Low-energy quasiparticle excitations in dirty d-wave superconductors and the Bogoliubov-de Gennes kicked rotator

İ. Adagideli¹ and Ph. Jacquod²

¹Instituut-Lorentz, Universiteit Leiden, P.O. Box 9506, 2300 RA Leiden, The Netherlands
²Département de Physique Théorique, Université de Genève, CH-1211 Genève 4, Switzerland

(Dated: November 8, 2018)

We investigate the quasiparticle density of states in disordered d-wave superconductors. By constructing a quantum map describing the quasiparticle dynamics in such a medium, we explore deviations of the density of states from its universal form $\rho(E)$, and show that additional low-energy quasiparticle states exist provided (i) the range of the impurity potential is much larger than the Fermi wavelength $|\Delta| > \lambda_F$; (ii) classical trajectories exist along which the pair-potential changes sign; and (iii) the diffusive scattering length is longer than the superconducting coherence length. In the classically chaotic regime, universal random matrix theory behavior is restored by quantum dynamical diffraction which shifts the low energy states away from zero energy, and the quasiparticle density of states exhibits a linear pseudogap below an energy threshold $E^* \ll \Delta_0$.

PACS numbers: 71.23.-k, 74.72.-h, 05.45.Mt

In recent years, considerable attention has been focused on the low-energy properties of the quasiparticle spectrum of disordered cuprate superconductors [1,2]. Because many of the cuprate superconductors are randomly chemically doped insulators and disorder is a pair-breaker for d-wave superconductors, the role of nonmagnetic impurities is particularly important for an understanding of the d-wave superconducting state, its quasiparticle spectral and transport properties. Of special interest is the low-energy behavior of the single-particle Density of States (DoS) $\rho(E)$.

In early work [3], the self consistent T-matrix approximation was shown to break down for 2-dimensional d-wave superconductors. This led to a series of papers using nonperturbative methods, which predicted (at first sight) contradictory results: vanishing [3,2,2,8], constant [7,8], and diverging [9,10] DoS as $E \to 0$. On the numerical side, several investigations also predicted both vanishing [11,12], and diverging [13,14] DoS at zero energy. It was soon argued, based on numerical analysis, that the reason behind these contradicting predictions is the fact that the microscopic details of disorder (i.e. details beyond the transport mean free path $\ell$, such as the density of scatterers or the correlation length $\zeta$ of the impurity potential) as well as the symmetries of the clean Hamiltonian matter both qualitatively and quantitatively [13,15] (see also [16]). An important feature shared by the numerics of Refs. [11,12,13,14,15] is that the disorder is introduced via isolated point-like scatterers. Long-wavelength disorder, which may arise due to chemical doping away from CuO$_2$ planes, or can be induced via ion radiation techniques [17] or via a STM tip [18], is thus ignored. Effects of long-wavelength disorder are expected to become dominant when the CuO$_2$ planes have almost no atomic disorder [19]; they are the main focus of this paper.

It has recently been realized, in the context of mesoscopic physics and weak localization, that $\zeta$ and $\ell$, together with the Fermi wavelength $\lambda_F$, define two classes of complex quantum systems: quantum disordered systems where $\lambda_F\ell/\zeta^2 > 1$ and quantum chaotic systems for which $\lambda_F\ell/\zeta^2 < 1$ [20]. The latter class is characterized by the emergence of a new diffusive scattering time scale, $\tau_E = \nu^{-1} \ln[1/\lambda_F]$, defined as the time it takes for the classically chaotic dynamics (with Lyapunov exponent $\nu$) to stretch a wavepacket of minimal initial extension $\lambda_F$ to a length $\zeta$. In contrast to quantum disordered systems, quantum chaotic systems exhibit nonuniversal properties due to their short-time classical (i.e. deterministic) dynamics. In particular, significant deviations from Random Matrix Theory (RMT) emerge, as was recently found by Adagideli et al. [10] in the context of impurities in d-wave superconductors. These authors used a semiclassical approach to calculate the low energy DoS for a collection of extended scatterers (a quantum chaotic system) and found an asymptotic behavior $\rho(E) \sim 1/\sqrt{E}$ [21,22] as $E \to 0$. They nevertheless argued that the RMT predictions of a linear pseudogap would be restored at lower energy $E < E^*$, i.e. that the singularity in the DoS would be cut off at an energy $E^*$ related to $\tau_E$, by diffusive (nonclassical) scattering occurring at larger times $\tau > \tau_E$. The purpose of the present paper is to investigate the modifications that the DoS undergoes as the correlation length of the impurity potential increases and $\tau_E$ becomes relevant. We will focus our attention on (i) providing for numerical checks of the theory of Adagideli et al. in the case of long-wavelength disorder [10]; (ii) finding out whether for some $E^*$ the DoS is suppressed for $E < E^*$, in agreement with RMT predictions [4,5,6], (this could also reconcile the above mentioned contradictory predictions); and (iii) investigate the transition region between extended and pointlike disorder.

We start by introducing a quantum map model for quasiparticle states in disordered d-wave superconduc-
tors. The main motivation behind this model is to investigate discrepancies (in low-energy DoS) between point-like vs. extended disorder as well as the transition region, i.e. the regime in which the impurity size is intermediate. To the best of our knowledge no disorder model which includes extended impurities as well as pointlike in d-wave superconductors has been studied numerically so far. This map has two additional advantages: First, as both the density and correlation length of impurities can be tuned independently, it is possible to interpolate between the two extreme regimes of strong disorder: unitary disorder (i.e. disorder due to dilute, pointlike scatterers, viz. quantum disorder) and quasi-classical disorder [10] (i.e. disorder due to extended scatterers, viz. quantum chaotic). Second, from a numerical point of view, it allows for the investigation of very large system sizes, i.e. lattice sizes of up to $256 \times 256$, which are necessary for both variations of the disorder correlation length and the numerical extraction of the parametric behavior of the DoS. Our reasons why a dynamical model is relevant are: (i) In absence of superconductivity, many properties of quasiparticles in disordered media (such as Anderson localization) are correctly described by 1-D maps [21]. In fact it has been shown by Altland and Zirnbauer that one of those maps, the 1-D kicked rotator, and quasi-one-dimensional metallic wires are described by the same effective field theory [22]. Recent numerical investigations in $D = 2$ suggest that this is also true in higher dimensions [23]. (ii) In presence of superconductivity, Andreev maps based on the kicked rotator have recently been shown to adequately describe quantum dots in contact with a superconductor [24].

We first briefly discuss generic properties of quantum maps for uncoupled quasiparticles. The dynamics corresponds to a succession of free propagations, interrupted by sudden kicks of period $\tau_0$, i.e. instantaneous perturbations. Quantum maps are conveniently represented by a unitary, Floquet operator $F$, giving the time-evolution after $p$ kicks as $u(p) = F^p u(0)$, for an initial wavefunction $u(0)$. The matrix $F$ has eigenvalues $\exp(-i \varepsilon_m)$, which define quasi-energies $\varepsilon_m \in (-\pi, \pi)$ (energies and quasi-energies are expressed in units of $\hbar/\tau_0$). While the energy is not conserved, the periodicity of the kick still preserves quasi-energies, much in the same way as a periodic potential breaks translational symmetry, but still preserves quasi-momentum. Time evolution of hole excitations (being the time-reversed of electronic excitations) is given by $v(p) = (F^*)^p v(0)$. Specializing to the $D$-dimensional kicked rotator, we write the Floquet operator as [27]

$$F = \exp\left(-i \frac{K I}{\hbar \tau_0} \prod_{j=1}^D \cos r_j \right) \exp\left(i \frac{\hbar \tau_0}{2T} \nabla^2 \right).$$

It describes the free motion of a particle with dimensionless coordinates $\{r_j\}$ (e.g. expressed in units of a lattice constant), which is interrupted at periodic time intervals $\tau_0$ by a kick of strength $K \prod_{j=1}^D \cos r_j$. $I$ is the moment of inertia of the particle, and $K$ is the kicking strength. For $D = 1$ and 2, increasing $K$ makes the classical dynamics evolve from integrable ($K = 0$) to fully chaotic [$K \gtrsim 7$, with Lyapunov exponent $\lambda \approx \ln(K/2)$]. For $0 < K < 7$ stable and unstable motion coexist (a so-called mixed phase space) [25]. Increasing $K$ is thus tantamount to increasing the amount of disorder, the fully chaotic regime corresponding to a finite density of impurities.

Electron and hole excitations inside a superconductor are however coupled by a nonvanishing pair-potential. Accordingly we extend the kicked rotator of Eq. (1) to a Bogoliubov-de Gennes form. We discuss this construction for the case $D = 2$. First, we replace the free quasiparticle motion by a coupled electron and hole dynamics,

$$\mathcal{F}_0 = \exp(-i\mathcal{H}\tau_0/\hbar),$$

$$\mathcal{H} = \mathcal{H}_0 + \Delta \sigma_z.$$ (2b)

Here, $H = -(\hbar^2 \nabla^2 / 2I) - E_F$, with $E_F$ the Fermi energy, $\sigma_x, \sigma_z$ are Pauli matrices acting in particle-hole space, and $\Delta$ is the superconducting pair potential. Second, the coupled quasiparticle motion is followed by a kick

$$\mathcal{F}_K = \exp(-i\mathcal{H}_K/\hbar),$$

$$\mathcal{H}_K = \frac{KI}{\tau_0} \cos x \cos y \sigma_z.$$ (3b)

Exponentiating the Pauli matrices, we end up with the Bogoliubov-de Gennes-Floquet (BdGF) operator

$$\mathcal{F} = \mathcal{F}_K \mathcal{F}_0,$$ (4a)

$$\mathcal{F}_0 = \cos \left(\frac{(H^2 + \Delta^2)(\tau_0/\hbar)^2}{I} \right) + i \sin \left(\frac{(H^2 + \Delta^2)(\tau_0/\hbar)^2}{\sqrt{H^2 + \Delta^2}} \right) [H\sigma_z + \Delta \sigma_x],$$ (4b)

$$\mathcal{F}_K = \cos \left(\frac{KI}{\hbar \tau_0} \cos x \cos y \right) + i \sin \left(\frac{KI}{\hbar \tau_0} \cos x \cos y \right) \sigma_z,$$ (4c)

with $\mathcal{I}$, the identity matrix in particle-hole space. For $\Delta = 0$, Eq. (4) describes uncoupled electron and hole excitations in a disordered 2D metal. Once this metal becomes superconducting, $\Delta$ couples these excitations during their free propagation, while it is neglected during the instantaneous kick. As in the case of a BdG eigenproblem, the $2M$ quasienergies [with average spacing $\delta \equiv \langle \varepsilon_{m+1} - \varepsilon_m \rangle = \pi/[M]$ of the BdGF equation $\mathcal{F}\phi_m = \exp(-i\varepsilon_m)\phi_m$, come in pairs with opposite sign $\varepsilon_m = -\varepsilon_{2M-m+1}$, similarly to the spectral properties of a BdG Hamiltonian. These considerations establish the correspondence between the map of Eq. (1) and quasi-particles in a dirty superconductor.

We next quantize the phase space on a 4-torus $\{x,y,p_x,p_y\}$, with dimensionless momentum $p_{x,y} = $
Fourier transform between real space and momentum co-
icically extract the quasienergy DoS from the eigenvalues
\[ L \text{ integer} \]
\[ M \]
\[ \zeta \]
\[ U \]
\[ \text{extended s-wave pair potentials away from half filling} \]
diagonalize using the Lanczos algorithm [26].
\[ E/\hbar \]
\text{The gap singularity at} \( E = \Delta_0 \) is provided in Fig. 2. It indicates that the characteristic semiclassical singularity exhibited by the DoS as \( E \rightarrow 0 \) is cut off at an energy \( E^* \ll \Delta_0 \), where a sharp drop occurs and \( \rho(E) \rightarrow 0 \). RMT predicts such a drop to occur over an energy scale given by the Thouless energy [2], which in our case is however significantly larger than \( \Delta_0 \) [28]. We thus attribute this drop to the emergence of diffractive scattering at times larger than \( \tau_\text{F} \) as follows. According to Ref. [10], the DoS corresponding to low energy semiclassical states can be estimated from a mapping onto a tight-binding chain with random hoppings, for which the eigenfunctions are localized with an energy-dependent localization length \( \xi(E) \) [28]. At low energies, \( \xi \) exceeds the diffusive scattering length \( v_F \tau_\text{F} \), in which case hoppings between otherwise uncoupled tight-binding chains (corresponding to different classical trajectories) have to be taken into account. The emergence of these processes signals the breakdown of semiclassics and the restoration of RMT. One thus expects the vanishing of the DoS below a threshold energy given by the condition \( \xi(E^*) \approx v_F \tau_\text{F} \).

In Fig. 1 we show the quasiequity DoS for d-wave and extended s-wave pair potentials away from half filling (\( E_F = 2\pi^2/5 < \pi^2/2 \)), as the kicking strength increases. In the clean case (\( K = 0 \)) the two DoS are the same. The gap singularity at \( E/\Delta_0 = 1 \) gets washed out as \( K \) increases in both cases, however, a peak emerges in the d-wave DoS around \( E = 0 \), while \( \rho(E) = 0 \) in the extended s-wave case. This is in agreement with Ref. [10], i.e. the existence of low-energy states requires a change in the sign of the pair potential. In the extended s-wave case, the low energy peak is shifted by an energy corresponding to the gap averaged over all momenta mixed by the impurity potential.

We focus on the d-wave symmetry from now on. A closer look at the DoS in the fully chaotic regime with \( K = 8 \), is provided in Fig. 2. It indicates that the characteristic semiclassical singularity exhibited by the DoS as \( E \rightarrow 0 \) is cut off at an energy \( E^* \ll \Delta_0 \), where a sharp drop occurs and \( \rho(E) \rightarrow 0 \). RMT predicts such a drop to occur over an energy scale given by the Thouless energy [2], which in our case is however significantly larger than \( \Delta_0 \) [28]. We thus attribute this drop to the emergence of diffractive scattering at times larger than \( \tau_\text{F} \) as follows. According to Ref. [10], the DoS corresponding to low energy semiclassical states can be estimated from a mapping onto a tight-binding chain with random hoppings, for which the eigenfunctions are localized with an energy-dependent localization length \( \xi(E) \) [28]. At low energies, \( \xi \) exceeds the diffusive scattering length \( v_F \tau_\text{F} \), in which case hoppings between otherwise uncoupled tight-binding chains (corresponding to different classical trajectories) have to be taken into account. The emergence of these processes signals the breakdown of semiclassics and the restoration of RMT. One thus expects the vanishing of the DoS below a threshold energy given by the condition \( \xi(E^*) \approx v_F \tau_\text{F} \).

In the inset to Fig. 2 we show the asymptotic behavior of the DoS on a log-log scale. Once abstraction is made of the drop in the DoS below \( E^* \), the semiclassical data exhibit a singular behavior slightly below \( \propto E^{-1} \).
FIG. 3: Low energy DoS for the d-wave Bogoliubov-de Gennes kicked rotator with decreasing disorder range as defined in Eqs. [2], [4], and [6], with $L_x \times L_y = 256 \times 256$, $\Delta_0 = 0.4$, $E_F = 2\pi^2/5$, $K = 8$, and $N_H = 1, 3, 5, 7, 17, \text{ and } 27$ (from top to bottom).

(black circles) which is in qualitative agreement with the prediction $\rho(E) \propto E^{-1} \ln|E|^{-3}$ of Ref. [10].

Having established the validity of semiclassical predictions in the quantum chaotic regime, we next decrease the range of the disorder and enter the quantum disordered regime. We accomplish this via the inclusion of higher harmonics to the kicking potential, and replace Eq. [6] by

$$\mathcal{H}_K = \frac{KI}{N_H^2 \tau_0} \sum_{l,m=1}^{N_H} \cos[lx] \cos[my] \sigma_z. \quad (6)$$

The typical impurity size decreases as $\zeta \propto N_H^{-1}$. Fig. 3 shows the disappearance of the low-energy peak in the DoS as $N_H$ increases. For the set of parameter considered, once $N_H \approx 27$ is reached, the DoS vanishes at $E = 0$. Note that the resolution used in Fig. 3 does not allow to see the opening of the RMT gap below $E^*$. A more precise look at the DoS for $N_H = 27$ is provided on the inset to Fig. 2 (empty diamonds). The data clearly indicate the expected RMT linear suppression of the DoS.

Our results thus clarify the competition between RMT and semiclassics. The next step is to investigate the transport properties and to explore the parametric dependence of $E^*$. Work along those lines is in progress.

This work was supported by the Dutch Science Foundation NWO/FOM and the Swiss National Science Foundation. We thank I. Gornyi, A. Yashenkin, M. Vojta, N. Trivedi, and P. M. Goldbart for interesting discussions and comments.

[26] This diagonalization is performed with only $O(M^2 \ln M)$ operations, if the multiplication with $U$ in Eq. [6] is performed with the Fast-Fourier-Transform algorithm, see: R. Ketzmerick, K. Kruse, and T. Geisel, Physica D 131, 247 (1999).
[28] Quantizing the kicked rotator on a torus corresponds to a ballistic chaotic system, so that the Thouless energy is of the order of the size of the band.