Mesoscopic Effects in Disordered Superconductors near H_{c2}

B. Spivak and Fei Zhou
Physics Department, University of Washington, Seattle, Washington 98195
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A theory of disordered superconductors near the upper critical field H_{c2}^0 is presented. It is shown that the magnetic field dependences of critical temperatures $T_c(H)$ in mesoscopic samples exhibit multiple reentrant superconductor–normal-metal transitions. In the framework of mean field theory T_c of bulk samples remains finite at any H and asymptotically approaches zero at large H. This means that at zero temperature the upper critical magnetic field does not exist. The role of quantum fluctuations of the phase of the order parameter is discussed.

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The theory of the upper critical field H_{c2}^0 in disordered superconductors was developed a long time ago [1,2]. In these papers H_{c2}^0 was found as a magnetic field at which the linear equation

$$\Delta(\mathbf{r}) = \alpha \int d\mathbf{r}'\, K(\mathbf{r},\mathbf{r}')\Delta(\mathbf{r}') \qquad (1)$$

has solutions. Here $\Delta(\mathbf{r})$ is the superconducting order parameter

$$K(\mathbf{r},\mathbf{r}') = T \sum_{n=-\infty}^{\infty} G_\epsilon(\mathbf{r},\mathbf{r}')G_{-\epsilon}(\mathbf{r},\mathbf{r}'), \qquad (2)$$

$G_\epsilon(\mathbf{r},\mathbf{r}')$ is the Matsubara exact electron Green's function [1] in the presence of the elastically scattering potential, $\epsilon = (2n+1)\pi T$, n is an integer number, α is the electron-electron interaction constant, and T is the temperature. Averaging over realizations of random potential leads to the equation

$$\langle\Delta(\mathbf{r})\rangle = \alpha \int d\mathbf{r}'\, \langle K(\mathbf{r},\mathbf{r}')\rangle\langle\Delta(\mathbf{r}')\rangle. \qquad (3)$$

Here angular brackets $\langle\,\rangle$ stand for averaging over realizations of random potential. To get Eq. (2) one has to assume the absence of correlations between the two random quantities $\Delta(\mathbf{r})$ and $K(\mathbf{r},\mathbf{r}')$. Quasiclassical treatment of magnetic field $\langle G_\epsilon(\mathbf{r},\mathbf{r}',H)\rangle = \langle G_\epsilon(\mathbf{r}-\mathbf{r}',0)\rangle \exp[(ie/c\hbar)\mathbf{A}(\mathbf{r})(\mathbf{r}-\mathbf{r}')]$ yields further simplification [1,2]. Here $\mathbf{A}(\mathbf{r})$ is the vector potential of the magnetic field. The justification for this procedure is that at $|T - T_c(0)| \ll T_c(0)$ and $H = 0$ the electron Green's function $\langle G_\epsilon(\mathbf{r}-\mathbf{r}',0)\rangle \sim \exp(-|\mathbf{r}-\mathbf{r}'|/l_0)$ decays exponentially as a function of $|\mathbf{r}-\mathbf{r}'|$ at $|\mathbf{r}-\mathbf{r}'| \gg l_0$, while $\mathbf{A}(\mathbf{r})$ changes over distances much larger than l_0. Here $l_0 = \min[v_F/T_c, l]$, $T_c(0)$ is the critical superconducting temperature at $H = 0$, v_F is the Fermi velocity, and l is the elastic electron mean free path.

In the case of clean samples with $l = \infty$ at $T = 0$ and $H = 0$ the electron Green's function decays relatively slowly (in the three-dimensional case it is inversely proportional to $|\mathbf{r}-\mathbf{r}'|$), which makes it impossible to treat the magnetic field quasiclassically and one has to take into account the Landau quantization [3–7]. This leads to the counterintuitive result [3,5–7] that at $T = 0$ superconducting solutions of the BCS equations exist at arbitrary high H.

On the other hand, in the case of disordered metals, at $T = 0$ the exponential decay of the averaged Green's function $\langle G_\epsilon(\mathbf{r}-\mathbf{r}')\rangle \sim \exp(-|\mathbf{r}-\mathbf{r}'|/l)$ originates from the fact that the phase of $G(\mathbf{r},\mathbf{r}')$ at $|\mathbf{r}-\mathbf{r}'| \gg l$ is random, while its modulus decays with the same power of $|\mathbf{r}-\mathbf{r}'|$ as in the clean case [8–11]. This implies that the above mentioned quasiclassical expression for the magnetic field dependence of the averaged Green's functions does not work well for the exact Green's function. The randomness in phase of $G_\epsilon(\mathbf{r},\mathbf{r}')$ leads to mesoscopic fluctuations of $\Delta(\mathbf{r})$. The mesoscopic fluctuations of superfluid density at $H = 0$, which originate from the fluctuations of $\Delta(\mathbf{r})$, have been calculated in [12–14]. In the case $p_F l/\hbar \gg 1$ they are small compared with the average superfluid density and give only small corrections to observable quantities.

In this paper we consider the mesoscopic effects in disordered superconductors at small temperatures $T \ll T_c(0)$ and under magnetic fields H close to H_{c2}^0. We show that, in this case, the picture of the superconductor–normal-metal transition is determined by the mesoscopic fluctuations. In small samples the mesoscopic fluctuations manifest themselves in multiple reentrant transitions between the superconducting and normal states. In this case the genuine upper critical field cannot be uniquely defined and there are many H_{c2}^i, which characterize these superconductor-metal and metal-superconductor transitions as a function of H (see Fig. 1). In the case of bulk samples at $T = 0$, due to the existence of the mesoscopic fluctuations, superconducting solutions of mean field theory exist in arbitrary H (see Fig. 2). This means that although $T_c(H)$ decreases with H it is never zero in the framework of mean field theory. To find the genuine H_{c2} at $T = 0$ one has to take into account the quantum fluctuations of the phase of the order parameter. The kernel $K(\mathbf{r},\mathbf{r}')$ is a random quantity, which depends on realizations of random potential. Although the "mesoscopic" part of the kernel $\delta K(\mathbf{r},\mathbf{r}') = [K(\mathbf{r},\mathbf{r}') - \langle K(\mathbf{r}-\mathbf{r}')\rangle] \ll K(\mathbf{r},\mathbf{r}')$ is relatively small, it nevertheless determines the properties of superconductors at $H \sim H_{c2}^0$.

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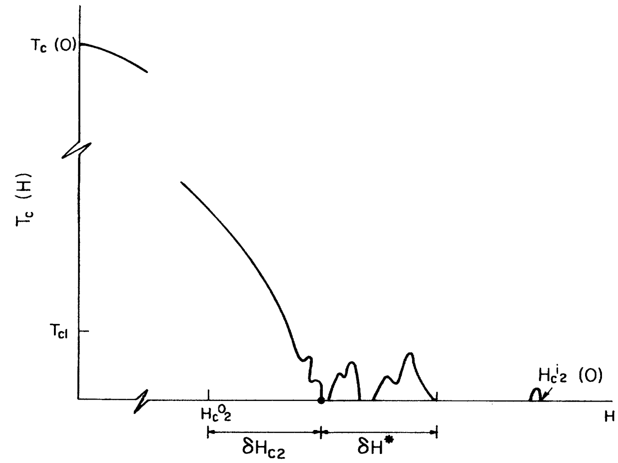

FIG. 1. Qualitative picture of $T_c(H)$ in mesoscopic samples.

Let us consider the case of a two-dimensional superconducting film with a thickness $L_z \ll \xi_0 = \sqrt{D/T_c(0)} \sim L_{H^0_{c2}}$. Here ξ_0 is the coherence length of a superconductor at $T = H = 0$, $L_{H^0_{c2}} = \sqrt{\pi c \hbar / e H^0_{c2}}$ is the magnetic length corresponding to H^0_{c2}, H^0_{c2} is the upper critical field for Eq. (3) calculated in Refs. [1,2], and $D = v_F l/3$ is the electron diffusion coefficient. Using perturbation theory for Eq. (1) one can estimate the correction to H^0_{c2} as follows ($\delta H_{c2} = H_{c2} - H^0_{c2}$):

$$\frac{\delta H_{c2}}{H^0_{c2}} = \int d\mathbf{r}\, d\mathbf{r}'\, \delta K(\mathbf{r}, \mathbf{r}', H^0_{c2} + \delta H_{c2}) \Delta_0(\mathbf{r}) \Delta_0(\mathbf{r}'), \quad (4)$$

$$\Delta_0(\mathbf{r}) = B \exp\left[-\left(\frac{|\mathbf{r} - \mathbf{r}_0|}{L_{H^0_{c2}}} \right)^2 \right]. \quad (5)$$

Here $\Delta_0(\mathbf{r})$ is the solution of Eq. (3), at $H = H^0_{c2}$, which describes a superconducting droplet of a size of the order of ξ_0, with the center at $\mathbf{r}_0$, and $B = \sqrt{1/L_z L_H^2}$. We define H_{c2} as the first point of superconductor-metal transition at $T = 0$ (see Fig. 1).

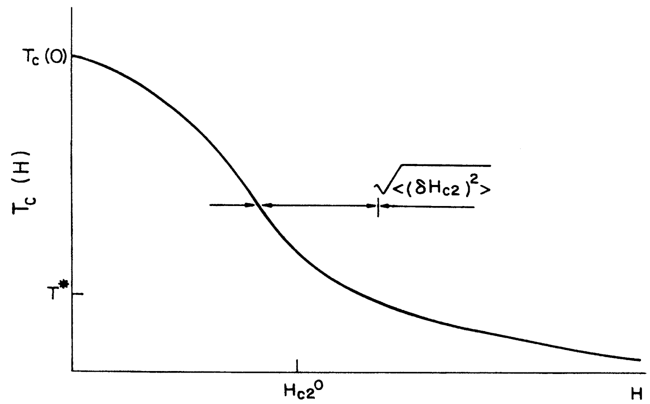

FIG. 2. Qualitative picture of $T_c(H)$ in bulk samples.

Thus the fluctuations of the H_{c2} can be expressed in terms of fluctuations of $G_\epsilon(\mathbf{r}, \mathbf{r}')$, and we can use the technique which has been developed for the description of mesoscopic fluctuations in normal metals [8,15–17]. In the first order of perturbation theory, with respect to $\hbar/p_F l$, the expression for the correlation function $\langle K(\mathbf{r}, \mathbf{r}') K(\mathbf{r}_1, \mathbf{r}'_1) \rangle$ is given by the diagrams in Figs. 3(a), 3(b), and 3(c) where solid lines correspond to $\langle G_\epsilon(\mathbf{r} - \mathbf{r}') \rangle$, dashed lines correspond to the correlation function of the scattering potential $(\pi/\tau\nu)\delta(\mathbf{r} - \mathbf{r}')$, $\tau = l/v_F$ is the elastic mean free path, and ν is the density of states in metal. The blocks of diagrams shown in Figs. 3(d) and 3(e) are called diffusion $D_\omega(\mathbf{r}, \mathbf{r}')$ and Cooperon $C_\omega(\mathbf{r}, \mathbf{r}')$, respectively. The equation for the Cooperon has the form [16]

$$\left\{ D\left(i\partial_{\mathbf{r}} + \frac{2e}{c} \mathbf{A}(\mathbf{r}) \right)^2 + i\omega \right\} C_\omega(\mathbf{r}, \mathbf{r}') = \frac{1}{\tau} \delta(\mathbf{r} - \mathbf{r}'). \quad (6)$$

The equation for $D_\omega(\mathbf{r}, \mathbf{r}')$ is the same as Eq. (6) with $\mathbf{A}(\mathbf{r}) = 0$. As a result we have

$$\begin{aligned} \frac{\langle (\delta H_{c2})^2 \rangle}{(H^0_{c2})^2} = \frac{T^2 \tau^6}{\nu^2} \sum_{n,n'=-\infty}^{\infty} \int & d\mathbf{r}\, d\mathbf{r}'\, d\mathbf{r}_1\, d\mathbf{r}'_1\, d\mathbf{r}_2\, d\mathbf{r}_3\, \Delta_0(\mathbf{r}) \Delta_0(\mathbf{r}') \Delta_0(\mathbf{r}_1) \Delta_0(\mathbf{r}'_1) \times [C_{\epsilon-\epsilon'}(\mathbf{r}_2, \mathbf{r}_3) C_{\epsilon'-\epsilon}(\mathbf{r}_2, \mathbf{r}_3) \\ & + D_{\epsilon-\epsilon'}(\mathbf{r}_2, \mathbf{r}_3) D_{\epsilon'-\epsilon}(\mathbf{r}_2, \mathbf{r}_3)] \times D\left[\left(i\partial_{\mathbf{r}_2} + \frac{2e}{c} \mathbf{A}(\mathbf{r}_2) \right) C_{2\epsilon}(\mathbf{r}, \mathbf{r}_2) \right] \left[\left(i\partial_{\mathbf{r}_2} + \frac{2e}{c} \mathbf{A}(\mathbf{r}_2) \right) C_{2\epsilon'}(\mathbf{r}_1, \mathbf{r}_2) \right] \\ & \times D\left[\left(i\partial_{\mathbf{r}_3} + \frac{2e}{c} \mathbf{A}(\mathbf{r}_3) \right) C_{2\epsilon}(\mathbf{r}', \mathbf{r}_3) \right] \left[\left(i\partial_{\mathbf{r}_3} + \frac{2e}{c} \mathbf{A}(\mathbf{r}_3) \right) C_{2\epsilon'}(\mathbf{r}'_1, \mathbf{r}_3) \right]. \quad (7) \end{aligned}$$

Using Eqs. (6) and (7) at $L_0 \gg \xi_0$ we get the estimate

$$\langle (\delta H_{c2})^2 \rangle / (H^0_{c2})^2 = \gamma (e^2/\hbar G_H)^2. \quad (8)$$

Here $L_0 = \min[L, L_T] = \sqrt{D/T}$, L is the sample size, G_H is the conductance of a sample of the size $L_{H^0_{c2}}$, which in the two-dimensional case does not depend on H^0_{c2} and of the order of $\sigma_0 L_z$, $\sigma_0 = e^2 \nu D$ is the Drude conductivity of the metal, and $\gamma \sim 1$. Equation (8) gives an estimate for fluctuations of H_{c2} between droplets which are separated by the distance larger than ξ_0. One can see it by calculating the correlation function at $|\mathbf{r} - \mathbf{r}'| \gg \xi_0$:

$$\langle \delta H_{c2}(\mathbf{r}_0) \delta H_{c2}(\mathbf{r}'_0) \rangle \sim \langle (\delta H_{c2})^2 \rangle \frac{(\xi_0)^4}{(|\mathbf{r}_0 - \mathbf{r}'_0|)^4} \times \exp\left(-\frac{|\mathbf{r}_0 - \mathbf{r}'_0|}{L_T} \right). \quad (9)$$

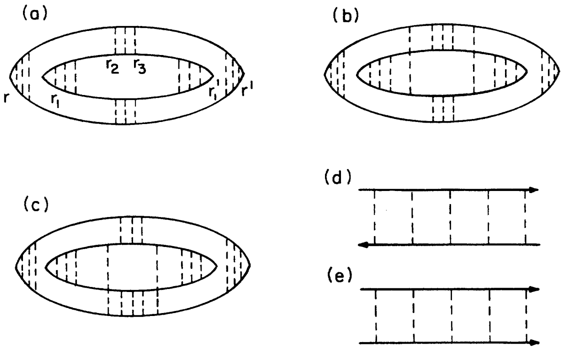

FIG. 3. Diagrams representing the correlation function $\langle \delta K(\mathbf{r},\mathbf{r}')\delta K(\mathbf{r},\mathbf{r}')\rangle$. Solid lines represent averaged electron Green's functions and dashed lines represent elastic scattering potential.

Since the expression on the right-hand side of Eq. (4) [which we will call $\delta k(H)$] is a random function of H, $T_c(H)$ is also a random function of H. This means that, generally speaking, the magnetic field dependence $T_c(H)$ corresponds to multiple superconductor–normal-metal transitions and there are many critical magnetic fields $H_{c2}^i(0)$ at $T = 0$. The dependence $T_c(H)$ is shown qualitatively in Fig. 1.

It is important that the dependence of $\delta k(H + \delta H)$ on δH contains random oscillations with all periods larger than Φ_0/L_0^2, $\Phi_0 = \hbar c/e$ being the flux quanta. One can see this, for example, from the power law decay of the correlation function as a function of δH at $L_H \ll L_{\delta H} \ll L_0$:

$$\left(\frac{[\langle \delta k(H) - \delta k(H + \delta H)\rangle]^2}{\langle (\delta k)^2\rangle}\right)^{1/2} = \gamma_1 \frac{L_H}{L_{\delta H}} \ln\!\left(\frac{L_{\delta H}}{l}\right). \tag{10}$$

Here $L_{\delta H} = \sqrt{\Phi_0/\delta H}$ and $\gamma_1 \sim 1$. Equation (10) has the same δH dependence as the correlation function of conductances of a point contact between semi-infinite metals at different values of magnetic field [18].

The physical meaning of Eq. (10) is the following. The random oscillations of $\delta k(H + \delta H)$ as a function of δH are determined by the random interference of different quasiclassical paths. The contribution into $\delta k(H)$ of paths traveling on a distance of the order of L_* decays at $L_* \gg \xi_0$ as L_*^2, while the sensitivity of this contribution to the change of the magnetic field increases with L_* as $H^2L_*^4/\Phi_0^2$ at $H^2L_*^4/\Phi_0^2 \ll 1$ and saturates at $H^2L_*^4/\Phi_0^2 > 1$. As a result the main contribution into $\delta k(\delta H)$ comes from the paths with $L_* \sim L_{\delta H} \gg \xi_0$ [9,18].

Therefore in large enough samples and at small temperatures, when $L_0 \gg L_{\delta H^*} = \xi_0(\hbar G_H/e^2)$, there exists an interval of magnetic fields near H_{c2}^0:

$$\delta H^* = H_{c2}^0(e^2/\hbar G_H)^2, \tag{11}$$

where the reentrant transitions corresponding to $T_c > 0$ take place with the probability of the order of unity (see Fig. 1). Equation (10) implies also that at $\delta H < \delta H^*$ the function $k(H)$ can be viewed as the diffusion trajectory of a particle with coordinate $k(H)$ and H plays the role of time. In this case the corresponding diffusion coefficient is of the order of $\langle(\delta k)^2\rangle/H_{c2}^0 = (e^2/\hbar G_H)^2/H_{c2}^0$ and the mean free path of the order of $\sqrt{\langle(\delta k)^2\rangle}\,\xi_0/L_T \sim (e^2/\hbar G_H)\xi_0/L_T$. Using the concept of diffusion of the quantity $k(H)$ in H space we can estimate the amplitude of the order parameter Δ^* at $T \ll T_{c1}$ and critical temperature T_{c1} corresponding to a reentrant transition $\Delta^* \sim T_{c1} \sim T_c(0)\,(e^2/hG_H)^2$ (see Fig. 1). The typical width of a reentrant transition is of the order of δH^*. At $|H - H_{c2}^0| \gg \delta H^*$, the reentrant transitions are rare events which are described by the tail of the distribution function of $k(H)$. Thus we arrive at a picture of superconducting droplets embedded into a normal metal and connected by Josephson links with a critical current J_c.The characteristic distance between the droplets is of the order of

$$R_0(H) = \xi(0)[f(H)]^{-1/2}. \tag{12}$$

Here $f(H)$ is the probability of finding a superconducting droplet.

We restrict ourselves to the consideration of high enough magnetic fields, when $R_0 > \xi_0$. In this case the average critical current $\langle J_c\rangle \sim \exp(-R_0/L_H - R_0/L_T)$ decays exponentially. It is important, however, that this effect originates from the fact that signs of J_c are random while the typical value of J_c does not decay exponentially if $R_0 > L_H$. One can estimate the typical value of $J_c(H)$ by calculating the variance with the help of diagrams shown in Figs. 3(a), 3(b), and 3(c).

$$[\langle (J_c - \langle J_c\rangle)^2\rangle]^{1/2} \sim (eD/R_0^2)\exp(-R_0/L_T). \tag{13}$$

To arrive at Eq. (13) we used the fact that $\Delta^* \gg D/R_0^2$. Since the values of critical Josephson currents between different superconducting droplets $J_c(R_0)$ have random signs, the network of such droplets should demonstrate spin glass behavior.One can estimate the critical transition temperature between the normal and superconducting glass states $T_c(H)$ as a temperature, which is of the order of typical Josephson coupling energy $E_J = \hbar|J_c|/2e$:

$$T_c \sim T_c^0 f(H). \tag{14}$$

The important consequence of Eq. (14) is that at $T_c(H) < T^* = T_0e^2/\hbar G_H$ the curvature of $T_c(H)$ becomes positive, as opposed to the conventional theory [1,2], where the curvature is negative. The dependence $T_c(H)$ in bulk samples is shown qualitatively in Fig. 2.

To calculate $F(k,H)$ we used the method of nonlinear σ model which was developed in Ref. [19] and has been used in Refs. [20,21] for the calculation of the distribution function of mesoscopic fluctuations of conductance of small samples. Since our calculations are basically the same as that in Refs. [20,21], we only present the results of the calculations. In the opposite limit $|\delta k|/\langle k\rangle \gg (\hbar G_H/e^2)^{1/2}$ we have

$$f(H) \sim \mathrm{Erf}\left[\frac{1}{\sqrt{4u}}\ln\left(\frac{|H - H_{c2}^0|}{H_{c2}^0}\frac{L_{H_{c2}^0}^2}{l_2}\right)\right]. \qquad (15)$$

Here $\mathrm{Erf}(x)$ is the error function, $u = \ln(\sigma/\sigma_0)$, and σ is the sample conductivity, which is different from σ_0 due to weak localization corrections.

Results presented above for the two-dimensional case can also be applied to the three-dimensional case with the precision of the order of 1. The only differences with the two-dimensional case are that in the three-dimensional case $G_H = \sigma_0\xi_0$ and that the preexponential factor in Eq. (13) should be replaced for $(eD/R_0^2)\xi_0/R_0$, which, however, does not change Eq. (14) significantly.

It is important to note that at $T = 0$ even Zeeman splitting of spin energy levels at $\mu H \gg T_c(0)$ does not suppress superconductivity in disordered samples entirely. One can see this from the fact that the value of $\langle[k(H)]^2\rangle$ given by diagrams shown in Figs. 3(a), 3(b), and 3(c) at $\mu H \gg T_c(0)$ decreases only a factor of 2 compared with the $\mu(H) = 0$ value considered above. Here μ is the Bohr magneton. The physical explanation of this result is the following. In the case of pure samples, $l = \infty$, the Zeeman splitting leads to the necessity of pairing electrons on the Fermi surface with different moduli of momentums. The wave functions composing the pairs at large enough splitting exhibit significantly different spatial dependences which lead to the suppression of superconductivity. However, if the g factor of electrons is equal to 2 and, in the case of Landau quantization, it is possible to pair electrons on the Fermi surface with opposite spins on different Landau levels [6]. In disordered samples with arbitrary g factor it is only a question of probability of finding regions of the size ξ_0 where at $\mu H \gg T_c(0)$ the wave functions of electrons of the Fermi surface with opposite spins exhibit similar spatial dependences.

The conclusion that at $T = 0$ there are superconducting solutions at arbitrary H looks paradoxical. We think that a genuine $H_{c2}(T = 0)$ is determined by quantum fluctuations of the phase of order parameter. The above considered mesoscopic fluctuations of the order parameter should be also reflected in mesoscopic fluctuations of magnetization of small superconducting samples.

The positive curvature of $T_c(H)$ has been observed both in high temperature superconductors [22,23] and in conventional disordered superconductors [24]. We think that the theory presented above may be useful for interpretation of these data.

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