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Itinerant ferromagnetism in one-dimensional two-component Fermi gases

Yuzhu Jiang,[1] D. V. Kurlov,[2] Xi-Wen Guan,[1,3] F. Schreck,[2] and G. V. Shlyapnikov[1,2,4,5]
[1]*State Key Laboratory of Magnetic Resonance and Atomic and Molecular Physics, Wuhan Institute of Physics and Mathematics, Chinese Academy of Sciences, Wuhan 430071, China*
[2]*Van der Waals-Zeeman Institute, Institute of Physics, University of Amsterdam, Science Park 904, 1098 XH Amsterdam, The Netherlands*
[3]*Department of Theoretical Physics, Research School of Physics and Engineering, Australian National University, Canberra ACT 0200, Australia*
[4]*LPTMS, CNRS, Univ. Paris-Sud, Université Paris-Saclay, 91405 Orsay, France*
[5]*Russian Quantum Center, Novaya Street 100, Skolkovo, Moscow Region 143025, Russia*

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We study a one-dimensional (1D) two-component atomic Fermi gas with an infinite intercomponent contact repulsion. It is found that adding an attractive nearly resonant odd-wave interaction breaking the rotational symmetry one can make the ground state ferromagnetic. A promising system for the observation of this itinerant ferromagnetic state is a 1D gas of ^{40}K atoms, where three dimensional s-wave and p-wave Feshbach resonances are very close to each other and the 1D confinement significantly reduces the inelastic decay.

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Itinerant ferromagnetism of degenerate spin-$1/2$ fermions is an intriguing problem promoting our understanding of strongly correlated systems [1]. Ultracold atomic gases attract a great interest for studying itinerant ferromagnetism because they are highly controllable and are tunable with respect to interactions. The origin of such ferromagnetic states is deeply rooted in quantum mechanics. In contrast to ultracold bosons, degenerate fermions try to avoid the ferromagnetic state because it requires them to have a significantly higher kinetic energy than in nonferromagnetic states. Ultracold gases of atomic fermions that are in two internal states can be mapped onto spin-$1/2$ fermions treating the internal energy levels as pseudospin states. The ferromagnetic phase is the one where all atoms are in the same superposition of the two internal states and one has a system of identical fermions. The kinetic energy is then higher than, for example, in the paramagnetic phase, which represents a statistical mixture of the two spin components.

Itinerant ferromagnetism for fermions has been studied since the 1930s, when the Stoner criterion for ferromagnetism in a free electron gas was introduced [2,3]. According to this criterion, in three dimensions (3D) the ground state can be ferromagnetic if there is a strong intercomponent repulsion in the paramagnetic state, which compensates for the large difference in the kinetic energies of these states. Developments in condensed matter physics [4,5] found that itinerant ferromagnetism is responsible for some of the properties of transition metals, such as cobalt, iron, and nickel. Theoretical studies were concentrated on the Hubbard model with strong contact interactions (spin independent with next to nearest neighbor hopping [6,7] or generic multiorbital on-site interactions [8]) and on the closely related quantum rotor model [9]. Except for the original Stoner model, such interactions are hard to be realized with ultracold atomic gases. The Stoner mechanism in Fermi gases was discussed in a number of papers [10], and Monte Carlo calculations [11,12] found an instability on approach to the strongly interacting regime. The efforts to stabilize the ferromagnetic state experimentally did not succeed [13,14]. In three dimensions a large intercomponent repulsion corresponds to a very large and positive s-wave scattering length and there is a weakly bound dimer of two fermions belonging to different internal states. In this situation atom loss by dimer formation is very fast at typical densities [13–15]. At the same time, the studies including the momentum dependence of the s-wave interaction (finite effective range) found that this process can be reduced near a narrow Feshbach resonance [16].

In one and two dimensions the difference in the kinetic energies of the ferro- and nonferromagnetic states is even larger than in 3D. Therefore, it looks like that in low dimensions making the ground state ferromagnetic is harder than in 3D [17]. However, in this paper we reveal that this can be done in a 1D Fermi gas. We use interactions that go far beyond the Stoner model and show that they can be realized with cold atoms, where one has a remarkable system of ^{40}K. It is characterized by a proximity of the s-wave (even wave in 1D) Feshbach resonance for the intercomponent interaction and p-wave (odd wave in 1D) resonance for one of the intracomponent interactions. Thus one can have simultaneously a strong or even infinite intercomponent repulsion and a significant momentum-dependent odd-wave interaction in one of the components. It is the latter one that drastically changes the situation and makes the ground state ferromagnetic.

The case of ^{40}K is really a "present from nature." The s-wave resonance for the interaction between $9/2,-7/2$ and $9/2,-9/2$ states occurs at a magnetic field of 202.1 G, and is very close to the p-wave resonance for the interaction between two $9/2,-7/2$ atoms at 198.8 G [18,19]. In the fields between 198.8 and 202.1 G the p-wave interaction is attractive, and the s-wave interaction is repulsive and can be made very strong by using a confinement-induced resonance in 1D. Moreover, the reduction of dimensionality to 1D decreases the inelastic decay even not far from the resonances, which is very promising for achieving itinerant ferromagnetism in a 1D gas of ^{40}K.

The odd-wave scattering of identical fermions occurs in the triplet spin channel, where the spin part of the wave function is symmetric and the coordinate part antisymmetric. If the odd-wave interaction is the same in all triplet states, and the even-wave repulsion (occurring in the singlet spin channel) is infinitely strong, then the problem is exactly solvable and can be mapped onto two-component bosons with SU(2) spin rotation symmetry. For the latter case the ground state is ferromagnetic [20–22]. However, the spin rotation symmetry breaks if the odd-wave interaction is nearly resonant, since it then depends on the spin projections of colliding particles. In this regime the exact solution is no longer available and we have to employ many-body perturbation theory.

We consider a 1D Fermi gas in free space and assume that the intercomponent contact interaction is infinitely repulsive. It takes place between two particles with zero total (pseudo)spin and is present only in the nonferromagnetic phases. Therefore, omitting the odd-wave interaction, the ferromagnetic phase represents an ideal single-component Fermi gas, with the Fermi momentum $k_F = \pi n$. The total energy E_f is equal to the kinetic energy $E_{\rm kin} = E_F N/3$, where n is the 1D density, N is the total number of particles, m is the particle mass, and $E_F = \hbar^2 k_F^2/2m$ is the Fermi energy. The nonferromagnetic phases in this case are described by the exactly solvable Yang-Gaudin model [23,24]. A finite contact repulsion leads to the antiferromagnetic ground state, which is a singlet-pair correlated phase. For an infinite repulsion all spin configurations are degenerate with the energy equal to $E_{\rm kin}$ [23–25].

We are interested in the regime where the odd-wave interaction is nearly resonant. Since for the most important case of ^{40}K atoms the resonance in the odd-wave channel is present only between two atoms in the $9/2,-7/2$ states we confine ourselves to the odd-wave interaction between these states. Below the state $9/2,-7/2$ is denoted as spin-$\uparrow$, and the state $9/2,-9/2$ as spin-$\downarrow$. Moreover, we assume that although the odd-wave interaction is nearly resonant, it is not too strong (a more precise condition will be given later), and still can be treated as perturbation. The fact that one can use a perturbative approach in a 1D odd-wave interacting system (in contrast to the even-wave interaction) finds its origin in the absence of a weakly bound state in a sufficiently shallow attractive potential. For the ferromagnetic many-body system our perturbative results perfectly agree with the existing Bethe ansatz solution [26].

Thus, to zero order the kinetic energy $E_{\rm kin}$ is the same in any spin configuration (as a consequence of the infinite even repulsion) and gives the main contribution to the total energy E of the system, while the odd-wave interaction provides a small correction, which we derive up to the second order in perturbation theory. For the nonferromagnetic phases we employ the single-component momentum distribution functions $N_\uparrow(k)$ and $N_\downarrow(k)$ that we obtain by solving numerically the Bethe ansatz equations for the Yang-Gaudin model at an infinite intercomponent repulsion [27]. Considering equally populated $\uparrow$ and $\downarrow$ internal states, $N_\uparrow(k) = N_\downarrow(k) = N(k)$, we have $\int_{-\infty}^{+\infty} N(k)dk/2\pi = n/2$ and

$$\int_{-\infty}^{+\infty} \frac{L\,dk}{2\pi}\frac{\hbar^2k^2}{2m}N(k) = E_{\rm kin}/2 = \frac{\pi^2\hbar^2n^2}{12m}N, \tag{1}$$

with L being the size of the system. For the ferromagnetic phase we use the Fermi step momentum distribution $N(k) = \theta(k_F - |k|)/2$.

In order to develop many-body perturbation theory we follow the method used in Refs. [28,29]. We define the off-shell scattering amplitude as

$$f(k',k) = \int_{-\infty}^{\infty} dx\, e^{-ik'x}\, V(x)\,\psi_k(x), \tag{2}$$

where $V(x)$ is the interaction potential, $\psi_k(x)$ is the true wave function of the relative motion with momentum $k = (k_1 - k_2)/2$, with k_1, k_2 and k_1', k_2' being the particle momenta in the incoming and outgoing channels. For $|k'| = |k_1' - k_2'|/2 = |k|$ we have the on-shell amplitude. The total energy is $E = E_{\rm kin} + \tilde{E}^{(1)} + \tilde{E}^{(2)}$, where the first- and second-order corrections are given by [27]

$$\tilde{E}^{(1)} = \frac{1}{L}\sum_{k_1,k_2}\tilde{f}_{\rm odd}(k)N(k_1)N(k_2), \tag{3}$$

$$\tilde{E}^{(2)} = -\frac{1}{L^2}\sum_{k_1,k_2,k_1'}\frac{4m}{\hbar^2}\frac{\tilde{f}_{\rm odd}(k',k)\tilde{f}_{\rm odd}(k,k')}{k_1^2+k_2^2-k_1'^2-k_2'^2} \times N(k_1)\,N(k_2)\,N(k_1'), \tag{4}$$

with $k_1 + k_2 = k_1' + k_2'$. The amplitude $\tilde{f}_{\rm odd}$ is different from the odd-wave part of (2) by the absence of the imaginary term in the denominator [27]. The terms $\tilde{E}^{(1)}$ and $\tilde{E}^{(2)}$ are the two-body (mean-field) and the many-body, or beyond mean-field, contributions to the interaction energy. As the 1D regime is obtained by tightly confining the motion of particles in two directions to zero point oscillations, the odd-wave off-shell scattering amplitude in the vicinity of the resonance is given by [27,30]

$$\tilde{f}_{\rm odd}(k',k) = \frac{2\hbar^2}{m}\frac{k'k\,l_p}{1+\xi_p l_p\,k^2}, \tag{5}$$

where the parameters l_p and ξ_p of the 1D odd-wave scattering can be expressed through the parameters of the 3D p-wave scattering as $l_p = 3a_\perp/(a_\perp^3/w_1 + 0.88)$; $\xi_p = \alpha_1 a_\perp^2/3$, where w_1 and α_1 are the 3D scattering volume and effective range, and $a_\perp = \sqrt{\hbar/(m\omega_\perp)}$ is the extension of the wave function in the directions tightly (harmonically) confined with frequency $\omega_\perp$. For ^{40}K atoms near the p-wave Feshbach resonance the magnetic field dependence of w_1 and α_1 has been measured in the JILA experiments [31]. Near the resonance in 3D the scattering volume w_1 changes from infinitely negative to infinitely positive, whereas the effective range α_1 remains practically constant and equal to $4\times10^6\ {\rm cm}^{-1}$. On the positive side of the resonance ($w_1 > 0$) in 3D one has the formation of rapidly decaying p-wave molecules [19,32–35], and a similar phenomenon is expected in 1D. The issue of inelastic losses is discussed in more detail below, but in what follows we consider only the case of attractive odd-wave interaction, i.e., $l_p < 0$. The energy corrections (3) and (4) can be rewritten as [27]

$$\tilde{E}^{(1)} = -E_{\rm kin}\left\{\frac{1}{2\pi}\eta + \frac{3}{16\pi}\kappa\,\eta^2\,\mathcal{I}(Q)\right\}, \tag{6}$$

$$\tilde{E}^{(2)} = \frac{3}{4\pi^2}\eta^2\mathcal{J}(Q)E_{\rm kin}, \tag{7}$$

where $\eta = k_F|l_p|$, $\kappa = k_F\xi_p$, $Q = \eta\kappa$, and we took into account that in any spin configuration the momentum distribution

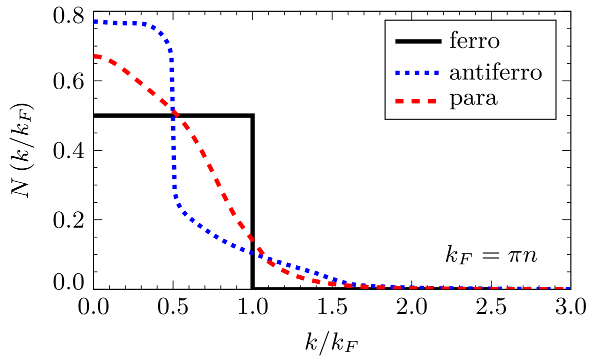

FIG. 1. Momentum distributions in the ferro-, antiferro-, and paramagnetic phases. The latter two are for the Yang-Gaudin model at an infinite repulsion.

is a universal function of k/k_F. The integrals $\mathcal{I}(Q)$ and $\mathcal{J}(Q)$ are given by

$$\mathcal{I}(Q) = \mathrm{P}\int_{-\infty}^{+\infty} dx_1 dx_2 N(x_1)N(x_2)\frac{(x_1-x_2)^4}{1-\frac{Q}{4}(x_1-x_2)^2}, \quad (8)$$

$$\mathcal{J}(Q) = \mathrm{P}\int_{-\infty}^{+\infty} dx_1 dx_2 dx_3 \frac{N(x_1)N(x_2)N(x_3)}{x_1-x_3} \times \frac{(x_1-x_2)^2}{\left[1-\frac{Q}{4}(x_1-x_2)^2\right]}\frac{x_1+x_2-2x_3}{\left[1-\frac{Q}{4}(x_1+x_2-2x_3)^2\right]}, \quad (9)$$

with $x_i = k_i/k_F$ being a dimensionless momentum, and the symbol P denoting the principal value of the integral. The choice of a particular spin configuration is encoded in the momentum distribution $N(k/k_F)$, and from Eqs. (6)–(9) it is evident that the odd-wave interaction splits the energies of different phases only if $\kappa \neq 0$. The unperturbed momentum distributions for the ferro-, antiferro-, and paramagnetic states are displayed in Fig. 1. At $k \gg k_F$ the momentum distributions in the nonferromagnetic states behave as $N(k) \to \mathcal{C}/k^4$, where $\mathcal{C}$ is Tan's contact [36,37]. In the antiferromagnetic phase we have $\mathcal{C}/k_F^4 = 2\ln(2)/3\pi^2 \approx 0.047$ [37], and in the paramagnetic phase we obtain $\mathcal{C}/k_F^4 \approx 0.016$.

For $Q \to 0$ we have $\mathcal{J} = 1/2$ and $\mathcal{I} = 2(1+3\mathcal{D})/3$, where $\mathcal{D} = \int_{-\infty}^{\infty} dx[N(x)x^4 - \mathcal{C}/k_F^4]$. For a finite Q we calculate the integrals (8) and (9) numerically.

Realization of the 1D regime requires the Fermi energy to be much smaller than the tight confinement frequency. For realistic confinement frequencies $\omega_\perp$ in the range from 50 to 150 kHz, the condition $E_F \ll \hbar\omega_\perp$ requires the Fermi momentum $k_F \lesssim 10^5$ cm^{-1} (which corresponds to densities $n \lesssim 3\times 10^4$ cm^{-1} and $E_F \lesssim 1$ μK). The confinement length $a_\perp$ for such frequencies is from 400 to 700 Å. Then, taking the potassium value $\alpha_1 \simeq 4\times 10^6$ cm^{-1} we see that the parameter $\kappa = \pi n \alpha_1 a_\perp^2/3$ ranges from 1 to 10. In the perturbative regime we require $\eta/\pi = n|l_p| \ll 1$, i.e., one should not be too close to the resonance, and in order to stay within the limits of perturbation theory we put $\eta < 0.8$.

We then calculate the total energy of the gas up to the second order in perturbation theory for the ferro-, para-, and antiferromagnetic phases. The results are presented in Fig. 2

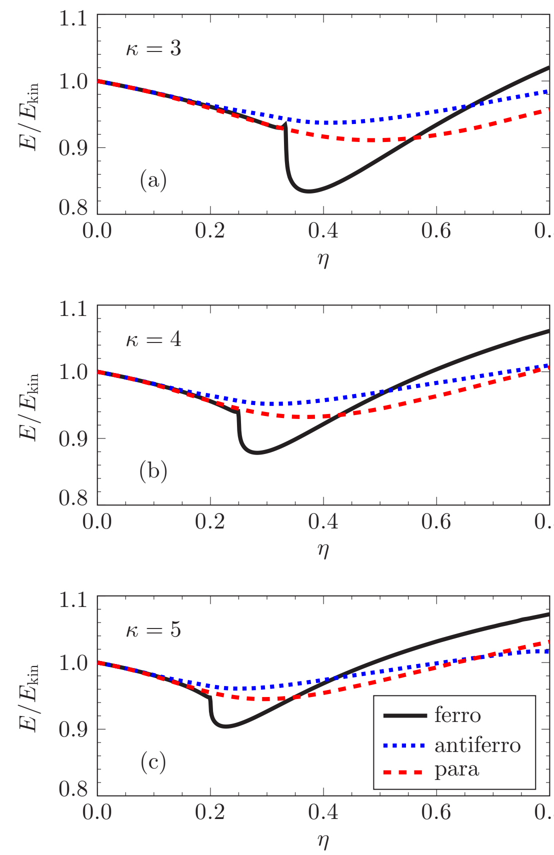

FIG. 2. The total energy (in units of $E_{\mathrm{kin}} = E_F N/3$) versus η in the ferro-, antiferro- and paramagnetic state for $\kappa = 3$, $\kappa = 4$, and $\kappa = 5$.

and can be explained as follows. The main contribution to the interaction energy comes from the first order correction given by Eq. (6). There is a characteristic momentum $\tilde{k} = 1/\sqrt{|l_p|\xi_p} = k_F/\sqrt{\eta\kappa}$ above which the odd-wave interaction turns from attractive to repulsive. For sufficiently small η (large $\tilde{k}$) only momenta $k \ll \tilde{k}$ contribute to the interaction energy, the term $\xi_p l_p k^2$ in the denominator of the interaction amplitude (5) is not important and all states have the same energy (left part in Fig. 2). For η such that $\tilde{k}$ is close to k_F, the main contribution to the interaction energy comes from $k \sim \tilde{k}$. In the ferromagnetic phase the interaction remains attractive if $\tilde{k} > k_F$ and becomes repulsive only in a narrow momentum interval $\tilde{k} < k < k_F$ if $\tilde{k} < k_F$. On the contrary, in the nonferromagnetic phases the distribution function $N(k)$ extends essentially to momenta k greater than $\tilde{k} \sim k_F$. This makes the repulsive contribution to the interaction energy larger than in the ferromagnetic phase, and the latter becomes the ground state (central part of Fig. 2). For $\tilde{k}$ significantly smaller than k_F, achieved for example by increasing η, the contribution of repulsive interactions dominates. It is the largest at momenta in the interval $k_F/2 \lesssim |k| \lesssim k_F$ where the distribution function $N(k)$ in the ferromagnetic state is larger than in nonferromagnetic phases. This makes the ferromagnetic energy the highest (right part of Fig. 2).

For a gas of ^{40}K atoms with a density $n \approx 3\times 10^4$ cm^{-1} ($E_F \approx 540$ nK) under the transverse confinement with frequency $\omega_\perp \approx 100$ kHz we have $\kappa \approx 3.1$, and at fields slightly lower than 199 G we obtain $\eta \approx 0.36$. Then the ferromagnetic state has the lowest energy, and the energies of nonferromagnetic states are close to each other. The energy difference $(E_p - E_f)/N$ is about $0.03E_F$ or 16 nK. For $\omega_\perp \approx 120$ kHz we obtain $\kappa \approx 2.6$, and with $\eta \approx 0.43$ at $B \approx 199$ G the energy difference is $(E_p - E_f)/N \approx 20$ nK. Thus, the ferromagnetic state can be observed at temperatures below 20 nK.

The regimes described above ensure that the even repulsion is infinitely strong, even though the corresponding magnetic fields are not too close to the s-wave resonance. However, in the 1D geometry obtained by tightly confining particles in two directions, the coupling constant for the even contact interaction is $g_{1D} = (2\hbar^2 a/ma_\perp)/(a_\perp - 1.03a)$ [38]. For confining frequencies of 100 and 150 kHz the length $a_\perp$ is 500 and 400 Å, respectively. In a field close to 199 G the scattering length is $a \approx 400$ Å. Thus, due to the confinement-induced resonance, one can achieve an infinite contact repulsion $g_{1D} \to \infty$.

In three dimensions p-wave Feshbach resonances are suffering rapid inelastic losses [19,32–35]. There are two types of inelastic collisional processes. The first one is three-body recombination, which is especially pronounced if there are weakly bound dimer states. However, weakly bound dimer p-wave states are expected only on the positive side of the resonance ($l_p > 0$), and on the negative side ($l_p < 0$) the three-body recombination should not be very dangerous, at least slightly away from the resonance. The absence of weakly bound p-wave states also prohibits the formation of dimers in two-body collisions, where the released binding energy goes to the creation of holes in the Fermi sea [15,39]. Another decay process is two-body relaxation. The state $9/2,-7/2$ can undergo collisional relaxation to the $9/2,-9/2$ state which has a lower energy.

In fields slightly higher than 199 G the measured rate constant of three-body recombination in 3D is $\alpha_{\rm rec}^{3D} \sim 10^{-25}$ cm^6/s and the rate constant of two-body relaxation is $\alpha_{\rm rel}^{3D} \sim 10^{-14}$ cm^3/s [19,40]. The measurements were done at temperatures from 1 to 3 μK, so that one expects about the same rate constants at $E_F \sim 1$ μK and much lower temperatures. In order to transform these results to 1D one should recall that the inelastic processes occur at atomic interparticle distances. We thus may integrate out the motion in the tightly confined directions [41]. This leads to $\alpha_{\rm rel}^{1D} \approx \alpha_{\rm rel}^{3D}/2\pi a_\perp^2$ [42]. With the above specified $\alpha_{\rm rel}^{3D}$ and densities $n \sim 10^4$ or 3×10^4 cm^{-1} we obtain a relaxation time of about a second.

For the three-body recombination of identical fermions in 1D one has an extra suppression by a factor of E_F/E_* compared to 3D [43,44]. The quantity E_* is a typical energy in the molecular problem, and one has $E_* \sim \hbar^2/mR_e^2$, where $R_e \sim 50$ Å is the radius of interaction between particles. So, $E_* \sim 1$ mK and there is an extra suppression by 3 orders of magnitude for $E_F \sim 1$ μK. Integrating out the particle motion in the tightly confined direction we obtain $\alpha_{\rm rec}^{1D} \sim \alpha_{\rm rec}^{3D}(E_F/E_*)/3\pi^2 a_\perp^4$ [42]. Again, for the above mentioned parameters we obtain a decay time of about a second at 1D densities in between 10^4 and 3×10^4 cm^{-1}. The ferromagnetic state can be viewed as a composition of identical fermions, and this estimate remains valid in this phase. Thus, in 1D inelastic decay processes are not as crucial as in 3D. In this respect the situation is somewhat similar to the one with strongly interacting bosons [45,46].

In conclusion, we showed that there is a realistic possibility to find itinerant ferromagnetic states in 1D two-component Fermi gases, and a promising system is the gas of ^{40}K atoms. This will require fine tuning of the interaction between particles by varying the magnetic field and the confinement strength. The required temperatures are about 10 to 20 nK, which is achievable with present facilities.

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