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Phase detection in an ultracold polarized Fermi gas via electromagnetically induced transparency

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ABSTRACT

An optical detection scheme based on electromagnetically induced transparency is proposed to investigate the nature of Fermi pairing in a trapped strongly interacting two-component atomic Fermi gas with population imbalance. We show that valuable information can be acquired from the *in situ* probe absorption spectrum to identify different phases existing in the system.

1. Introduction

The experimental realization of ultracold degenerate atomic Fermi gas, together with the unprecedented control over the interatomic interaction by the means of Feshbach resonance, have provided a powerful playground for the study of pairing and superfluidity. For a two-component Fermi gas with equal population, the pairing mechanism have been well understood since the pioneering work of Bardeen, Cooper and Schrieffer (BCS) [1]. On the other hand, many recent investigations focused on polarized Fermi gas with imbalanced spin populations, which is expected to shed light on many different fields of physics, for example, magnetized superconductors and QCD [2].

For a polarized Fermi gas, population imbalance together with the resonant strong interaction have introduced a few non-BCS exotic superfluid phases into the system such as the inhomogeneous Fulde–Ferrell–Larkin–Ovchinnikov (FFLO) state [3] with non-zero-momentum pairing and the breached-pair (BP) state [4] in which the particles are paired except for a finite region of momentum shell. Polarized Fermi gas is predicted to exhibit a far richer phase diagram than the equal spin case [5–12] and have been realized in experiment [13–17]. In the presence of the trap, due to the effective chemical potentials become space-dependent, spatial phase separation will take place. An important question is how to unambiguously differentiate these different phases.

In previous experiment, people use phase-contrast imaging technique to map the phase diagram of a polarized Fermi gas [18]. In this detection scheme, by appropriately choosing the probe frequency, the probe light dispersively interacts with the two atomic components simultaneously and results in a phase shift proportional to the density difference between the two atomic components. They use the discontinuity in the spin polarization profile as the signature for phase transition, which can only indicate the first-order superfluid-to-normal phase transition. The higher-order phase transitions, such as that between the BCS superfluid and BP phase, is of at least third-order nature and cannot be detected.

In order to clearly identify different superfluid phases, one will have to go beyond first-order density–density correlation to gain pairing information. This is the idea lies at the heart of early proposals on the optical detection of BCS phase transition [19–21] as well as other methods used to study strongly interacting atomic Fermi gases such as the spatial noise correlations [22], the momentum-resolved stimulated Raman technique [23–25] and radio-frequency (RF) spectroscopy [26–30].

In this work we discuss the application of an alternative detection scheme based on electromagnetically induced transparency (EIT), which was proposed recently for the detection of BCS pairing [31], to detect various phases in a trapped polarized Fermi gas. The probe absorption spectrum for different quantum phases in the system is studied. We demonstrate that, the probe spectra will display unique properties for different phases, which can serve as the proof for unambiguous phase discrimination.

Compared to the previous detection methods (absorption imaging or phase-contrast imaging), only one additional control light field is required for the EIT-based detection scheme. In the meanwhile, it is principally non-destructive and does not induce much heating of the atomic sample, which makes the continuous

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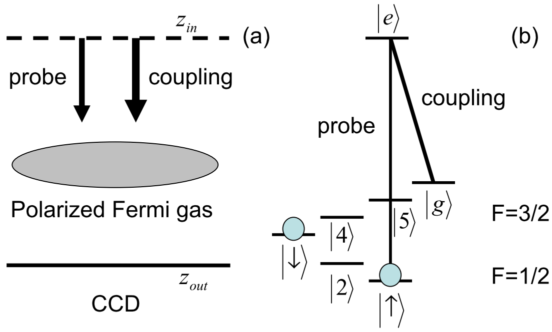

Fig. 1. (Color online.) Phase detection scheme for trapped two-component Fermi gas with population imbalance. (a) We consider the two light fields to be copropagating and input in the $+z$ direction at z_{in}. The output intensity at z_{out} can be detected with a CCD camera. The coupling beam is cw. (b) The level configuration for the ^{6}Li atoms near the Feshbach resonance point $B_0 = 690$ G. The probe light couples the atoms of the lowest level ($|\uparrow\rangle$) to the electronically excited state $|e\rangle$, while the coupling light couples the $|e\rangle$ with the ground-state sublevel $|g\rangle$.

measurements possible without the need to prepare the atomic sample repeatedly.

The Letter is organized as follows: In Section 2, we illustrate the basic idea of the detection scheme and the Hamiltonian describing the system will be given. The response of the gas to the probe light is derived in Section 3, the probe spectra for a few typical different phases will be speculated in detail. In Section 4, we provide the numerical examples that showcase the reliability of our detection scheme both at zero and finite temperature. Finally, a summary is given in Section 5.

2. Model and Hamiltonian

The detection scheme is illustrated in Fig. 1. To be more specific, we consider the case of ^{6}Li, atoms are initially prepared in the ground-state manifold $|\uparrow\rangle$ and $|\downarrow\rangle$, which correspond to the ground-state manifold $|F = 1/2,\ m_F = 1/2\rangle$ and $|F = 3/2,\ m_F = -3/2\rangle$ respectively, just the same case as [26,32]. The two atomic components are resonantly coupled to each other via Feshbach resonance. In the case of broad Feshbach resonance, the molecular population can be safely neglected and the atomic system is described by the single-channel model. We assume atoms with spin $\uparrow$ to be the majority component. A weak probe field with linear polarization parallel to the magnetic field drives the transition between $|\uparrow\rangle$ and the electronically excited state $|e\rangle$, whereas a strong control field with σ_- polarization couples $|e\rangle$ with another initially empty ground state manifold $|g\rangle$ ($|F = 3/2,\ m_F = 3/2\rangle$). The light field is characterized by the Rabi frequency $\Omega_{p(c)} \exp[i(\mathbf{k}_{p(c)} \cdot \mathbf{r} - \omega_{p(c)}t)]$ with wave-vector $\mathbf{k}_{p(c)}$ and frequency $\omega_{p(c)}$. We assume that the two lasers are in the co-propagation configuration so that the wave-vectors of the two laser fields remain essentially the same and are denoted as $\mathbf{k}_L$. In the essence of local density approximation (LDA), we assume the chemical potential to be coordinate-dependent and treat the system as homogeneous at each position. Let $\hat{a}_{\mathbf{k},i}$ be the annihilation operator for a particle of state i with momentum $\mathbf{k}$, in the rotating-wave approximation the Hamiltonian reads

$$\hat{H} = \hat{H}_\sigma + \sum_{\mathbf{k}} \Big[(\epsilon'_{\mathbf{k},\uparrow} - \delta_{1p}) \hat{a}^\dagger_{\mathbf{k},e} \hat{a}_{\mathbf{k},e} + (\epsilon'_{\mathbf{k},\uparrow} - \delta_2) \hat{a}^\dagger_{\mathbf{k},g} \hat{a}_{\mathbf{k},g} - \frac{\Omega_c}{2} (\hat{a}^\dagger_{\mathbf{k}+\mathbf{k}_L,e} \hat{a}_{\mathbf{k},g} + \text{h.c.}) - \frac{\Omega_p}{2} (\hat{a}^\dagger_{\mathbf{k}+\mathbf{k}_L,e} \hat{a}_{\mathbf{k},\uparrow} + \text{h.c.}) \Big], \tag{1}$$

with

$$\hat{H}_\sigma = \sum_{\mathbf{k},\sigma=\uparrow,\downarrow} \epsilon'_{\mathbf{k},\sigma} \hat{a}^\dagger_{\mathbf{k},\sigma} \hat{a}_{\mathbf{k},\sigma} - \sum_{\mathbf{k}} (\Delta \hat{a}^\dagger_{\mathbf{k},\uparrow} \hat{a}^\dagger_{-\mathbf{k},\downarrow} + \text{h.c.}), \tag{2}$$

where $\epsilon'_{\mathbf{k},\sigma} = \epsilon_{\mathbf{k}} - \mu_\sigma$, with $\epsilon_{\mathbf{k}} = \mathbf{k}^2/2m$ (m is the atomic mass and $\hbar = 1$) and μ_σ being the chemical potential for the spin σ component. Under LDA we have $\mu_\uparrow = \mu_{\mathbf{r}} + h$, $\mu_\downarrow = \mu_{\mathbf{r}} - h$, $\mu_{\mathbf{r}} = \mu - V(\mathbf{r})$, where $V(\mathbf{r})$ is the trap potential, and μ, h are the chemical potential at the trap center and the chemical potential difference, respectively. $\delta_{1p} = \omega_p - \omega_{e\uparrow}$ and $\delta_{1c} = \omega_c - \omega_{eg}$ are single-photon detunings, $\delta_2 = \delta_{1p} - \delta_{1c}$ is the two-photon detuning, with ω_{ij} being the atomic transition frequency between level i and j. The pairing gap Δ at zero temperature is given by $\Delta_s = -U \sum_{\mathbf{k}} \langle \hat{a}_{-\mathbf{k},\downarrow} \hat{a}_{\mathbf{k},\uparrow} \rangle / \mathcal{V}$, where $U = 4\pi a/m$ is the bare background scattering rate with the s-wave scattering length a and $\mathcal{V}$ is the quantization volume. At finite temperature Δ should be understood as the total gap with $\Delta^2 = \Delta_s^2 + \Delta_{pg}^2$, where Δ_{pg} is the pseudogap coming from the contribution of noncondensed Cooper pairs [40]. Without loss of the essential physics, we have ignored the Hartree potential term in the atomic Hamiltonian (2).

Similar to [31], for the sake of simplicity, we have neglected the collisions involving states $|g\rangle$ and $|e\rangle$ in Hamiltonian (1). This is made possible by appropriately choosing the pairing states. For the specific case considered in Fig. 1, in which the pairing states are just the same as those adopted in [26,32], at the resonance position of $B_0 = 690$ G, one can expect that the final state scattering length is small and it is a safe treatment to ignore the collision interactions of $|g\rangle$ ($|e\rangle$) with either of the pairing state.

In typical EIT, the probe field is much weaker than the control field ($\Omega_p \ll \Omega_c$), the depletion of ground-state atoms is small. So we can assume that during the detection process, the atoms are kept in the state determined by Hamiltonian (2) alone. The ground state of (2) have been investigated in detail [6–10], where a few different phases can exist and one may cross several different phases from the trap center to the edge.

In order to gain physical insights, we introduce two internal dressed states $|+\rangle$ and $|-\rangle$ and define the corresponding field operators by the canonical transformation $\hat{\alpha}_{\mathbf{k},+} = \hat{\alpha}_{\mathbf{k},\uparrow}$ and $\hat{\alpha}_{\mathbf{k},-} = \hat{\alpha}^\dagger_{-\mathbf{k},\downarrow}$. Here $\hat{\alpha}_{\mathbf{k},\uparrow}$ and $\hat{\alpha}^\dagger_{-\mathbf{k},\downarrow}$ are the quasiparticle operators defined in the usual Bogoliubov transformation

$$\begin{pmatrix} \hat{a}_{\mathbf{k},\uparrow} \\ \hat{a}^\dagger_{-\mathbf{k},\downarrow} \end{pmatrix} = \begin{pmatrix} u_{\mathbf{k}} & v_{\mathbf{k}} \\ -v_{\mathbf{k}} & u_{\mathbf{k}} \end{pmatrix} \begin{pmatrix} \hat{\alpha}_{\mathbf{k},\uparrow} \\ \hat{\alpha}^\dagger_{-\mathbf{k},\downarrow} \end{pmatrix}, \tag{3}$$

with

$$u_{\mathbf{k}}^2 = \left(1 + \frac{\epsilon_{\mathbf{k}} - \mu_{\mathbf{r}}}{E_{\mathbf{k}}}\right)\Big/2, \qquad v_{\mathbf{k}}^2 = \left(1 - \frac{\epsilon_{\mathbf{k}} - \mu_{\mathbf{r}}}{E_{\mathbf{k}}}\right)\Big/2,$$
$$E_{\mathbf{k}} = \sqrt{(\epsilon_{\mathbf{k}} - \mu_{\mathbf{r}})^2 + \Delta^2}. \tag{4}$$

The grand Hamiltonian (1) now becomes

$$\hat{H} = \sum_{\mathbf{k}} \Big[E_{\mathbf{k}+} \hat{\alpha}^\dagger_{\mathbf{k},+} \hat{\alpha}_{\mathbf{k},+} + E_{\mathbf{k}-} \hat{\alpha}^\dagger_{\mathbf{k},-} \hat{\alpha}_{\mathbf{k},-} + (\epsilon'_{\mathbf{k},\uparrow} - \delta_{1p}) \hat{a}^\dagger_{\mathbf{k},e} \hat{a}_{\mathbf{k},e} + (\epsilon'_{\mathbf{k},\uparrow} - \delta_2) \hat{a}^\dagger_{\mathbf{k},g} \hat{a}_{\mathbf{k},g} - \frac{\Omega_c}{2} (\hat{a}^\dagger_{\mathbf{k}+\mathbf{k}_L,e} \hat{a}_{\mathbf{k},g} + \text{h.c.}) - \frac{u_{\mathbf{k}} \Omega_p}{2} (\hat{a}^\dagger_{\mathbf{k}+\mathbf{k}_L,e} \hat{\alpha}_{\mathbf{k},+} + \text{h.c.}) - \frac{v_{\mathbf{k}} \Omega_p}{2} (\hat{a}^\dagger_{\mathbf{k}+\mathbf{k}_L,e} \hat{\alpha}_{\mathbf{k},-} + \text{h.c.}) \Big], \tag{5}$$

where $E_{\mathbf{k}\pm} = \pm(E_{\mathbf{k}} \mp h)$. The physical picture emerges from this Hamiltonian is that, the state $|\pm\rangle$ has the energy dispersion $E_{\mathbf{k}\pm}$ and coupled to the excited state with an effective Rabi frequency $u_{\mathbf{k}} \Omega_p$ ($v_{\mathbf{k}} \Omega_p$). In the dressed state picture, our model becomes a double Λ system. We'll take advantage of the unique spectroscopic features of this double EIT for phase detection, as will be discussed in detail in the following section.