Non-Hermitian Ferromagnetism in an Ultracold Fermi Gas

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(Dated: April 14, 2020)

We develop a non-Hermitian effective theory for a repulsively interacting Fermi gas in the excited branch. The on-shell T-matrix is employed as a complex-valued interaction term, which describes a repulsive interaction between atoms in the excited branch and a two-body inelastic decay to the attractive branch. To see the feature of this model, we have addressed, in the weak coupling regime, the excitation properties of a repulsive Fermi polaron as well as the time-dependent number density. The analytic expressions obtained for these quantities qualitatively show a good agreement with recent experiments. By calculating the dynamical transverse spin susceptibility in the random phase approximation, we show that a ferromagnetic system with nonzero polarization undergoes a dynamical instability and tends towards a heterogeneous phase.

PACS numbers: 03.75.Ss, 03.75.-b, 03.70.+k

Introduction—An ultracold atomic gas has attracted much attention in modern physics because various quantum many-body phenomena manifest themselves due to controllability of such physical parameters as scattering length and density [1, 2]. For example, the realization of crossover from the Bardeen-Cooper-Schrieffer (BCS) Fermi superfluid to the molecular Bose-Einstein condensation (BEC) or vice versa [3, 4] by tuning the attractive interaction via a Feshbach resonance [5] has opened a new frontier for studies of strongly correlated quantum systems [6–8]. Nowadays, the ultracold atomic gas acts as a quantum simulator in various research fields ranging from condensed matter to nuclear physics [7–12].

In cold atomic physics, what kind of many-body state of repulsively interacting Fermi gases occurs is a long-standing problem [13]. A ferromagnetic phase transition is expected to occur due to the intrinsically short-range repulsion, given the analogy with the so-called Stoner system [9–12]. For example, the realization of crossover from the Bardeen-Cooper-Schrieffer (BCS) Fermi superfluid to the molecular Bose-Einstein condensation (BEC) or vice versa [3, 4] by tuning the attractive interaction via a Feshbach resonance [5] has opened a new frontier for studies of strongly correlated quantum systems [6–8]. Nowadays, the ultracold atomic gas acts as a quantum simulator in various research fields ranging from condensed matter to nuclear physics [7–12].

We develop a non-Hermitian effective theory for a repulsively interacting Fermi gas in the excited branch. The on-shell T-matrix is employed as a complex-valued interaction term, which describes a repulsive interaction between atoms in the excited branch and a two-body inelastic decay to the attractive branch. To see the feature of this model, we have addressed, in the weak coupling regime, the excitation properties of a repulsive Fermi polaron as well as the time-dependent number density. The analytic expressions obtained for these quantities qualitatively show a good agreement with recent experiments. By calculating the dynamical transverse spin susceptibility in the random phase approximation, we show that a ferromagnetic system with nonzero polarization undergoes a dynamical instability and tends towards a heterogeneous phase.

In this work, we develop a non-Hermitian effective theory for repulsively interacting two-spin-component Fermi gases in the excited branch. We model the non-Hermitian interaction for the inelastic two-body decay from the on-shell two-body T-matrix at finite momenta. To see how to connect between the complex-valued interaction and the existing experimental data, we derive excitation properties of a repulsive Fermi polaron in the polarized limit and also time-dependent number density in the unpolarized case. We show that the dynamical transverse spin susceptibility calculated within the random phase approximation (RPA) exhibits a dynamical instability in the long wavelength limit due to the non-Hermitian interaction in a metastable ferromagnetic state.
Non-Hermitian effective Hamiltonian—For a two-spin-component Fermi gas with short-range interactions, the Pauli principle allows us to write the effective Hamiltonian as

$$
H_{\text{eff}} = \sum_{p, \sigma} \varepsilon_p c_{p, \sigma}^\dagger c_{p, \sigma} + \sum_{k, k', p}\, U_R(k, k', P) \times (c_{P+2/k', 2/\sigma} c_{P-k/2, 2/\sigma} c_{P-k/2, 2/\sigma})^\dagger,
$$

(1)

where $\varepsilon_p = \varepsilon_p - \mu_\sigma \equiv p^2/2m - \mu_\sigma$ is the kinetic energy of a Fermi atom with momentum $p$, pseudospin $\sigma = \uparrow, \downarrow$, and mass $m$ measured from the chemical potential $\mu_\sigma$, and $c_{p, \sigma}$ is the annihilation operator. $U_R(k, k', P)$ is the $\uparrow \downarrow$ repulsive interaction that will be specified below as function of the relative momenta $k$ and $k'$ of incoming and outgoing two particles, as well as the center of mass momentum $P$.

We proceed to derive the repulsive interaction $U_R$ from the on-shell two-body $T$-matrix with an attractive contact interaction $U(<0)$ as

$$
T(k, k; 2\varepsilon_k + i\delta) = \left[ \frac{1}{U} - \sum_p \frac{1}{2\varepsilon_k + i\delta - 2\varepsilon_p} \right]^{-1} = \frac{4\pi a}{m} \frac{1 - ika}{1 + k^2a^2},
$$

(2)

where we set the effective range to zero by taking an infinitely large cutoff $\Lambda$, $\delta$ is a positive infinitesimal, and $a = (\frac{4\pi}{m\Gamma} + \frac{2}{Z})^{-1}$ is the $s$-wave scattering length. While the constant repulsive interaction $U_R = T(0, 0; 0) = \frac{4\pi a}{m}$ in the excited branch is usually employed to study possible ferromagnetism, we here keep the incoming momentum dependence of the $T$-matrix and obtain $U_R(k) = \frac{4\pi a}{m} \frac{1 - ika}{1 + k^2a^2}$, which is complex-valued. In $U_R(k)$ we have set $k' = k$ and $P = 0$ for simplicity. This complex-valued effective interaction, which behaves as $1/(1 + ika)$, reflects the fact that the excited branch is unstable against inelastic decay to the molecular state of energy $-1/m\alpha^2$ in the attractive branch. While this imaginary part is negligible when $a$ is small, it would play a significant role near the ferromagnetic transition expected to occur in the strongly interacting regime.

Repulsive Fermi polaron—To see the feature of our non-Hermitian model, we first consider the zero-temperature ($T = 0$), highly polarized case in which an impurity atom ($\sigma = \downarrow$) immersed in a Fermi sea of noninteracting majority atoms ($\sigma = \uparrow$) forms a repulsive Fermi polaron. The retarded Green’s function of such an impurity is given by $G_R^R(p, \omega) = [\omega - \varepsilon_p - \Sigma_R^R(p, \omega)]^{-1}$. Here,

$$
\Sigma_R^R(p) = \sum_{p'} U_R\left(\left|\frac{p - p'}{2}\right|\right) f(\xi_{p', \uparrow}),
$$

(3)

with the Fermi-Dirac distribution function $f(\omega) = (e^{\beta\omega} + 1)^{-1}$, is the impurity self-energy that retains only the lowest-order diagram (Hartree correction). This self-energy can be regarded as an approximate form of the $T$-matrix approximation (TMA) [47, 53]. While TMA allows for the Pauli blocking of majority atoms in intermediate states that are included in the ladder diagrams, the present approach neglects such diagrams and hence reduces to a mean-field approximation to the medium effect. We can obtain the simple analytical expressions for the repulsive polaron energy $E_r = \text{Re}\Sigma_R^R(0)$, the decay rate $\Gamma = -2\text{Im}\Sigma_R^R(0)$, and the effective mass $m/m^* = 1 + m \frac{\partial^2 \text{Re}\Sigma_R^R(p)}{\partial p^2} \bigg|_{p=0}$ as

$$
E_r = \frac{1}{\pi} \frac{16}{k_F \gamma} \left[ 1 - \frac{2\tan^{-1}\left(\frac{k_F \gamma}{2}\right)}{k_F \gamma^2} \right],
$$

(4)

$$
\Gamma = \frac{32}{\pi} \left(\frac{k_F \gamma}{2}\right)^2 \ln \left[ 1 + \left(\frac{k_F \gamma}{2}\right)^2 \right],
$$

(5)

and

$$
m/m^* = 1 - \frac{16}{3\pi} \left(\frac{k_F \gamma}{2}\right)^3 \left(4 + k_F \gamma^2/2\right),
$$

(6)

where $\varepsilon_F, \gamma$ and $k_F \gamma$ are the majority Fermi energy and momentum, respectively.

In Fig. [H] we plot the obtained results for (a) $E_r$, (b) $\Gamma$, and (c) $m^*$, together with the experimental results [21] and earlier calculations [43, 54, 55]. Although we have used a rather simple method, our result for $E_r$ shows a good agreement with the experimental and TMA results, particularly in the weakly repulsive regime $k_F \gamma \lesssim 1$. This is consistent with the fact that our approach, a simplified version of TMA, is what TMA reduces to in the weak coupling limit. The overestimation of $\Gamma$ in our approach as compared with the empirical values even in the weak coupling regime, on the other hand, can be understood as lack of the Pauli-blocking effect and of the three-body recombination process [21, 56] in the two-body $T$-matrix. Incidentally, near the unitarity limit, our results underestimate the repulsive polaron energy. From the comparison between our results and TMA, this may possibly be because our approach ignores the role played by the Pauli blocking in the attractive branch in indirect suppression of the two-body loss. Fluctuation corrections beyond the present mean-field approach may also be responsible for such underestimate. Although the Hartree correction contains no energy dependence and hence keeps the quasi-particle residue $Z$ at unity, in some cases, even the lowest order correction due to the interaction requires the $Z$ contribution. For example, the effective mass undergoes a leading order modification by $Z$ as $m/m^* = Z\left[1 + m \frac{\partial^2 \text{Re}\Sigma_R^R(p)}{\partial p^2} \bigg|_{p=0}\right]$, where
FIG. 1: (a) Repulsive polaron energy $E_r$, (b) decay rate $\Gamma$, and (c) effective mass $m^*$, calculated from the non-Hermitian model as function of $1/k_{F,\uparrow}a$. For comparison, we also show the experimental results and theoretical curves of TMA, variational method (VAR) and functional renormalization group (FRG). In panel (b), the theoretical results for the decay rates associated with polaron-to-free particle decay ($\Gamma_{\text{PF}}$), polaron-to-polaron decay ($\Gamma_{\text{PP}}$), and three-body recombination ($\Gamma_3$) are also plotted. The dash-dotted curve in panel (c) shows $m^*$ modified by the quasi-particle residue $Z$ shown in the inset.

$$Z = \left[ 1 - \text{Re} \frac{\partial \gamma_\sigma}{\partial \omega} \bigg|_{\omega=0,\omega} \right]^{-1}$$

The resultant $m^*$ agrees well with the experimental and TMA results.

**Lindblad equation and number densities**— We next show how the number densities $N_\sigma$ behaves via the two-body loss suffered by the repulsive branch due to the non-Hermitian interaction. We assume that the system considered here is described by an open quantum system in which the main system in the repulsive branch interacts with the bath in the attractive branch via a density-dependent interaction $\hat{U}_R = U_R(k = k_\alpha)$, where $k_\alpha = N_\uparrow N_\downarrow \frac{1}{N_\uparrow N_\downarrow} \sum |k| \leq k_{F,\uparrow} \sum |k'| \leq k_{F,\downarrow} \frac{|k-k'|}{2}$ is the averaged relative momentum with the Fermi momentum of $\sigma$ atoms.

In describing the quantum dynamics of the system in the repulsive branch, we utilize the Lindblad equation for the reduced density matrix $\rho$ in the repulsive branch, which is given by

$$i \frac{d \rho}{dt} = [H, \rho] - i \{K, \rho\} + i \int d^3 r L(r) \rho L^\dagger(r), \quad (8)$$

where $H$ and $K$ are the Hermitian and non-Hermitian parts of the effective Hamiltonian $H$, i.e., $H_{\text{eff}} = H - iK$, and $L(r)$ is the local Lindblad operator as given by

$L(r) = \sqrt{-2\text{Im} \hat{U}_R \psi_\uparrow(r) \psi_\uparrow^\dagger(r)}$ in such a way as to satisfy $K = \frac{1}{2} \int d^3 r L^\dagger(r)L(r)$. Using Eq. (5), one can obtain the time derivative of $N_\sigma = \text{Tr}(\rho N_\sigma)$ with the density operator $N_\sigma = \sum_p \bar{p}_p c_{p,\sigma}$ as

$$\frac{d N_\sigma}{dt} = 2\text{Im} \hat{U}_R \int d^3 r (\psi_\uparrow^\dagger(r) \psi_\uparrow^\dagger(r) \psi_\uparrow(r) \psi_\uparrow^\dagger(r)) \approx 2\text{Im} \hat{U}_R N_\uparrow N_\downarrow, \quad (9)$$

where we have used a commutation relation $[\hat{N}_\sigma, L(r)] = -L(r)$ as well as the weak coupling approximation in the second line. From Eq. (9), one can find $\frac{dN_\sigma}{dt} \approx \frac{\partial}{\partial t} (N_\uparrow - N_\downarrow) = 0$, which indicates that the two-body loss itself would not directly suppress the magnetization $M$, if any at all, but would drive such a ferromagnetic system to phase separation between the ferromagnetic and molecular states due to decrease of the number density in the repulsive branch.

Figure 2 shows the calculated $N_\sigma(t)$ for an unpolarized repulsive Fermi gas that has an initial number density $N_\sigma(0) = \frac{k_{F,\sigma}^3}{3\pi^2}$, which is plotted as function of $t \varepsilon_F \sigma$ with the common Fermi energy $\varepsilon_F, \varepsilon_F \uparrow = \varepsilon_F \downarrow$. As is consistent with the empirical behavior, the resulting decay rate tends to increase with $\alpha$. We note, however, that our results overestimate the particle loss during the time evolution. While we here use the density-dependent interaction $\hat{U}_R$ relevant for the $T = 0$ Fermi degeneracy, no temperature dependence has been considered. In the realistic case of nonzero temperature, the Fermi distribution has a thermal diffuseness, which acts to weaken the interaction by allowing scattering between two atoms of lower momenta to occur. In addition, the inverse process, namely, the pumping to the excited branch due to dissociation of a molecule in the attractive branch,
where \( \chi \) in Eq. (11), \( G \) use the Keldysh component in thermal equilibrium as another particle is added to the system, i.e., \( \mu_\sigma = \frac{\partial E}{\partial N_\sigma} \), reads \( \mu_\sigma = \epsilon_{F,\sigma} + \bar{U}_R N_{-\sigma} \in \mathbb{C} \). This is because within the mean-field approximation, \( E = \frac{3}{2} N_\uparrow \epsilon_{F,\uparrow} + \frac{2}{5} N_\downarrow \epsilon_{F,\downarrow} + \bar{U}_R N_{-\sigma} \in \mathbb{C} \). By substituting this \( \mu_\sigma \) to \( G^R_\sigma(p, \omega) = (\omega - \epsilon_p + \mu_\sigma - \bar{U}_R N_{-\sigma} + i\delta)^{-1} \), one can find that \( \text{Im} \bar{U}_R N_{-\sigma} \) is compensated by the imaginary part of \( \mu_\sigma \). Consequently, \( \chi^R_\sigma(q, \omega) \) reduces to the well-known Lindhard function,

\[
\chi^R_\sigma(q, \omega) = - \sum_k \frac{f(\xi_{k,\sigma}^\ast + q, \omega) - f(\xi_{k,\sigma})}{\omega + \xi_{k,\sigma}^\ast - \xi_{k+q,\sigma}^\ast + i\delta},
\]

where \( \xi_{k,\sigma}^\ast = \epsilon_k - \mu_\sigma^\ast \) with the renormalized chemical potential \( \mu_\sigma^\ast = \mu_\sigma - \bar{U}_R N_{-\sigma} \) (corresponding to the Fermi energy \( \epsilon_{F,\sigma} \)) is the excitation energy. It is to be noted that \( \mu_\sigma^\ast \) is now real-valued, a feature that is also consistent with the number equation (12).

Finally, we address what kind of ferromagnetic state is realized in the presence of the inelastic two-body decay. To this end, for simplicity, we focus on the zero-momentum pole \( \Omega \) of the RPA susceptibility \( \chi^R(q, \omega) \), Eq. (10), which fulfills

\[
1 + \bar{U}_R \chi^R_\sigma(q \to 0, \Omega) = 0.
\]

In the unpolarized case with \( \Omega = 0 \), the critical density that satisfies Eq. (14) corresponds to the condition for the Stoner instability, i.e., spontaneous polarization occurs statically and uniformly. In the presence of such polarization along which we take the direction of \( \uparrow \), we obtain a solution to Eq. (14) as \( \mu^\ast \)

\[
\text{Re} \Omega = \mu^\ast_\uparrow - \mu^\ast_\downarrow - \text{Re} \bar{U}_R(N_\uparrow - N_\downarrow),
\]

\[
\text{Im} \Omega = -\text{Im} \bar{U}_R(N_\uparrow - N_\downarrow) > 0.
\]

The pole in the upper complex plane of frequency indicates that the system undergoes a dynamical instability once the system has nonzero polarization. This is because the inverse Fourier transformation of \( \chi^R_\sigma(0, \omega) \) with respect to frequency leads to the time dependence like \( \exp(-i\Omega t) = \exp(-i\text{Re} \Omega t + i\text{Im} \Omega t) \). This exponential growth of spin fluctuations, together with the above argument based on Eq. (4), suggests that the ferromagnetic state does not undergo a phase transition back to the homogeneous paramagnetic (unpolarized) phase, but to a qualitatively different state that is reminiscent of a

\[
N_\sigma = -\frac{1}{\pi} \sum_p \int d\omega f(\omega) \text{Im} G^R_\sigma(p, \omega).
\]
heterogeneous phase observed in recent experiments [23]. We remark in passing that within the present model, a homogeneous transition from the paramagnetic to ferromagnetic phase with increasing density is unlikely to occur since the inelastic decay to the molecular state is designed to keep decreasing the density without any feedback. Recall that this caveat holds also for the study of the repulsive polaron energy as well as the decay of the atom number density. Further investigations beyond the present framework would be desired.

**Summary**— In this work, we have developed a non-Hermitian effective theory to describe the polaronic and magnetic properties of a repulsively interacting Fermi gas in the excited branch. This theory incorporates the complex-valued interaction obtained from the on-shell two-body \(T\)-matrix in such a way as to characterize the two-body inelastic decay to the molecular state in the attractive branch. Within the weak coupling approximation, we have derived simple analytical formulas for the repulsive polaron properties and the differential equation for the time-dependent atomic number density, which explain the experimental results fairly well given the overestimated two-body inelastic loss. By building the present non-Hermitian framework into the analysis of the dynamical transverse spin susceptibility within the RPA, we show that the complex-valued two-body interaction drives a uniform ferromagnetic system unstable to heterogeneity as observed in recent experiments. To investigate the spatial scale of the heterogeneous phase, the momentum dependence of the time-dependent susceptibility would have to be clarified.

The authors thank F. Scavia for providing us with their data in Ref. [21] and E. Nakano, J. Takahashi, K. Nishimura, T. Hata, T. M. Doi, and S. Tsutsui for useful discussions. This work is supported by Grants-in-Aid for JSPS fellows (No. 17J03975) and for Scientific Research from JSPS (Nos. 18H01211 and 18H05406).

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[62] The solution of Eq. (14) at nonzero ReΩ would mean the presence of a gapped collective mode (magnon) in the system with a broken spin-inversion symmetry. Then, the solution at vanishing ReΩ would give a condition at which the magnon becomes gapless and hence domains of itinerant ferromagnetism occur. Once an impurity atom (↓) is put into a resultant magnetic domain (↑), the repulsive polaron formed has an energy $E_r = \varepsilon_{F,↑}$, which can be obtained by substituting into Eq. (15) the relevant conditions ReΩ = 0, $N_\downarrow = 0$, and $\mu_\downarrow^* = 0$ as well as using Eq. (3). This just corresponds to the critical stability condition of the configuration that consists of magnetic domains as given in Ref. [49].