



Pressure induced superconductor quantum critical point in multi-band systems

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ABSTRACT

In multi-band superconductors as inter-metallic systems and heavy fermions, external pressure can reduce the critical temperature and eventually destroy superconductivity driving these systems to the normal state. In many cases this transition is continuous and is associated with a superconducting quantum critical point (SQCP). In this work we study a two-band superconductor in the presence of hybridization V . This one-body mixing term is due to the overlap of the different wave-functions. It can be tuned by external pressure and turns out as an important control parameter to study the phase diagram and the nature of the phase transitions. We use a BCS approximation and include both inter- and intra-band attractive interactions. For negligible inter-band interactions, as hybridization (pressure) increases we find a SQCP separating a superconductor from a normal state at a critical value of the hybridization V_c . We obtain the behavior of the electronic specific heat close to the SQCP and the shape of the critical line as V approaches V_c .

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1. Introduction

The study of asymmetric superconductivity, i.e., of superconductivity in systems where different types of quasi-particles coexist at a common Fermi surface has raised a lot of interest in the last years. This in part is due to the relevance of this problem for many different areas in physics. It arises in cold atomic systems with superfluid phases [1], in color superconductivity in the core of neutron stars [2–4] and in condensed matter physics [5]. Furthermore it is closely related to inhomogeneous superconductivity, as FFLO phases [6], since this is a possible ground state for asymmetric systems. In condensed matter, as inter-metallic materials, there are electrons from different orbitals at the Fermi surface. If these quasi-particles have different effective masses, or occur in distinct numbers per atom there is a natural mismatch of their Fermi wave-vectors.

Then, in multi-band systems, even in the absence of external magnetic fields, one has to consider the possibility of inhomogeneous superconductivity or other types of exotic ground states as gapless superconducting phases [7,8] or phase separation [9].

In this paper we focus on the problem of driving a multi-band superconductor to the normal state by applying external pressure. When this occurs continuously, this transition is associated with a superconducting quantum critical point (SQCP). The theories which have been proposed for the SQCP rely in general on the presence of disorder or magnetic impurities [11]. However, there

is no reason to expect that this should play a role in clean systems driven to the normal state by external pressure. Here we discuss a model in which mixing transfers electrons that participate in Cooper pairing to a normal band eventually destroying superconductivity. Since there is no dissipation in the normal band, superconductivity disappears due to a loss of coherence in the system. Our model considers two hybridized bands in the presence of inter- and intra-band attractive interactions. These interactions are competing and determine the nature of the zero temperature phase transitions to the normal state as hybridization (pressure) increases. We show that only when intra-band interactions are dominant this transition is continuous. Otherwise, it is of first order and accompanied by phase separation as usual in this case [5].

The problem of superconductivity in systems with overlapping bands has been treated originally by Suhl et al. [13]. The relevance of the different interactions [14] has been discussed in terms of an energy associated with the Fermi surface mismatch, $\Delta(\delta_{k_f}) = v_F \delta_{k_f}$ and the critical temperature, $k_B T_c$. Only when the former is much smaller than the latter, inter-band interactions become important. In our approach, the mismatch δ_{k_f} depends on hybridization and can be controlled by pressure.

2. Model and formalism

We consider a model with two types of quasi-particles, a and b , with an attractive inter-band interaction [14] $\tilde{g}$, an attractive intra-band interaction $\tilde{U}$ and a hybridization term $\tilde{V}$ that mixes

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different quasi-particles states [5]. This one-body mixing term V is related to the overlap of the wave-functions and can be tuned by external parameters, like pressure, allowing to explore the phase diagram and quantum phase transitions of the model. The Hamiltonian is given by

$$H = \sum_{k\sigma} \tilde{\epsilon}_k^a a_{k\sigma}^\dagger a_{k\sigma} + \sum_{k\sigma} \tilde{\epsilon}_k^b b_{k\sigma}^\dagger b_{k\sigma} + \tilde{g} \sum_{kk'\sigma} a_{k'\sigma}^\dagger b_{-k-\sigma}^\dagger b_{-k-\sigma} a_{k\sigma} + \tilde{U} \sum_{kk'\sigma} b_{k'\sigma}^\dagger b_{-k'-\sigma}^\dagger b_{-k-\sigma} b_{k\sigma} + \sum_{k\sigma} \tilde{V}_k (a_{k\sigma}^\dagger b_{k\sigma} + b_{k\sigma}^\dagger a_{k\sigma}) \quad (1)$$

where $a_{k\sigma}^\dagger$ and $b_{k\sigma}^\dagger$ are creation operators for the light a and the heavy b quasi-particles, respectively. The index $l = a, b$. The dispersion relations $\tilde{\epsilon}_k^l = (\hbar^2 k^2 / 2m_l) - \mu_l$ and the ratio between effective masses is taken as $\alpha = m_a / m_b < 1$. For simplicity, we renormalize all the energies in the problem by the chemical potential μ_a of band a (the *non-tilde* quantities). Furthermore, we take $\hbar^2 / (2m_a \mu_a) = 1$. In this case, the dispersion relations can be written as, $\epsilon_k^a = \tilde{\epsilon}_k^a / \mu_a = k^2 - 1$ and $\epsilon_k^b = \tilde{\epsilon}_k^b / \mu_a = \alpha k^2 - b$, with $b = \mu_b / \mu_a$.

The V -term is responsible for the transmutation among the quasi-particles. In metallic systems, as transition metals [15], inter-metallic compounds and heavy fermions [16], it is due to the mixing of the wave-functions of the quasi-particles in different orbitals through the crystalline potential. In the problem of color superconductivity, it is the weak interaction that allows the transformation between up- and down-quarks [2,4,12]. For a system of cold fermionic atoms in an optical lattice with two atomic states (a and b), the V -term is due to Raman transitions with an effective Rabi frequency which is directly proportional to V [10]. Then, the physical origin of the V -term is different for each case. The main point is that at least in inter-metallic systems, hybridization can be easily controlled by pressure or doping [17] allowing to explore their phase diagram using these quantities as external parameters. Notice that since hybridization transforms a quasi-particle into another, in its presence only the total number of particles is conserved.

There are two order parameters that characterize the different superconducting phases of the system described by the above Hamiltonian. They are, $\Delta_{ab} = -g \sum_{k\sigma} \langle b_{-k-\sigma} a_{k\sigma} \rangle$ and $\Delta = -U \sum_{k\sigma} \langle b_{-k-\sigma} b_{k\sigma} \rangle$. These are related to inter-band and intra-band superconductivity, respectively. The anomalous correlation functions $\langle b_{-k-\sigma} a_{k\sigma} \rangle$ and $\langle b_{-k-\sigma} b_{k\sigma} \rangle$ can be obtained from the corresponding Greens functions which also yield the spectrum of excitations in the superconducting phases. We use the equation of motion method to calculate the anomalous Greens functions [18]. Excitonic types of correlations that just renormalize the hybridization [19] are neglected. The relevant anomalous, frequency dependent Greens functions are, $\langle \langle a_{k\sigma}; b_{-k-\sigma} \rangle \rangle_\omega$ and $\langle \langle b_{k\sigma}; b_{-k-\sigma} \rangle \rangle_\omega$. Here we use the notation of Ref. [20]. When we write the equations of motion for these Greens functions, new Greens functions are generated [18]. Some of these are of higher order, as they contain a larger number of creation and annihilation operators than just the two of the initials Greens functions. For these, we apply a BCS type of decoupling [18] to reduce them to the order of the originals propagators. Writing the equations of motion for the new Greens functions, we obtain a closed system of equations that can be solved [5]. Finally, the anomalous frequency dependent propagators, from which the order parameters can be self-consistently obtained, are given by [5]

$$\langle \langle a_{k\sigma}; b_{-k-\sigma} \rangle \rangle_\omega = \frac{\Delta_{ab} N(\omega)}{\omega^4 + C_2 \omega^2 + C_1 \omega + C_0} \quad (2)$$

with

$$N(\omega) = \Delta_{ab}^2 - V^2 - (\omega - \epsilon_k^b)(\omega - \epsilon_k^a) + \frac{\Delta}{\Delta_{ab}} (\omega + \epsilon_k^a)$$

and

$$\langle \langle b_{k\sigma}; b_{-k-\sigma} \rangle \rangle_\omega = \frac{\Delta \left[(\omega^2 - \epsilon_k^{a2}) + 2 \frac{\Delta_{ab}}{\Delta} V \omega \right]}{\omega^4 + C_2 \omega^2 + C_1 \omega + C_0} \quad (3)$$

where

$$C_2 = -[\epsilon_k^{a2} + \epsilon_k^{b2} + \Delta^2 + 2(\Delta_{ab}^2 + V^2)]$$

$$C_1 = 4\Delta_{ab}\Delta$$

$$C_0 = [\epsilon_k^a \epsilon_k^b - (V^2 - \Delta_{ab}^2)]^2 + \Delta^2 \epsilon_k^{a2} \quad (4)$$

As mentioned before, the poles of these propagators give the energies (ω) of the excitations in the superconducting phase. Also from the discontinuity of the Greens functions on the real axis we can obtain the anomalous correlation functions characterizing the superconducting state. In general, the appearance of exotic superconducting phases is related to the existence of soft modes in the spectrum of excitations [7]. In the present case, for the energy of the excitations to vanish, it is required that $[\epsilon_k^a \epsilon_k^b - (V^2 - \Delta_{ab}^2)]^2 + \Delta^2 \epsilon_k^{a2} = 0$. This can occur by tuning the hybridization parameter, such that, $V = \Delta_{ab}$ in which case gapless excitations appear at $k = k_F^a$ where $\epsilon_k^a = 0$. Without this fine tuning there are no gapless modes. However, in case the intra-band interaction vanishes, there is a zero energy mode for the wave-vector k , such that, $\epsilon_k^a \epsilon_k^b - (V^2 - \Delta_{ab}^2) = 0$. We will see the effects of this behavior in the next section. If, for symmetry reasons, we neglect the term linear in ω ($C_1 = 0$), we obtain the energy of the excitations in the form

$$\omega_{1,2}(k) = \sqrt{A_k \pm \sqrt{B_k}} \quad (5)$$

with

$$A_k = \frac{\epsilon_k^{a2} + \epsilon_k^{b2}}{2} + \Delta_{ab}^2 + V^2 + \frac{\Delta^2}{2} \quad (6)$$

and

$$B_k = \left(\frac{\epsilon_k^{a2} - \epsilon_k^{b2}}{2} \right)^2 + V^2 (\epsilon_k^a + \epsilon_k^b)^2 + \Delta_{ab}^2 (\epsilon_k^a - \epsilon_k^b)^2 + 4V^2 \Delta_{ab}^2 + \frac{\Delta^4}{4} - \frac{\Delta^2}{2} (\epsilon_k^{a2} - \epsilon_k^{b2}) + \Delta^2 (V^2 + \Delta_{ab}^2) \quad (7)$$

The order parameters are determined self-consistently by a set of two coupled equations which for finite temperatures are given by [5]

$$\frac{1}{g\rho_a} = \sum_{j=1}^2 \int_{-\omega_D}^{\omega_D} \frac{(-1)^j d\epsilon}{2\sqrt{B(\epsilon)}} \left[\frac{\omega_j^2(\epsilon) - \lambda^2(\epsilon)}{2\omega_j(\epsilon)} \right] \tanh \frac{\beta\omega_j(\epsilon)}{2} \quad (8)$$

with

$$\lambda^2(\epsilon) = \left[\frac{\epsilon + (\alpha\epsilon - b)}{2} \right]^2 + (\Delta_{ab}^2 - V^2) + \frac{\Delta V}{4} \left[\Delta V + 4 \left(\frac{\epsilon + (\alpha\epsilon - b)}{2} \right) \right] - \left[\frac{\epsilon + (\alpha\epsilon - b)}{2} - \frac{\Delta V}{2} \right]^2 \quad (9)$$

where in B_k and $\omega_j(k)$, we substituted $\epsilon_k^a = \epsilon$ and $\epsilon_k^b = \alpha + (\alpha\epsilon - b)$.

$$\frac{1}{U\rho_b} = \sum_{j=1}^2 \int_{-\omega_D}^{\omega_D} \frac{(-1)^j d\epsilon}{2\sqrt{B(\epsilon)}} \left[\frac{\alpha^2 \omega_j^2(\epsilon) - (\epsilon + b - \alpha)^2}{2\alpha^2 \omega_j(\epsilon)} \right] \tanh \frac{\beta\omega_j(\epsilon)}{2} \quad (10)$$

In this equation, we substituted $\epsilon_k^b = \epsilon$ and $\epsilon_k^a = (\epsilon + b - \alpha)/\alpha$ in B_k and $\omega_j(k)$. The quantities ρ_a and ρ_b are the density of states at the

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