electrical transport:

- Identify charge carriers: electrons
pute the scattering rate of these ch
ϕ fluctuations.

- Analogous to electron-phonon scat
"Bloch's law": a resistivity $\rho(T) \sim$

- "Bloch's law" for the Ising-nemati
$\rho(T) \sim T^{4/3}$.



However, this ignores "phonon drag"

PHONON DRAG

Peierls[28] pointed out a way in which the low temperature resistivity might decline more rapidly than T^5.

[28] R. E. Peierls, *Ann. Phys.* (5) **12**, 154 (1932).

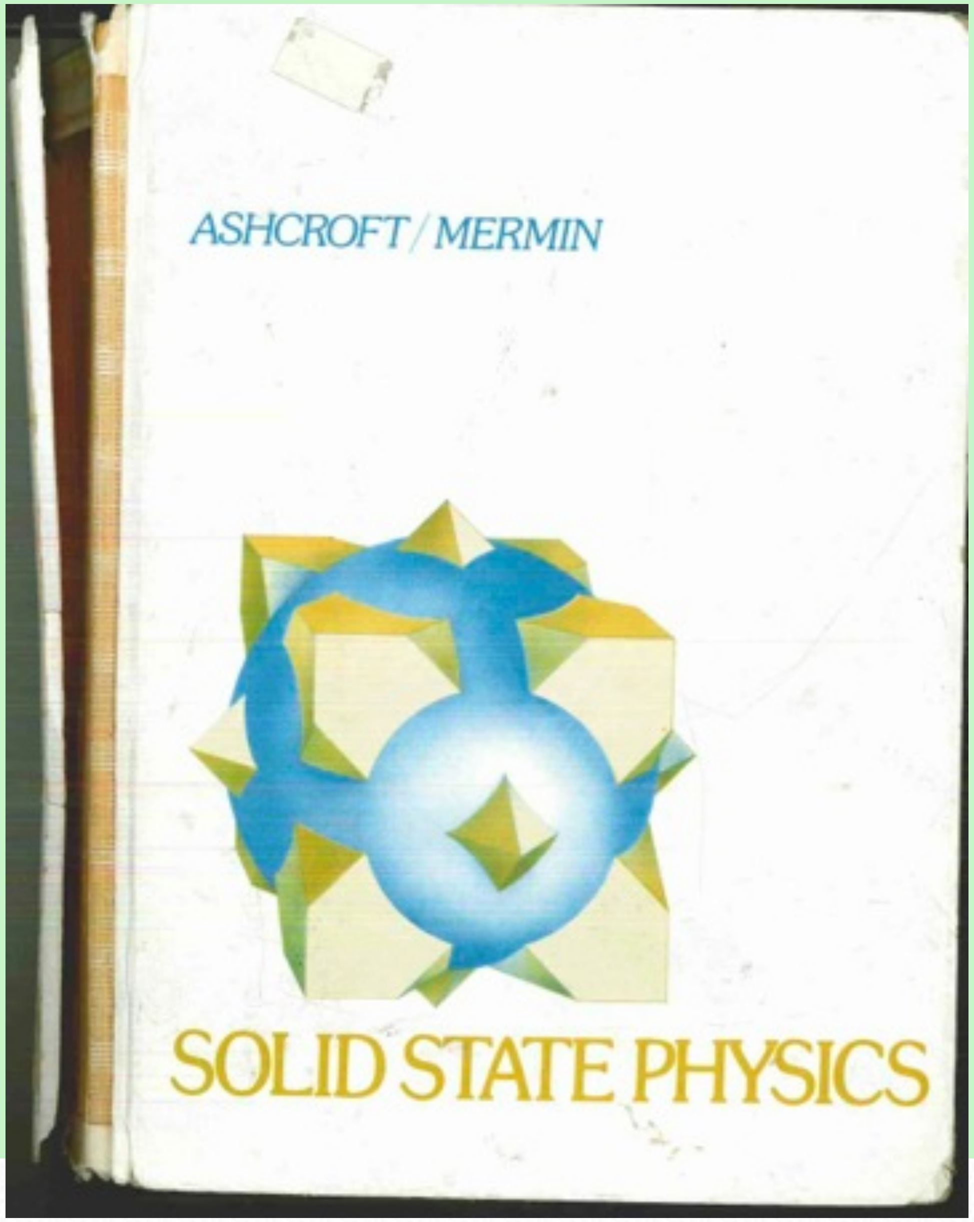
Quantum criticality of Ising-nematic ordering in a metal

Boltzmann view of electrical transport:

- Identify charge carriers: electrons
pute the scattering rate of these ch
ϕ fluctuations.

- Analogous to electron-phonon scat
"Bloch's law": a resistivity $\rho(T) \sim$

- "Bloch's law" for the Ising-nemati
$\rho(T) \sim T^{4/3}$.



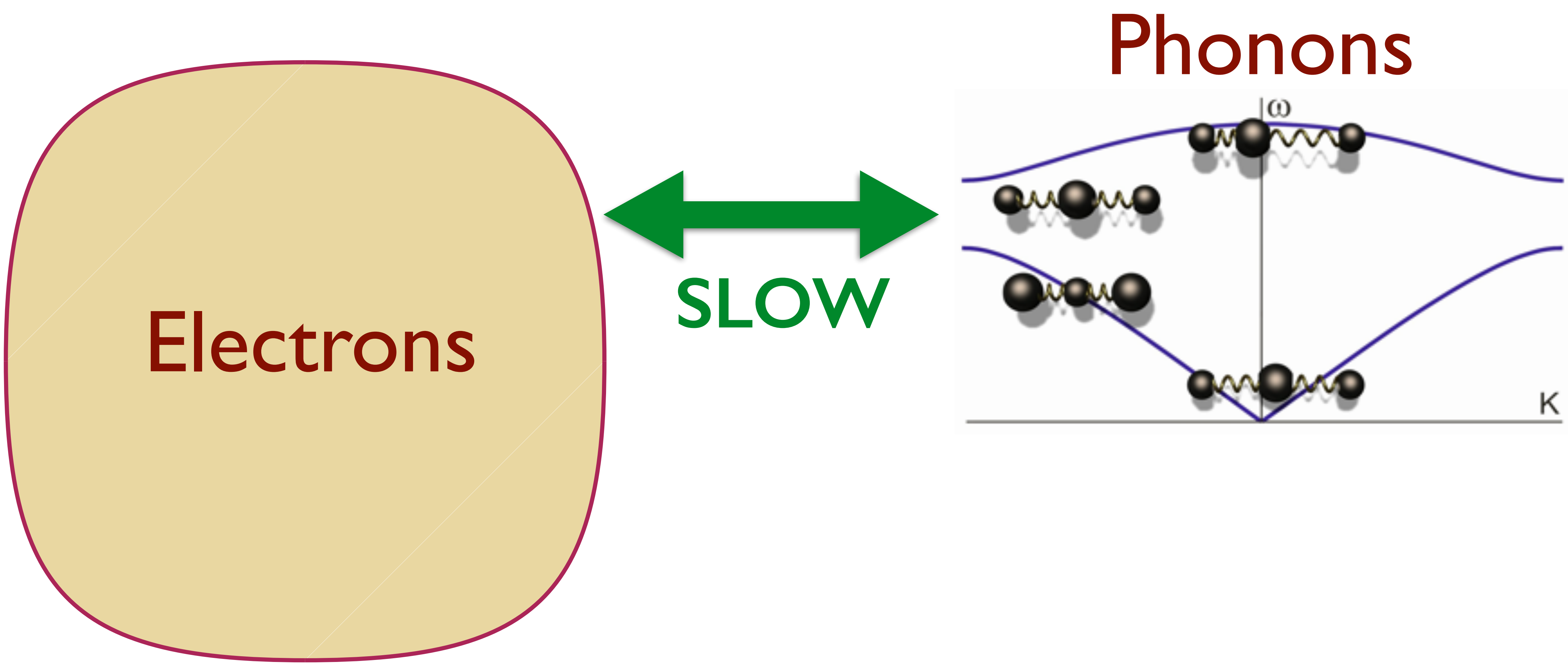
However, this ignores "phonon drag"

PHONON DRAG

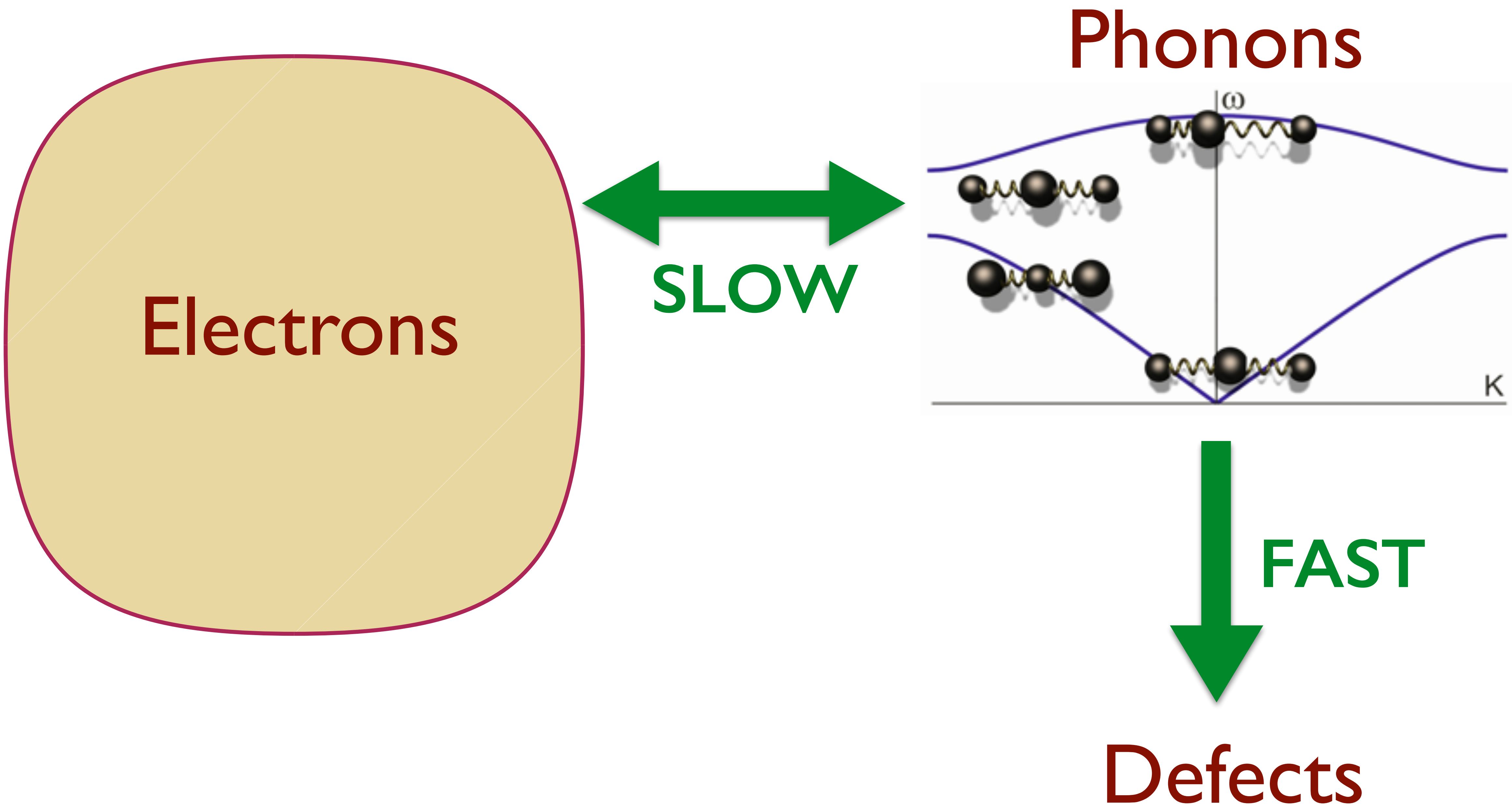
Peierls[28] pointed out a way in which the low temperature resistivity might decline more rapidly than T^5. This behavior has yet to be observed.

[28] R. E. Peierls, *Ann. Phys.* (5) **12**, 154 (1932).

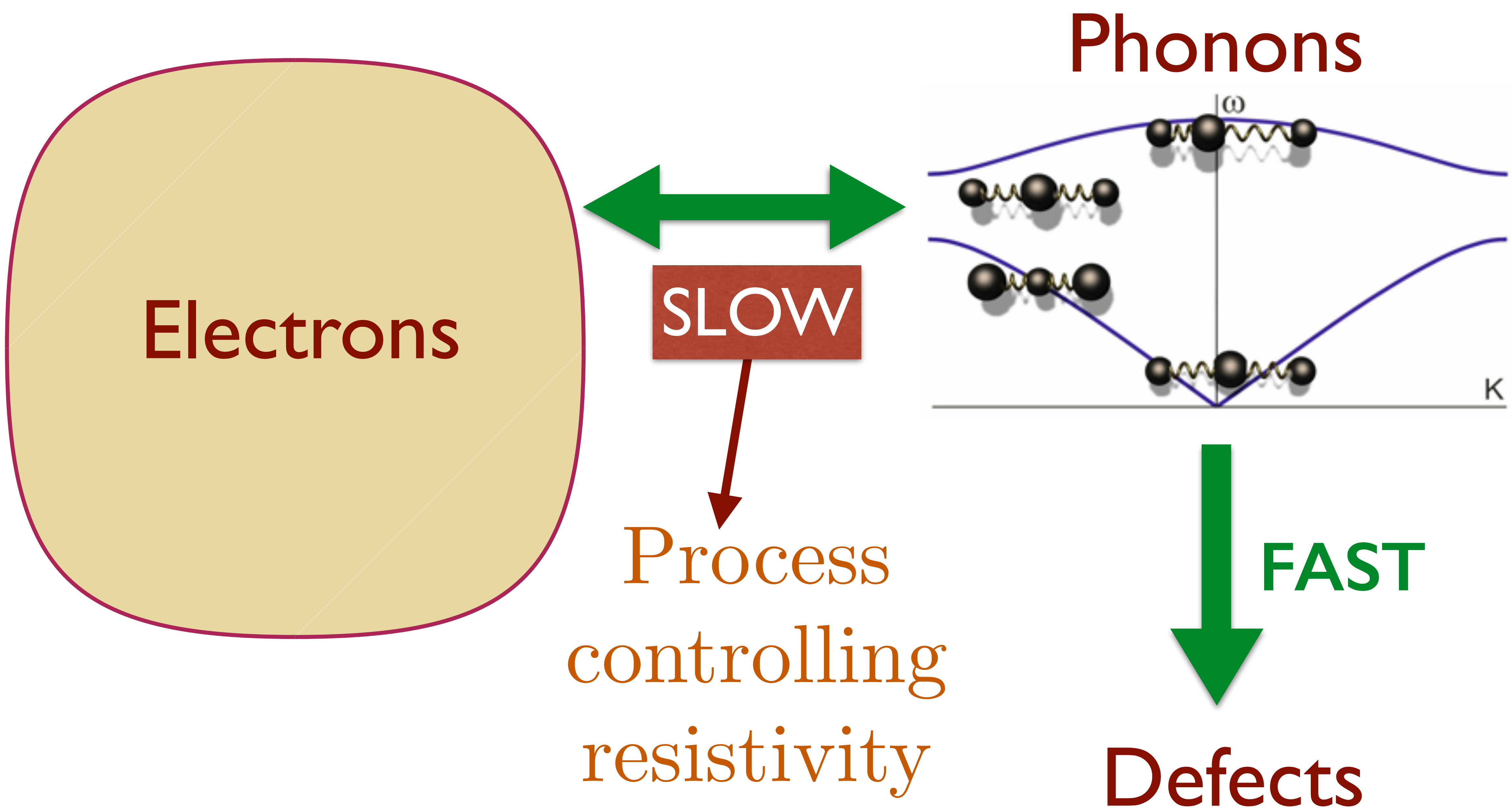
Rates of Momentum Flow



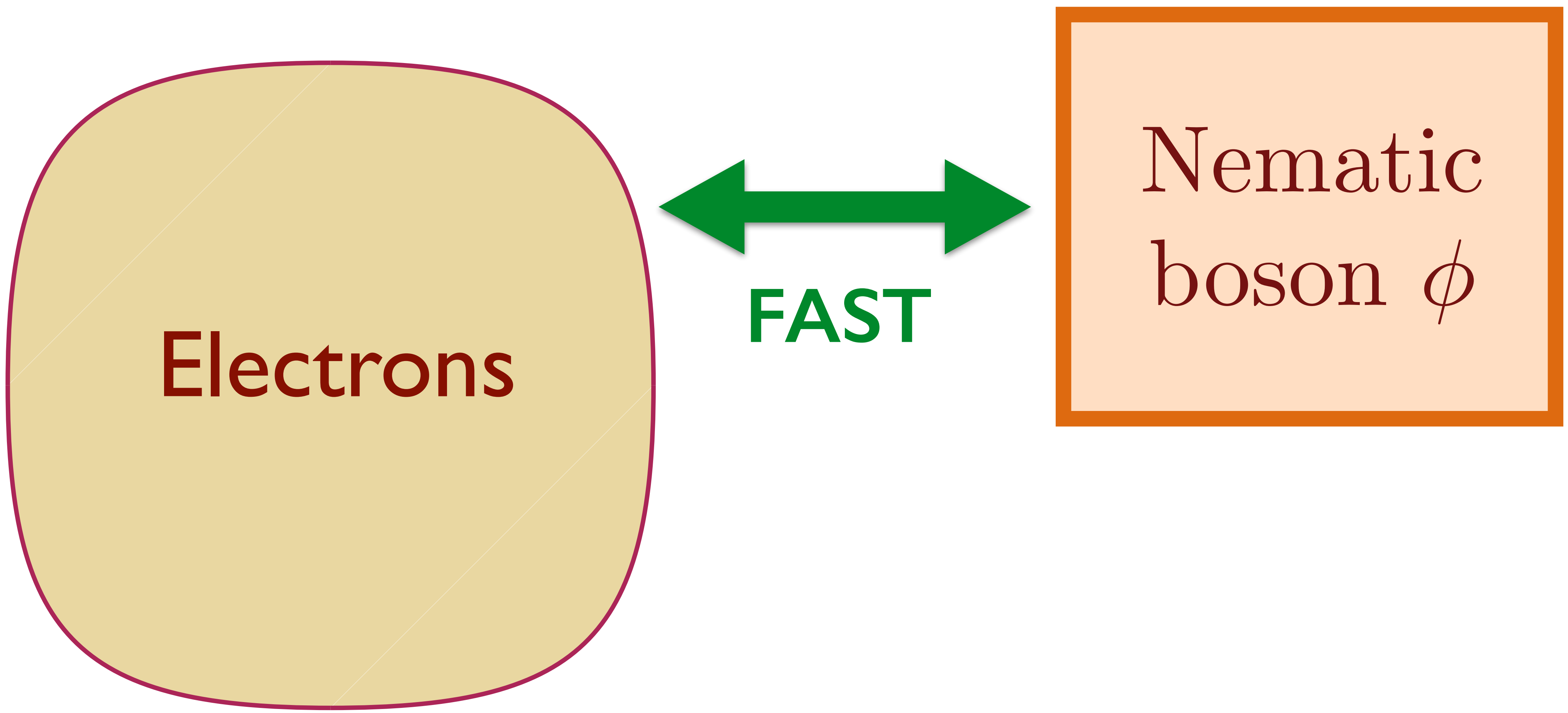
Rates of Momentum Flow



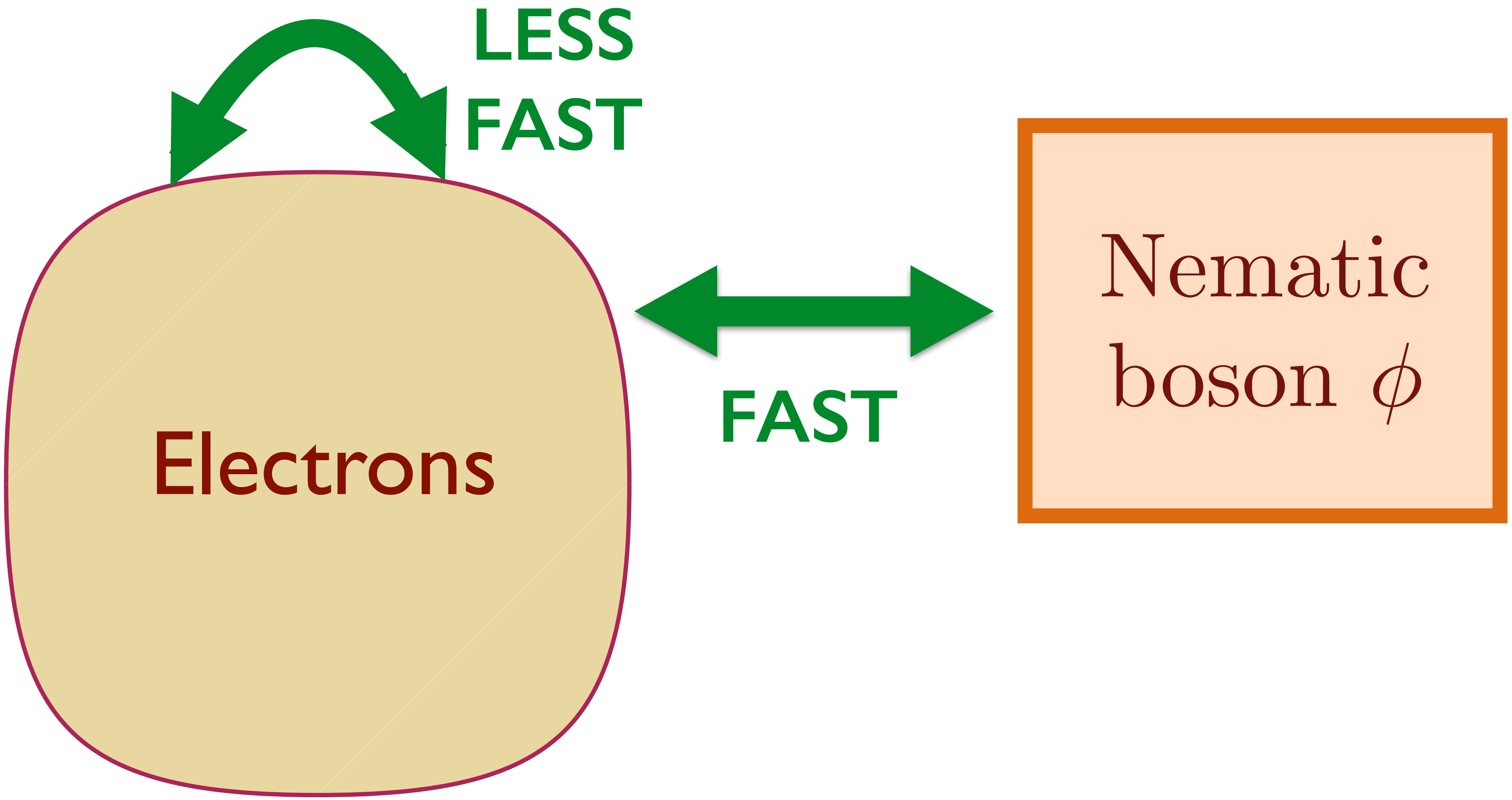
Rates of Momentum Flow



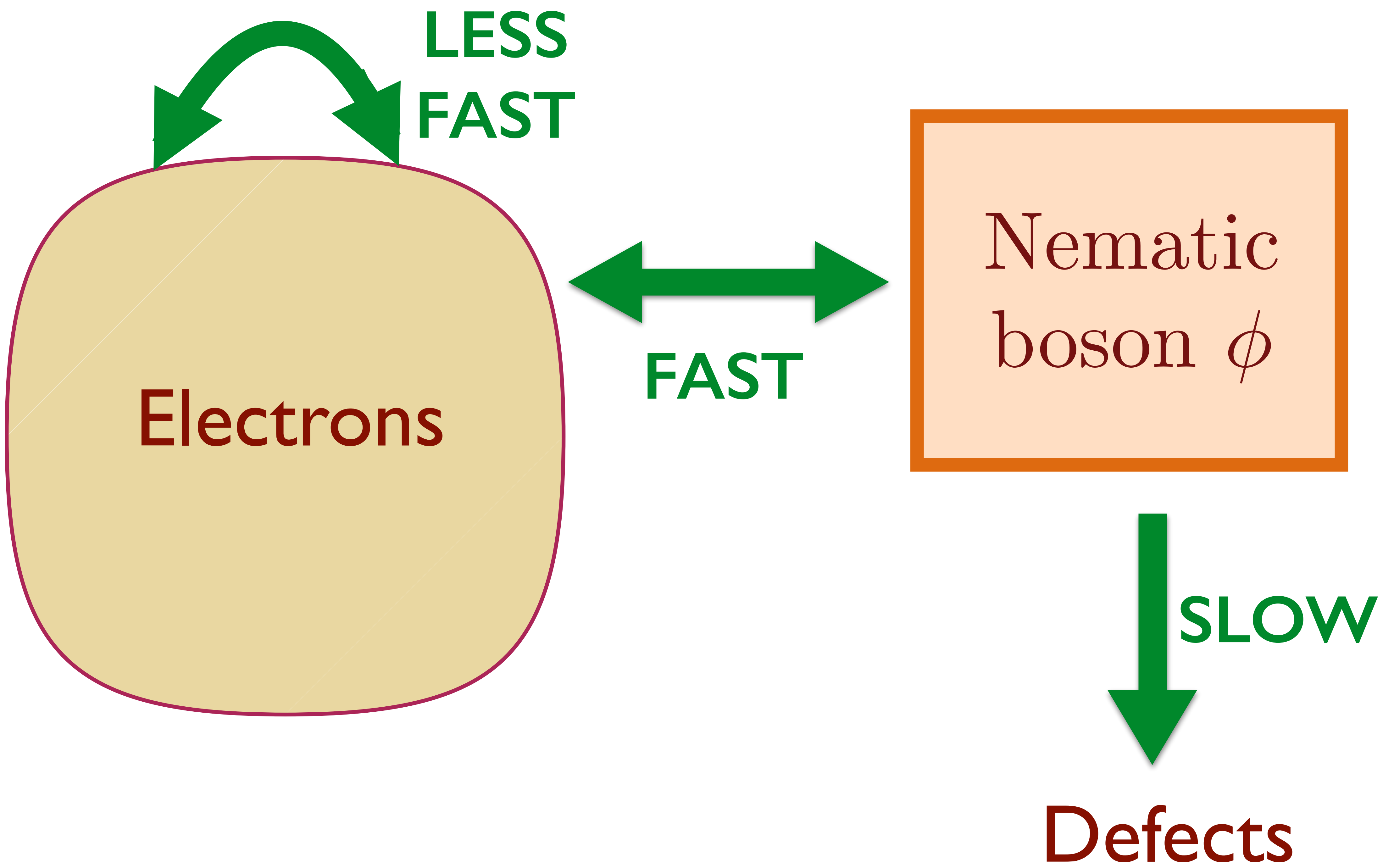
Rates of Momentum Flow



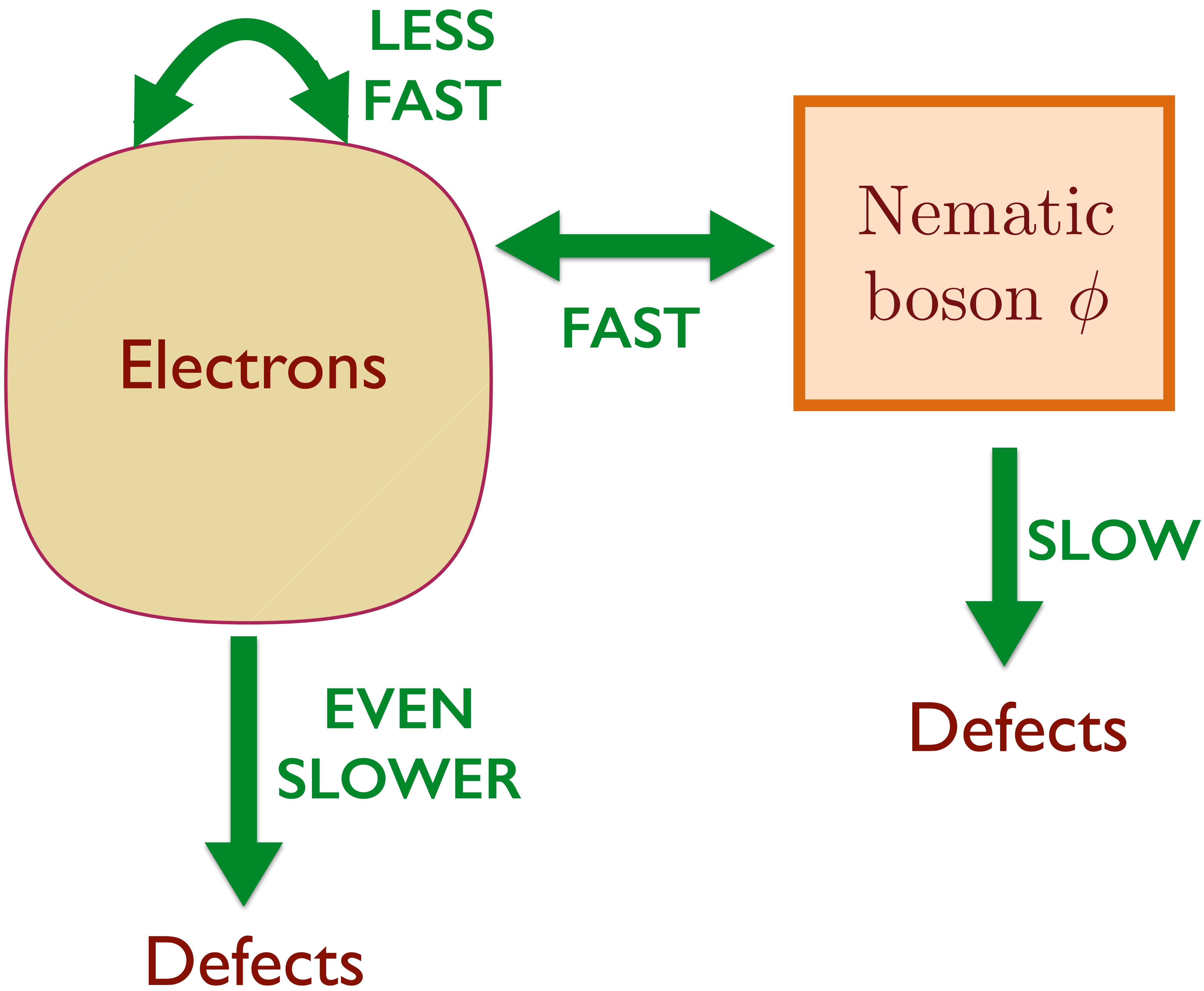
Rates of Momentum Flow



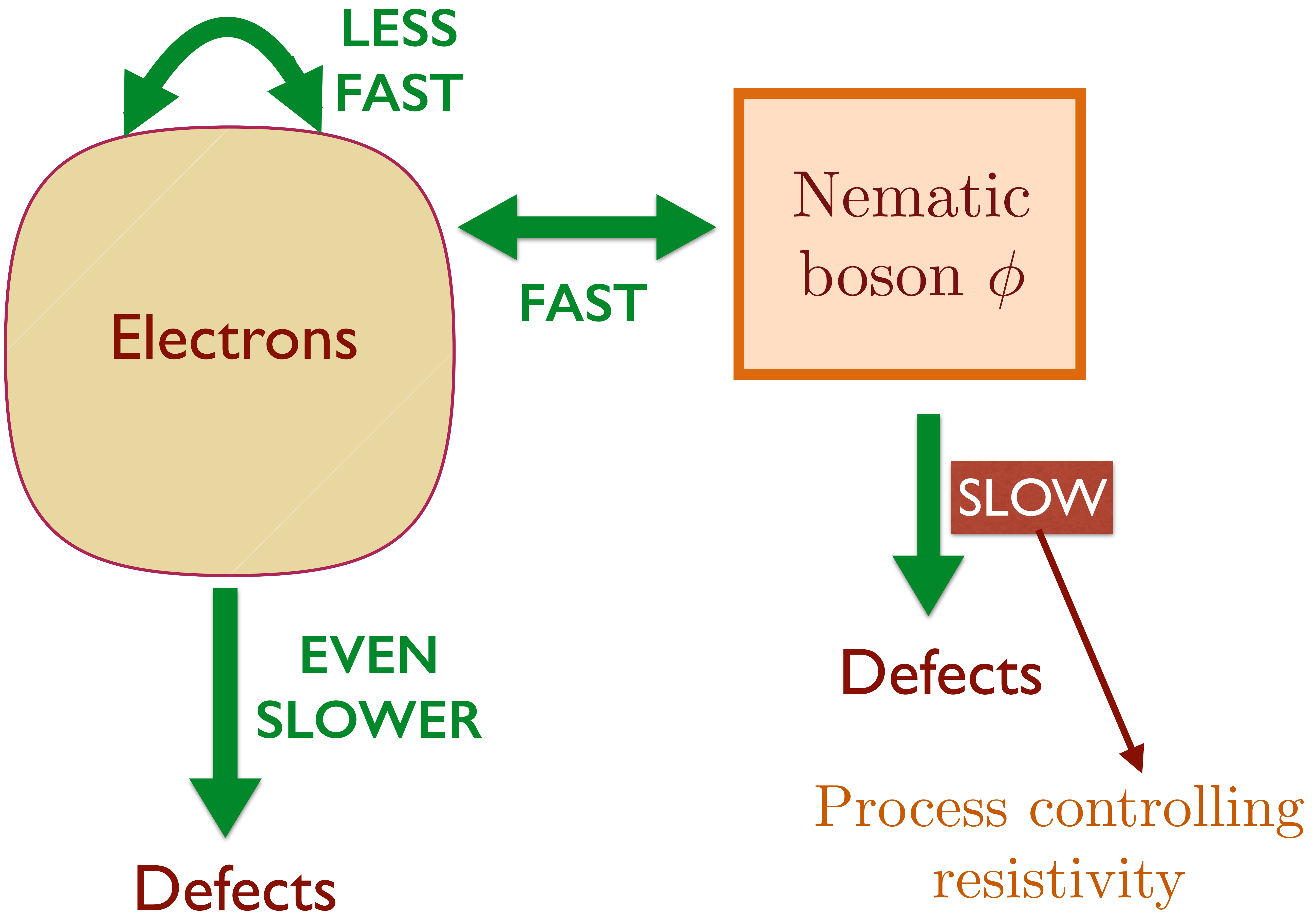
Rates of Momentum Flow



Rates of Momentum Flow



Rates of Momentum Flow



Quantum criticality of Ising-nematic ordering in a metal

Transport without quasiparticles:

- Focus on the interplay between J_μ and $T_{\mu\nu}$!

The most-probable state with a non-zero current $\mathbf{J}$ has a non-zero momentum $\mathbf{P}$ (and vice versa). At non-zero density, $\mathbf{J}$ “drags” $\mathbf{P}$.

Quantum criticality of Ising-nematic ordering in a metal

Transport without quasiparticles:

- Focus on the interplay between J_μ and $T_{\mu\nu}$!

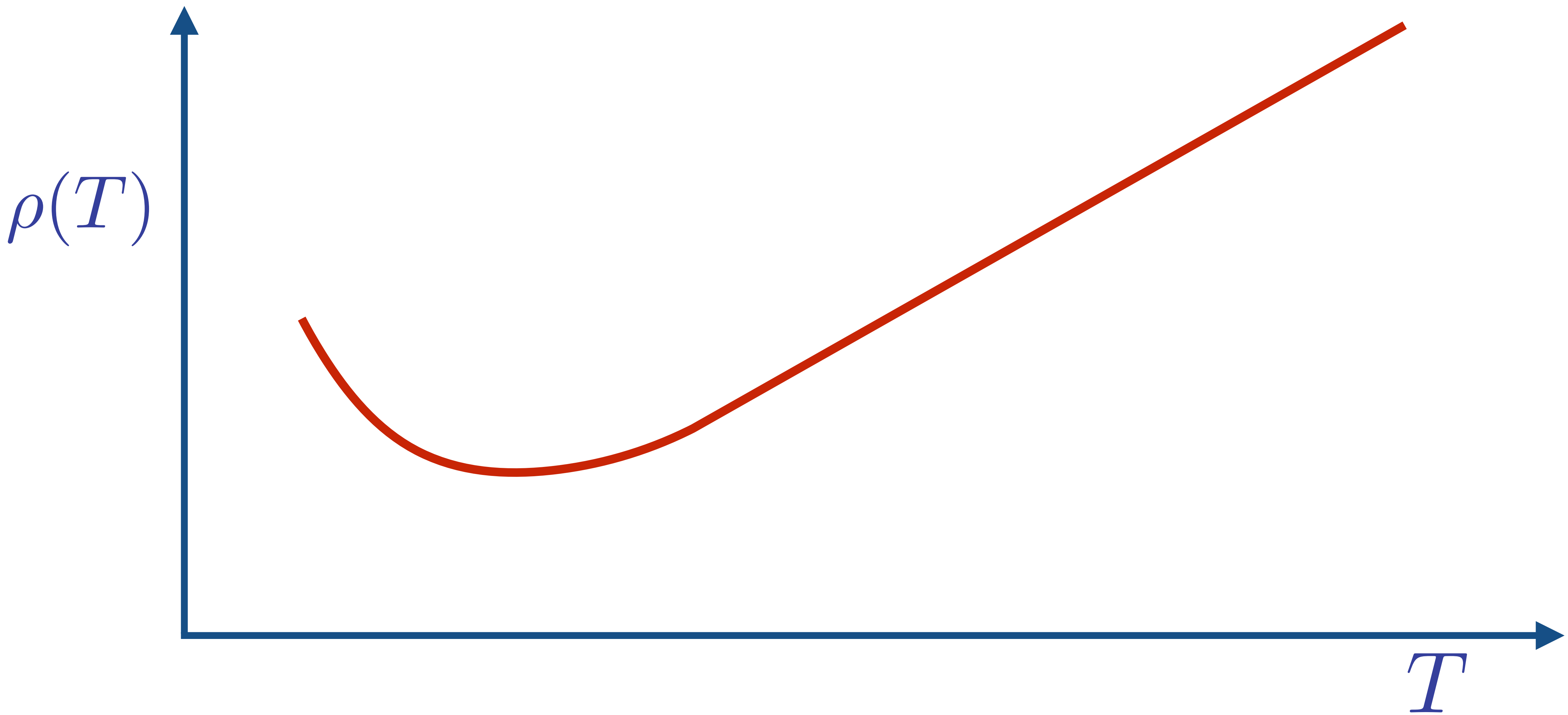
The most-probable state with a non-zero current **J** has a non-zero momentum **P** (and vice versa). At non-zero density, **J** "drags" **P**.

The resistivity of this metal is *not* determined by the scattering rate of charged excitations near the Fermi surface, but by the dominant rate of momentum loss by *any* excitation, whether neutral or charged, or fermionic or bosonic

Quantum criticality of Ising-nematic ordering in a metal

Transport without quasiparticles:

Resistivity from random-field disorder

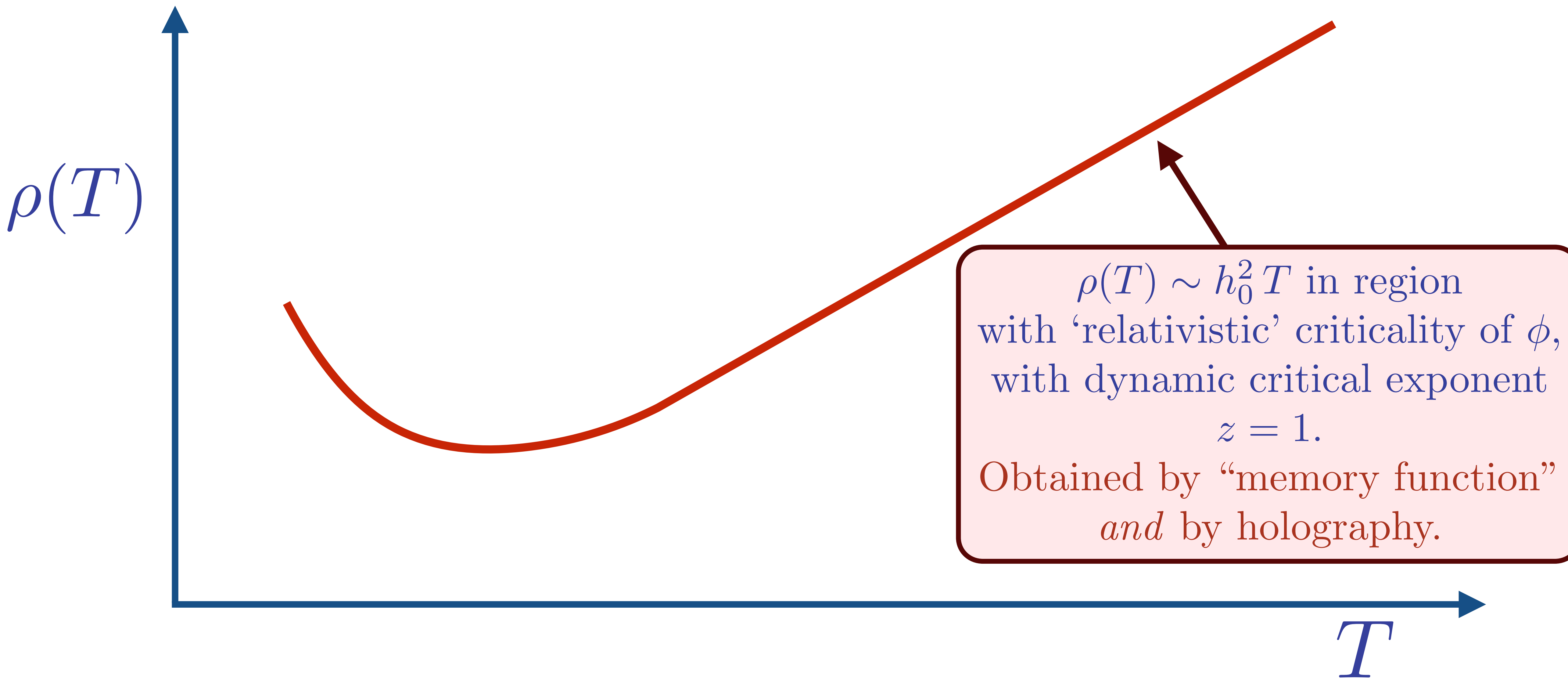


S. A. Hartnoll, R. Mahajan, M. Punk and S. Sachdev, Physical Review B **89**, 155130 (2014)
A. Lucas, S. Sachdev, and K. Schalm, Physical Review D **89**, 066018 (2014)

Quantum criticality of Ising-nematic ordering in a metal

Transport without quasiparticles:

Resistivity from random-field disorder

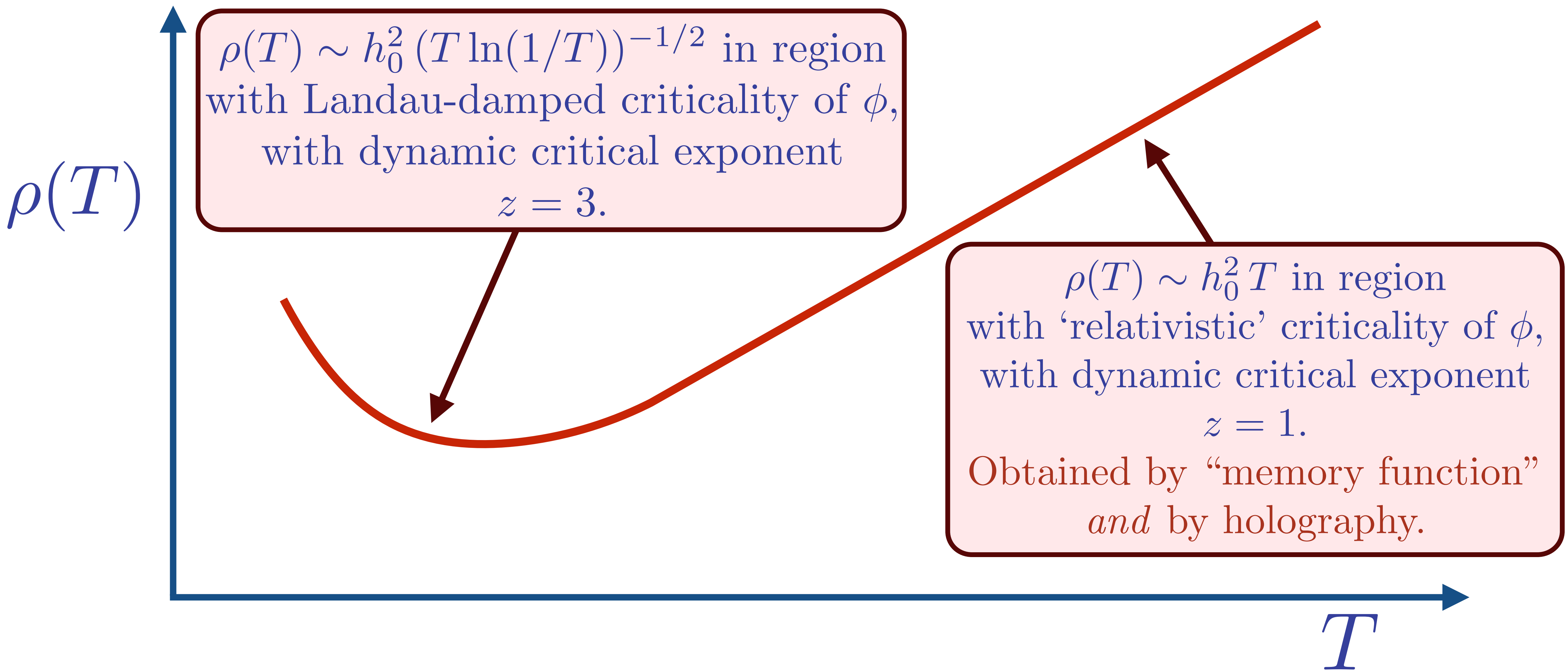


S. A. Hartnoll, R. Mahajan, M. Punk and S. Sachdev, Physical Review B **89**, 155130 (2014)
A. Lucas, S. Sachdev, and K. Schalm, Physical Review D **89**, 066018 (2014)

Quantum criticality of Ising-nematic ordering in a metal

Transport without quasiparticles:

Resistivity from random-field disorder

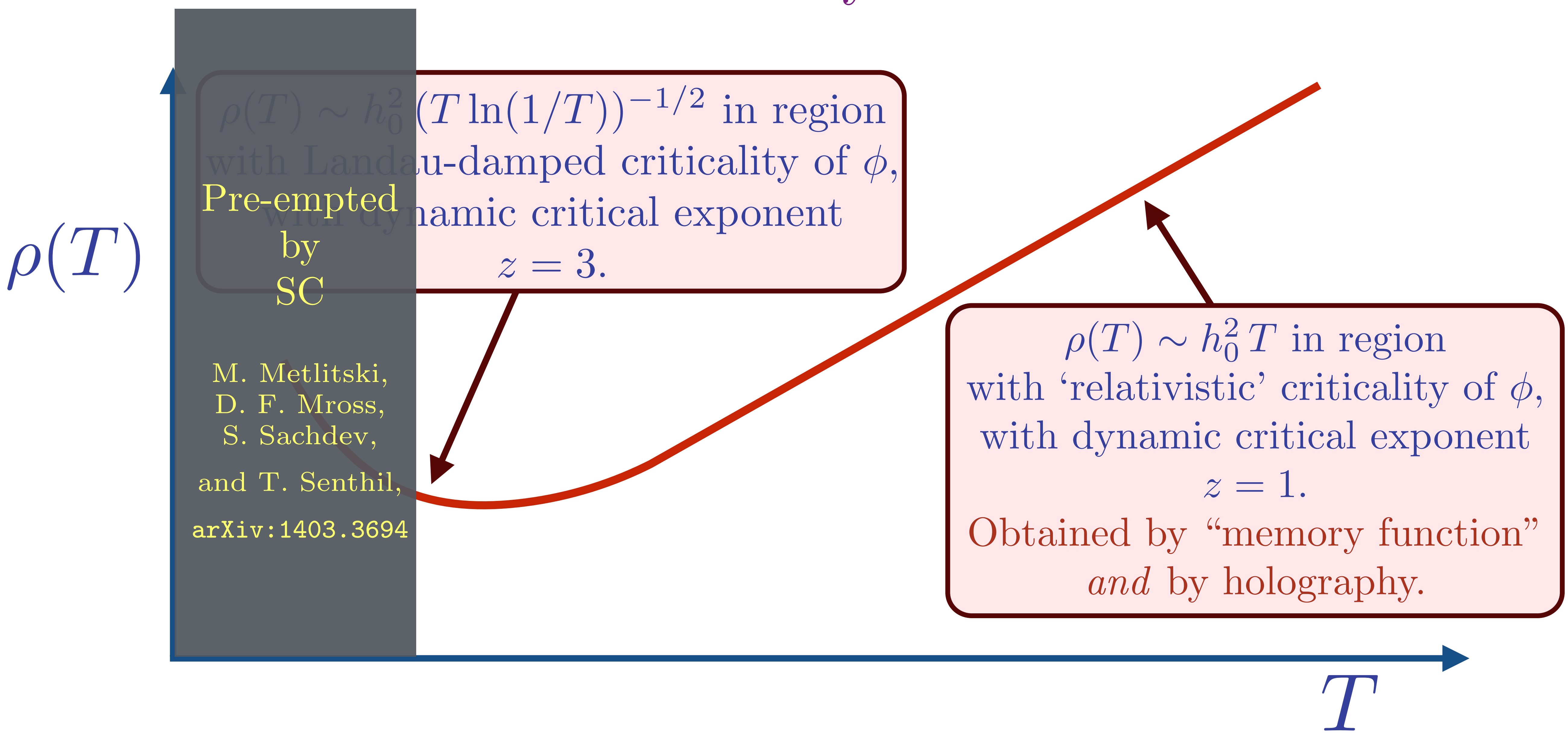


S. A. Hartnoll, R. Mahajan, M. Punk and S. Sachdev, Physical Review B **89**, 155130 (2014)
A. Lucas, S. Sachdev, and K. Schalm, Physical Review D **89**, 066018 (2014)

Quantum criticality of Ising-nematic ordering in a metal

Transport without quasiparticles:

Resistivity from random-field disorder



S. A. Hartnoll, R. Mahajan, M. Punk and S. Sachdev, Physical Review B **89**, 155130 (2014)
A. Lucas, S. Sachdev, and K. Schalm, Physical Review D **89**, 066018 (2014)

- Strongly-coupled quantum criticality leads to a novel regime of quantum dynamics without quasiparticles.

- Strongly-coupled quantum criticality leads to a novel regime of quantum dynamics without quasiparticles.

- The simplest examples are conformal field theories in 2+1 dimensions, realized by ultracold atoms in optical lattices. Quantitative predictions for transport by combining quantum Monte Carlo and holography.

- Strongly-coupled quantum criticality leads to a novel regime of quantum dynamics without quasiparticles.

- The simplest examples are conformal field theories in 2+1 dimensions, realized by ultracold atoms in optical lattices. Quantitative predictions for transport by combining quantum Monte Carlo and holography.

- Exciting recent progress on the description of transport in metallic states without quasiparticles, via field theory and holography.