Eigenfunction fractality and pseudogap state near superconductor-insulator transition

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We develop a theory of a pseudogap state appearing near the superconductor-insulator transition in strongly disordered metals with attractive interaction. We show that such an interaction combined with the fractal nature of the single particle wave functions near the mobility edge leads to an anomalously large single particle gap in the superconducting state near SI transition that persists and even increases in the insulating state long after the superconductivity is destroyed. We give analytic expressions for the value of the pseudogap in terms of the inverse participation ratio of the corresponding localization problem.

Rapidly growing number of experiments\textsuperscript{1,2,3,4,5,6,7,8,9} on various disordered superconductors show that a novel phase often appears on the insulating side of the superconductor-insulator transition. As the disorder strength is increased the superconductivity is suppressed leading to a strange insulator characterized by a large thermally assisted resistance with a small but hard gap (Fig. 1). Further increasing the disorder one gets the usual variable range hopping behavior. Experimentally, the phase diagram of disordered superconductors is often explored by varying the applied magnetic field. On the superconducting side of the transition a relatively small field destroys the superconductivity resulting in a hard gap insulating state. At larger magnetic fields the resistance and the gap drop\textsuperscript{4,5,6,7,8}. This is observed only in a narrow window of disorder strengths, away from this window on a superconducting side the application of magnetic field converts the superconductor into a normal metal as usual.

It is tempting to explain these data by the formation of localized Cooper pairs\textsuperscript{9,10}. In this picture the superconductivity is due to a fragile coherence between localized Cooper pairs, while the energy to break the pair is much larger and remains finite even when the coherence (and thus the superconductivity) is destroyed. The hypothesis of preformed Cooper pairs is further confirmed by the behavior of these superconductors at higher temperatures. On the insulating side of the transition in thick (effectively 3D) films one observes\textsuperscript{9} Arrhenius temperature behavior of the resistivity, $R(T) \propto \exp(T_I/T)$, at low temperatures. The experimental value of the activation energy, $T_I$, is somewhat larger than the superconducting gap in less disordered samples and grows with the disorder\textsuperscript{9,10}. However, at higher temperatures this behavior is replaced\textsuperscript{9} by Mott’s variable range hopping $R(T) \sim \exp[T_M/T]^{1/4}$. This can be understood if the insulating pseudogap is due to preformed Cooper pairs with a relatively large pairing energy $T_I$.

In this Letter we show that preformed Cooper pairs appear in the semi-microscopic three-dimensional model that contains only low energy electrons with weak BCS-type attraction and a strong random potential that leads to Anderson localization of single-particle states.

In the presence of preformed Cooper pairs the parity effect should arise – the ground state energy with even number of electrons is lower then that for the closest odd number. The corresponding parity gap $\Delta_p \sim T_I \sim \delta/\ln(\delta/\Delta)$ has been calculated by Matveev and Larkin\textsuperscript{11} for small superconducting grains. Here $\delta = 1/(\nu_0 L^3)$ is the mean level spacing in the grain, $\Delta < < \delta$ is the energy gap in the bulk superconductor, and $\nu_0$ and $L^3$ are the density of orbital states and the volume of the grain, respectively. In this Letter we argue that the result of Ref.\textsuperscript{11} can be generalized to bulk Anderson insulators. In this case, $L$ is replaced by the localization radius $L_{loc}$, so that $\delta \rightarrow \delta_L = 1/(\nu_0 L^3_{loc})$, and the BSC $\Delta$ is replaced by the superconducting gap $\Delta_{crit}$ at the Anderson transition point. The fractal nature of near-critical wavefunctions characterized by the fractal dimension $D_f < 3$ determines the reduction factor $(\delta_L/\Delta_{crit})^{1-D_f/3} \gg 1$ that replaces $\ln(\delta/\Delta)$ in the Matveev-Larkin formula.

\begin{figure}[h]
\centering
\includegraphics[width=0.5\textwidth]{fig1.png}
\caption{Schematics of low-temperature phase diagram of disordered superconductors in the vicinity of the superconductor-insulator transition as a function of the magnetic field, $B$, and disorder strength.}
\end{figure}
We consider two different regimes. In more disordered materials where \( \delta_L \gg \Delta_{\text{crit}} \) the Cooper instability and superconductive long-range order disappear. However, the attraction between electrons persists as long as \( \delta_L \) remains smaller than the Debye frequency \( \omega_D \) and results in the “local” pairing of electrons with opposite spins occupying the same localized state. We show below that in this regime the hard-gap insulator is formed with properties similar to those observed in. We further argue that in the less localized regime, when \( \delta_L \leq \Delta_{\text{crit}} \), the unusual superconductive state with a pseudogap is formed. The features of this state are: (i) single-electron excitation gap \( \Delta_1 \) is larger than superconductive gap \( \Delta \) so that the ratio \( \Delta_1/\Delta_{\text{c}} \) is anomalously high, and (ii) insulating trend in the \( R(T) \) curves exists above \( T_{\text{c}} \).

We assume that superconductivity is due to attraction between electrons that originates at high energy scales \( \sim \omega_D \) and that it is affected by localization of electron wave functions only for a very large disorder \( \delta_L \geq \omega_D \). In a fermion system with weak attraction one can leave only the pair interaction terms in the Hamiltonian leading to the usual BCS model in the basis of localized states:\(^{12}\)

\[
H = \sum_{j\sigma} \epsilon_j c_{j\sigma}^\dagger c_{j\sigma} - \frac{\lambda}{\nu_0} \sum_{j,k} M_{jk} c_{j\uparrow}^\dagger c_{k\downarrow},
\]

where

\[
M_{jk} = \int d^3r \frac{\psi_j^2(r)}{\psi_k^2(r)},
\]

\( \lambda \) is dimensionless Cooper coupling constant, \( \epsilon_j \) is the single-particle energy of the state \( j \), and \( c_{j\sigma} \) is the corresponding electron operator for the spin projection \( \sigma \).

Physical properties of the electron system are controlled by the electrons near the Fermi level, so a very important implicit ingredient of the model is the statistics of matrix elements \( M_{jk} \) between eigenstates in the vicinity of Anderson mobility edge. The key feature of these nearly-critical wavefunctions is their fractal structure\(^ {12}\) that shows in the anomalous scaling of diagonal matrix elements \( M_{jj} \approx M_j \) (“inverse participation ratios”, IPRs) with localization length: typical IPR \( M \propto L_{\text{loc}}^{-D_2} \), where fractal dimension \( D_2 < 3 \). Numerical studies\(^ {13,15}\) indicate that \( D_2 = 1.30 \pm 0.05 \) for the standard 3D Anderson transition. The IPR distribution function \( P(M_j) \) has been studied in\(^ {15}\). Scaling theory of localization predicts that near the mobility edge \( P(M_j) \) acquires a scale-invariant form and this is indeed what was observed\(^ {15}\). The same data demonstrate that \( P(M_j) \) decreases fast for atypically extended states, i.e. at \( M_j/M \ll 1 \). This allows us to use the average value

\[
\mathcal{M} = L_0^{-3}(L_{\text{loc}}/L_0)^{-D_2},
\]

where \( L_0 \) is the short-scale cutoff length of the fractal behaviour. The associated energy scale \( E_0 = 1/(\nu_0 L_0^3) \) depends on the microscopic details of the model of disorder and might be small compared to Fermi-energy \( E_F \). Localization length depends on Fermi-energy (in the scaling region \( L_{\text{loc}} \gg L_0 \)) as \( L_{\text{loc}} \approx L_0(E_0/(E_m - E_F))^\nu \), where \( E_m \) is the position of the mobility edge and \( \nu \) is the localization length exponent.

Another important property of the nearly-critical eigenstates is their strong correlation in energy and real space\(^ {13,15,17}\) even in the limit of strong fractality \( D_2 \ll 3 \). It results in the scaling dependence of the average matrix elements \( \overline{M_{jk}} \) on the energy difference \( E_j - E_k = \omega \):

\[
\mathcal{V} \overline{M_{jk}} = M(\omega) \approx \left\{ \begin{array}{ll}
(\omega_0/\omega)^{\gamma} & \text{at } \omega < \delta_L \\
(\omega_0/\omega)^{\gamma} & \text{at } \delta_L < \omega \leq E_0
\end{array} \right.
\]

(3)

where \( \gamma = 1 - D_2/3 \) and \( \mathcal{V} \) is the total system’s volume. Note that in the critical region \( M(\omega) \gg 1 \) in contrast to both a metal and to a deep insulator \( L_{\text{loc}} \sim L_0 \), where \( M(\omega) \approx 1 \).

We begin with the insulating region \( \Delta \ll \delta_L \ll \omega_D \) where Cooper interaction can be treated perturbatively. In the first order of the perturbation theory we take into account only diagonal terms \( j = k \) of the interaction similar to the case of ultrasmall grain\(^ {14}\). Then the energy (counted from \( E_F \)) required to break a bound pair of electrons sitting in the \( j \)-th orbital state is \( 2\Delta_0^{(j)} = \Delta_{\text{loc}}^{(j)} \).

Typical value of this “parity gap” (cf.\(^ {11}\)) scales then as

\[
\Delta_0^{(j)} = \frac{\lambda}{2} E_0 \left( \frac{L_0}{L_{\text{loc}}} \right)^{D_2} \propto (E_m - E_F)^{\nu D_2}
\]

(4)

Neglecting fluctuations in the local values of \( \Delta_0^{(j)} \), one finds that all states occupied by single electrons are shifted up by the amount \( \Delta_0 \) that leads to the electron DoS \( \nu(\varepsilon) = \nu_0 \delta(\varepsilon - \Delta_0) \). In fact, values of \( \Delta_0^{(j)} \) differ for different localized states, and the average density of states \( \nu_0(\varepsilon) \) in a large sample is determined by the IPR distribution \( P(M) \):

\[
\nu(\varepsilon) = \nu_0 \int_0^{2\varepsilon/\nu_0} \mathcal{P}(M)dM
\]

(5)

As mentioned above, numerical data for \( P(M) \) indicate its very fast decrease at \( M/M \rightarrow 0 \). Thus the DoS shape\(^ {10}\) is not far from a rectangular sharp gap, with the gap value given by Eq.\(^ {11}\). We emphasize that i) parity gap \( \Delta_0 \) is much larger than level spacing \( \delta_L \) at \( L_{\text{loc}}/L_0 > (2/\lambda)^{D_2/2} \), and ii) the DoS\(^ {10} \) does not contain any “coherence peak” above the gap (cf. Ref.\(^ {18}\)).

We associated\(^ {10} \) the spectral gap \( \Delta_0 \) with the measured\(^ {10} \) activation energy \( T_i \). The external parameter \( (E_m - E_F) \) representing the disorder strength in Eq.\(^ {11}\) can be replaced with an experimentally more accessible parameter \( (\sigma - \sigma) \propto (E_m - E_F) \). Here \( \sigma \) is the high temperature conductivity and \( \sigma \) is the value of the conductivity where the parity gap \( \Delta_0 \) first develops. We obtain

\[
T_i = A(1 - \sigma/\sigma)^{\nu D_2},
\]

(6)

where \( A \) is conductivity-independent. This equation predicts a moderate increase of \( T_i \) with disorder strength in agreement with the experimental data, see Fig.\(^ {2}\).
Fermi level is very close to the Anderson mobility edge $E_m$ and level spacing $\delta L \sim E_0((E_m - E_F)/E_0)^{3\nu}$ is negligibly small, one can use for $M(\omega)$ the Eq. (3). Then the critical temperature is given by

$$T_c(\lambda, \gamma) = E_0 \lambda^{1/\gamma} C(\gamma)$$

(8)

where dimensionless function $C(\gamma)$ can be computed numerically. At small $\lambda$ the value given by Eq. (8) exceeds the BCS value $T_{BCS} \sim \omega_F \epsilon^{-1/\lambda}$. This may lead to a maximum in $T_c$ near the critical disorder. The zero-temperature energy gap $\Delta_T = 0$ in the same limit $\delta L \to 0$ is given by Eq. (5) with $C(\gamma)$ replaced by another function $D(\gamma)$. We plot $D(\gamma)$ and $2D(\gamma)/C(\gamma)$ in Figure 3.

![Figure 3: Suppression of $T_c$ as a function of level spacing within localization volume. The insert shows $\gamma$-dependence of the dimensionless gap and $2 \Delta(0)/T_c$ for $\delta L = 0$.](image-url)

Using Eq. (8) we eliminate the interaction constant $\lambda$ and the cut-off parameter $E_0$ from Eq. (8) and arrive at

$$\Delta_P = \frac{1}{2(C(\gamma))^2} \cdot \delta L \left( \frac{T_c(\delta L)}{T_c(0)} \right)^\gamma$$

(9)

This formula (applied at $\delta L \gg T_c^0$) generalizes the results of Ref. [11] to bulk strongly disordered superconductors. Note that in contrast to Ref. [11] here the reduction of $\Delta_P$ compared to $\delta L$ is not due to the renormalization of attractive interactions, but to the enhancement of the matrix elements $M_j$ due to fractality.

To study the effect of nonzero level spacing $\delta L$ upon $T_c$, we approximate $M(\omega)$ by a simple interpolation formula $M(\omega) = E_0(\omega^2 + \delta L^2)^{-1/2}$ and solve the linearized version of Eq. (6) for $T_c(\delta L)$ numerically. Since $M(\omega)$ is a uniform function of $T$ and $\delta L$, while the coupling constant $\lambda$ and $E_0$ enter Eq. (6) only in a combination $E_0 \lambda^{1/\gamma}$, it is possible to present the dependence $T_c(\delta L)$ in the form

$$T_c(\delta L) = T_c^0 T_c \left( \frac{T_c^0}{\delta L} \right)$$

(10)

where $T_c^0$ is defined in Eq. (5) and scaling function $T_c(x)$ does not depend on $\lambda$. This universal function was found numerically for $\gamma = 0.57$ (corresponding to $D_2 = 1.3$ for 3D Anderson transition), the result is plotted in Fig. 3. Function $T_c(x)$ approaches very slowly its asymptotic value 1 at large $x$. The actual $T_c$ is suppressed as compared to the mean-field result due to fluctuations. We

We now turn to the parameter region $\delta L \ll \Delta$ where one expects a global superconductive coherence to exist at low enough $T$. Indeed, in this regime a given localized single-particle state typically overlaps in a real space with a large number $\sim \Delta/\delta L$ of eigenstates $\psi_j$ in the same energy strip $|\epsilon_j| \leq \Delta$. It is natural to expect that in this case the mean-field approximation can provide some guidance, while not necessarily being quantitatively accurate (e.g. because fluctuations of the overlap matrix elements $M_{ij}$ are large). To test the validity of the mean-field, we compared its prediction for $T_c(\delta L)$ (see below) with the transition temperature that was found numerically by computing the first terms of the virial expansion applied to [11] with $M_{ij}$ determined by exact diagonalization. Reasonably good agreement was found.

To proceed with the mean-field analysis, we introduce averaged energy-dependent pairing amplitudes $F(\epsilon_j) = (\epsilon_j^2 + \epsilon_j^2)$ and the gap function $\Delta(\epsilon) = \epsilon \int d\epsilon_M(\epsilon - \epsilon_1)F(\epsilon_1)$. Following standard steps, we decouple interaction term in the Hamiltonian (6) via the gap function $\Delta(\epsilon)$, calculate anomalous averages $F(\epsilon)$, and arrive at the modified BSC gap equation in the form

$$\Delta(\epsilon) = \frac{\lambda}{2} \int^{-\infty}_{+\infty} d\epsilon_1 \frac{M(\epsilon - \epsilon_1)\Delta(\epsilon_1)}{\sqrt{\epsilon_1^2 + \Delta^2(\epsilon_1)}} \tan h \frac{\sqrt{\epsilon_1^2 + \Delta^2(\epsilon_1)}}{2\epsilon}$$

(7)

Gap function $\Delta(\epsilon)$ obeys Eq. (7) is an even function of $\epsilon$ with the maximum value $\Delta(\epsilon = 0) = \Delta_0$.

Superconducting transition temperature $T_c$ is determined by linearization of Eq. (6) with respect to $\Delta(\epsilon)$. Due to power-law decrease of $M(\omega)$ at large arguments, the integral in Eq. (6) converges and is dominated by $\omega \sim T$ ("infrared superconductivity") so no upper cut-off is needed, contrary to usual BSC problem. When

FIG. 2: Experimental values of the gap from Ref.[1] $T_f$ (boxes) and a fit to the equation [12] with $\nu = 1, D_2 = 1.3$. The only fitting parameter was the constant $A = (\lambda/2)E_0$; the data points of Ref.[1] correspond to $E_0 \approx 100K$ at $\lambda \approx 0.2$ extracted from the BCS value of $T_c \approx 3K$ for less disordered samples and $\omega_f \approx 500K$. The value of $\sigma_x$ was determined from high $T$ data. Application of scaling formulas to the data shown here is justified by the large value of the localization length of the most disordered sample shown in this plot: $\lambda_{max} > 30Å$ which was deduced from the Mott temperature characterizing the resistivity of similar samples at intermediate temperatures.

FIG. 3: Suppression of $T_c$ as a function of level spacing within localization volume. The insert shows $\gamma$-dependence of the dimensionless gap and $2 \Delta(0)/T_c$ for $\delta L = 0$. Using Eq. (8) we eliminate the interaction constant $\lambda$ and the cut-off parameter $E_0$ from Eq. (8) and arrive at

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where $T_c^0$ is defined in Eq. (5) and scaling function $T_c(x)$ does not depend on $\lambda$. This universal function was found numerically for $\gamma = 0.57$ (corresponding to $D_2 = 1.3$ for 3D Anderson transition), the result is plotted in Fig. 3. Function $T_c(x)$ approaches very slowly its asymptotic value 1 at large $x$. The actual $T_c$ is suppressed as compared to the mean-field result due to fluctuations. We
expect it still has the scaling form \(10\) but with a different scaling function that vanishes at \(x \leq x_c \sim 1\).

Although gap equation \(7\) is similar to the conventional one, the real-space properties of the state that it describes are unusual. The local pairing amplitude \(F(r) = \sum_j (c_j^T c_j^\dagger) \psi_j^2(r)\) is extremely inhomogeneous in space, populating only small fraction \(\propto \{T_c/\delta_f\}/E_0\)\(^2\) of the total volume\(^2\). Diamagnetic response of such a superconductor differs strongly from that of usual "dirty-limit" materials with uniform \(|F(r)|\), but reminds that of weakly-coupled Josephson junctions arrays. Qualitatively, we expect (i) extremely weak Meissner effect and considerably stronger linear shielding effect, and (ii) superconductor \(\rightarrow\) gaped insulator \(\rightarrow\) Mott insulator sequence of transitions upon magnetic field increase.

We now discuss the effect of "local pairing" considered previously (see Eq. \(4\)) in insulating state. We have seen that in insulating state when \(\delta_L \gg \Delta\) single-particle excitation (carrying spin \(\frac{1}{2}\)) have a gap \(\Delta_1 = \Delta_p\). On the other hand, excitations that involve only hopping of paired electrons between localized levels and do not involve breaking pairs are gapless, i.e. the energy gap \(\Delta\) vanishes. On the other hand, in superconducting state within mean-field we have \(\Delta_1 = \Delta + \Delta_p\), while the gap for pair excitations\(^2\) (without pair breaking) is \(\Delta_2 = 2\Delta\). When \(L_{loc}/L_0 \rightarrow \infty\), the parity gap \(\Delta_p\) becomes much larger than \(\delta_L\). Therefore, we expect that there is a regime \(\Delta_p \gtrsim \Delta\) in the superconducting state where the spin gap \(\Delta_1\) is larger than the energy gap.

In the derivation of the gap equation \(7\) we neglected the effects of the parity gap which is correct for \(\Delta_p \ll T\). A large parity gap pushes up all unpaired states with respect to the paired ones which modifies the thermal distribution factor in the Eq. \(4\). In the limit of \(\Delta_p \gg T\) the contribution of unpaired electrons vanishes and the problem becomes equivalent to the pseudospin model\(^2\). This effectively doubles the value of the gap and thus replaces \(2T \rightarrow T\) in the \(\tanh()\) argument in \(7\). For a general \(\Delta_p \sim T\) the crossover formula between these two regimes can be obtained in the mean-field approximation. The result is (for derivation cf.\(^2\)) that \(\tanh()\) in Eq. \(7\) should be replaced by

\[
F(\epsilon_1, T) = \frac{\sinh B}{\cosh B + e^{-\Delta_p/T}},
\]

where \(B = \sqrt{\epsilon_1^2 + \Delta^2(\epsilon_1)/T}\), and \(\epsilon_1 = \epsilon - \Delta_p\).

Note that most experiments probing the superconducting gap, such as tunnelling conductance, optical conductivity or NMR, measure the spin gap \(\Delta_1\). One thus expects that these data would show anomalously large (compared to BCS value of 1.76) ratio of spectral gap to transition temperature near \(T_s\) transition. An additional suppression of \(T_c\) in comparison with \(\Delta_p\) can be due to electron-electron interaction in the density channel (irrespectively of the interaction sign) that is not included into the model \(4\).

An anomalously large ratio \(\Delta_1/T_c\) leads to the insulating trend of the resistivity versus temperature behavior in the intermediate temperature range \(T_c < T < \Delta_1\). This was observed in strongly disordered superconductors, see e.g.\(^2\), and is well-known as a pseudogap phenomenon in underdoped cuprates\(^2\). The quantitative similarity between \(R(T, B)\) behavior in \(\text{InO}_x\) films and underdoped cuprates\(^2\) allows one to speculate that the pseudogap in underdoped cuprates might also be related to pairing of electrons on localized states. The important difference of the cuprates is the d-wave symmetry of the pairing.

The effect of magnetic field on the fractal superconductor deserves a separate study. Here we only remark that the available magnetoresistance data are not consistent with a conventional picture of weakly coupled superconducting grains but point towards fractal superconductivity. Typically\(^2, 5, 6, 7, 9\), the magnetoresistance is negative at high fields and persists up to very high field values. Consider, for instance, the magnetoresistance data on amorphous \(\text{InO}_x\) shown in Fig.1 of Ref. \(4\) (the samples were of the same type and origin as those studied in\(^2\)). Activation energy for this sample is \(T_I \approx 15\) K (as determined in the temperature range \(1.3 - 5\) Kelvin); application of 8 Tesla field at the temperature of 1.58 K leads to negative magnetoresistance (MR) \(R(H = 8T)/R(0) \approx 0.5\). Interpreting this effect as being due to magnetic-field suppression of the pairing gap in small grains, we find \(T_I - T_f(H = 8T) \approx 0.7\) K, which is about 5\% of \(T_I\) only. Interpolating this dependence we estimate that the field necessary to destroy completely superconductivity in each grain is about \(H_{cg}^{exp} \sim 50 - 80\) Tesla which is much larger than one expects for realistic grain sizes. Indeed, orbital critical magnetic field for a small (radius \(R < \xi\), where \(\xi = \sqrt{R/\Delta}\) is the coherence length) superconductive grain is equal\(^2, 27\) to \(H_{cg}^{est} \approx \frac{1000 T}{R^2}\) where \(R\) and \(\xi\) are measured in nanometers. Using a typical diffusion constant for a poor metal \(D \approx 1\) cm\(^2\)/s and allowing for a very high gap value \(\Delta = 10\) K, we find \(\xi = 8.5\) nm. Together with the lowest bound for the grain radius \(R = 6\) nm (determined via the condition \(\Delta = (4\pi\epsilon_0 R^3)\sim \Delta < \Delta)\) it leads to \(H_{cg}^{est} \leq 20\) Tesla which is still much smaller than \(H_{cg}^{exp}\) above. These estimates did not take into account the spin effect of magnetic field that would further decrease the value of \(H_{cg}^{est}\). This discrepancy between the experimental value of the characteristic magnetic field, \(H_{cg}^{exp}\), and its estimate, \(H_{cg}^{est}\), can be accounted for by the fractal nature of the paired states because the effect of magnetic field on the local pairing gets much weaker due to low value of fractal dimension \(D_2 = 1.3\) of these states that indicates "almost one-dimensional" nature of these eigenstates.

In conclusion, weak Anderson insulators with Cooper attraction are shown to possess hard insulating gap whose magnitude is determined by the IPR statistics near the mobility edge. Although this gap is due to electron pairing, it does not lead to a coherence peak. In the ground-state of this insulator all electrons are paired on individual localized eigenfunctions. When the Fermi-level gets closer to the mobility edge, superconductive correlations develop between localized pairs. Key
features of the predicted superconductive ground-state are extreme inhomogeneity of superconductive correlations in real space, an unusually large (compared to \(T_c\)) single-particle excitation gap (spin gap), and pseudogaped regime at temperatures about \(T_c\). All these unusual features are due to the fractal nature of localized eigenstates near the mobility edge.

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