Stability of superfluid phases in the 2D spin-polarized attractive
Hubbard model

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Abstract – We study the evolution from the weak coupling (BCS-like limit) to the strong coupling limit of tightly bound local pairs (LPs) with increasing attraction, in the presence of the Zeeman magnetic field (h) for d = 2, within the spin-polarized attractive Hubbard model. The broken symmetry Hartree approximation as well as the strong coupling expansion are used. We also apply the Kosterlitz-Thouless (KT) scenario to determine the phase coherence temperatures. For spin-independent hopping integrals (t^↑ = t^↓), we find no stable homogeneous polarized superfluid (SC_M) state in the ground state for the strong attraction and obtain that for a two-component Fermi system on a 2D lattice with population imbalance, phase separation (PS) is favoured for a fixed particle concentration, even on the LP (BEC) side. We also examine the influence of spin-dependent hopping integrals (mass imbalance) on the stability of the SC_M phase. We find a topological quantum phase transition (Lifshitz type) from the unpolarized superfluid phase (SC_0) to SC_M and tricritical points in the h-|U| and t^↑/t^↓-|U| ground-state phase diagrams. We also construct the finite temperature phase diagrams for both t^↑=t^↓ and t^↑≠t^↓ and analyze the possibility of occurrence of a spin-polarized KT superfluid.

First theoretical studies of Fermi condensates in systems with spin and mass imbalances have shown that the BP state can have excess fermions with two FSs at T = 0 (BP-2 or interior gap state) [5–7]. However, the problem of stability of the BP-2 state is still open. According to some investigations, the interior gap state proposed by Liu and Wilczek [5] is always unstable even for large mass ratio r and PS is favoured [8].

At T = 0, for strong attraction, the SC_M phase occurs in the three-dimensional imbalanced Fermi gases [4,8] as well as in the 3D spin-polarized attractive Hubbard model in the dilute limit (for r = 1 and r ≠ 1 [9]). The SC_M phase is a specific superfluid state consisting of a coherent mixture of LPs (hard-core bosons) and excess spin-up fermions (Bose-Fermi mixture). This state can only have one FS.

In this paper we study the superfluid phases in the attractive Hubbard model (AHM) (U < 0) in a magnetic field with spin-dependent hopping:

\[ H = \sum_{ij\sigma}(t_{ij}^{\sigma} - \mu \delta_{ij})c_{i\sigma}^\dagger c_{j\sigma} + U \sum_i n_{i\uparrow} n_{i\downarrow} - h \sum_i (n_{i\uparrow} - n_{i\downarrow}), \]

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Introduction. – Unconventional superconductivity with non-trivial Cooper pairing and spin-polarized superfluidity have been investigated in recent years. The development of experimental techniques in cold atomic Fermi gases with tunable attractive interactions (through Feshbach resonance) has allowed the study of the BCS-BEC crossover and the properties of exotic states in these systems [1].

The presence of a magnetic field (h), population imbalance or mass imbalance (spin-dependent hopping integrals (t^↑ ≠ t^↓, \( t^\prime / t^\dagger \equiv r \) – mass ratio) introduces a mismatch between Fermi surfaces (FS). This makes possible the formation of Cooper pairs across the spin-split Fermi surface with non-zero total momentum (κ↑, –κ↑ + q↓) (Fulde, Ferrell, Larkin and Ovchinnikov [2] (FFLO) state). Another kind of pairing and phase coherence that can appear is the spatially homogeneous spin-polarized superfluidity (SC_M) (called breached pair (BP) state or Sarma phase [3]), which has a gapless spectrum for the majority spin species [4].
\[ \rho_s(T) = \frac{1}{4N} \sum_\vec{k} \left\{ \frac{\partial^2 \epsilon^+_\vec{k}}{\partial k_x^2} + \frac{\partial^2 \epsilon^-_\vec{k}}{\partial k_y^2} \right\} + \frac{1}{2} \left( \frac{\partial^2 \epsilon^+_\vec{k}}{\partial k_x^2} + \frac{\partial^2 \epsilon^-_\vec{k}}{\partial k_y^2} \right) + \left( \frac{\partial \epsilon^+_\vec{k}}{\partial k_x} \right)^2 \left( \frac{\partial \epsilon^-_\vec{k}}{\partial k_y} \right)^2 \tan \left( \frac{\beta E_{k^\perp}}{2} \right). \]

where: \( \sigma = \uparrow, \downarrow \), \( n_{i\uparrow} = c_{i\uparrow}^{\dagger} c_{i\uparrow} \), \( n_{i\downarrow} = c_{i\downarrow}^{\dagger} c_{i\downarrow} \), \( \epsilon^\sigma_l \) are the hopping integrals, \( U \) the on-site interaction, \( \mu \) the chemical potential. The gap parameter is defined by: \( \Delta = -\frac{\mu}{N} \sum_\vec{k} \sigma_\vec{k} \). Applying the broken symmetry Hartree (BCS-Stoner) approximation, we obtain the grand canonical potential and the free energy \( F \). Using the free energy expression, one gets the equations for the gap:

The KT temperature \( T_K \) is mapped (via the canonical transformation) to the unit.

We also calculate the superfluid density \( \rho_s(T) \) which for \( \epsilon^\sigma_l \) takes the form

\[ \rho_s(T) = \frac{1}{4N} \sum_\vec{k} \left\{ \frac{\partial^2 \epsilon^+_\vec{k}}{\partial k_x^2} + \frac{\partial^2 \epsilon^-_\vec{k}}{\partial k_y^2} \right\} + \left( \frac{\partial \epsilon^+_\vec{k}}{\partial k_x} \right)^2 \left( \frac{\partial \epsilon^-_\vec{k}}{\partial k_y} \right)^2 \tan \left( \frac{\beta E_{k^\perp}}{2} \right). \]

and: \( n = \frac{1}{N} \sum_i (2\rho_i^+ + 1) \), \( J_0 = \sum_{ij} J_{ij} \), \( J_{ij} = 2 \frac{t_{ij}^1 t_{ij}^2}{|\vec{r}_i - \vec{r}_j|} \), \( K_{ij} = 2 \frac{(t_{ij}^1)^2 + (t_{ij}^2)^2}{|\vec{r}_i - \vec{r}_j|} \), \( \bar{\mu} = \mu + \frac{U}{2} \), the pseudo-spin operators are: \( \rho_i^e = c_{i\uparrow}^{\dagger} c_{i\downarrow} \), \( \rho_i^o = c_{i\downarrow}^{\dagger} c_{i\uparrow} \), \( \rho_i^t = \frac{1}{2}(n_i^t + n_i^o - 1) \), the pseudo-spin operators satisfy the commutation rules of \( s = \frac{1}{2} \) operators. It is worth noting that Hamiltonian (4) operates on the unit.

With the spin-dependent hopping integrals, it is possible that the charge density wave ordered (CO) state can develop for any particle concentration. The SC to CO is a first-order transition at \( h = 0 \), \( t^t \neq t^o \) and \( n \neq 1 \). The critical \( n_{c} \) (within the mean-field approximation) above which superconductivity can coexist with commensurate CO is given by [12,15]: \( n_{c} = 1 - \frac{\sqrt{2}}{K_0 + K_0^*} \). Substituting expressions for \( J \) and \( K \), one obtains

\[ n_{c} = 1 \pm \frac{r - 1}{r + 1}. \]

Away from half filling the quantum fluctuations can extend the region of stability of the SC phase at \( T = 0 \) and enhance \( n_{c} \) [14,16]. In further considerations we fix \( n < n_{c} \) and consider mostly low \( n \).

Results. – We have performed an analysis of the evolution of superconducting properties from the weak to the strong coupling limit with increasing \( |U| \), within the spin-polarized AHM. The system of self-consistent equations has been solved numerically for the 2D square lattice. First, the chemical potential has been fixed. The first-order transition lines were determined from the condition \( \Omega^{SC} = \Omega^{NO} \). Next, \( \Omega^{SC}, \Omega^{NO} \) is the grand canonical potential of superconducting and normal state, respectively. Then, these results have been mapped onto the case of fixed \( n \). The stability of all phases has been examined very thoroughly. In the following, we use \( t \) as the unit.

According to the Leggett criterion [17], the BCS-BEC crossover takes place when the chemical potential \( \bar{\mu} \) reaches the lower band edge. In the limiting case of two
The diagrams for an arbitrary value of on-site attraction $U$ consist of two critical fields defining the PS region. The green point shows the BCS-LP crossover point $(U/t = -0.01959)$. For $d = 2$ AHM at $h = 0$, the SC-NO transition is of the KT type, i.e., below $T_c^{KT}$ a system has a quasi-long-range (algebraic) order which is characterized by a power law decay of the order parameter correlation function and non-zero superfluid stiffness. According to eq. (3), the KT transition temperature is found from the intersection point of the straight line $\frac{\hbar}{2}k_BT$ with the curve $\rho_\pi(T)$. In such a way, we can estimate the phase coherence temperatures and extend the analysis of crossover from weak to strong coupling to finite $T$ [21].

Fig. 2 shows $T$-$E_b$ phase diagrams in units of Fermi energy, for fixed values of magnetic field: (a) $h/E_F = 0.35$ (on the BCS side), (b) $h/E_F = 1$ (in the intermediate couplings) and (c) the $T$-$P$ diagram at fixed $E_b/E_F = 0.131$. The solid lines (2nd-order transition lines) and PS regions are obtained within the Hartree approximation. The thick dash–double-dotted line (red color) denotes the KT transition determined from eqs. (2), (3). The system is a quasiparticle quasistate (qSC) below $T_c^{KT}$. The polarization can be induced by thermal excitations of quasiparticles. Above $T_c^{KT}$ but below $T_c^{HF}$ (pair breaking temperature), pairs exist but without a long-range phase coherence. In this region one has a pseudo-gap behaviour. The temperatures $T_c^{KT}$ are generally much smaller than $T_c^{HF}$, but reducing the attraction, in the absence of magnetic field, the difference between $T_c^{KT}$ and $T_c^{HF}$ decreases in a weak coupling limit. At $h = 0$ and $E_b \ll 1$, $T_c^{HF} \neq 0$ and $T_c^{KT} \neq 0$. When the magnetic field increases, $T_c^{HF} = 0$, below a definite value of the binding energy. This critical $E_b$ increases with $h$ (fig. 1). In the strict BCS-MFA diagram the tricritical point (TCP) exists at finite magnetic fields. In this mean-field TCP, the SC MF phase, the NO state and PS coexist. There is also “TCP” on the KT curve in which three states meet: the qSC phase, the state of incoherent pairs and PS (fig. 2(a), (b)). The PS range widens with increasing $h$ and the distance between the MF TCP and KT “TCP” is larger. As shown in fig. 2(c), the effect of finite $P$ on the KT superfluid state is strong. For $r = 1$, the KT phase is restricted to
the weak coupling region and low values of \( P \). Increasing polarization favours the phase of incoherent pairs. The range of occurrence of qSC in the presence of \( P \) widens in the weak coupling regime with increasing \( n \). In turn, the qSC state is highly reduced with increasing attractive interaction even for low population imbalance.

By the analysis of the quasiparticle excitation spectrum, we also find a gapless region (yellow color in diagrams), for \( h > \Delta \) (the BCS side) and for \( h > E_g/2 \), where \( E_g = 2\sqrt{\langle \bar{\mu} - \epsilon_0 \rangle^2 + \Delta f^2} \) (on the LP side). If \( r = 1 \) this gapless region can only be realized at \( T > 0 \) and has excess fermions with two FS in the weak coupling limit. In fig. (2) the gapless region is distinguished within the state of incoherent pairs, i.e. is non-superfluid. In the strong coupling regime, the temperature can induce the spin-polarized gapless region (in a state of incoherent pairs) with one FS. This is in contrast to the 3D case where the BP-1 phase can be stable even without mass imbalance at \( T = 0 \) and low \( n \) [9]. In the strong coupling limit, \( T^*_c \) does not depend on magnetic field and approximately approaches \( k_B T^*_c = E_F \approx \frac{\hbar}{T^*F} \left( 1 - \frac{\Delta}{T^*_F} \right) \) for \( |U| \gg t \), \( E_F = 2tn \), in contrast to the continuum model which yields in that limit \( k_B T^*_c/E_F = \frac{1}{2} \) [21]. For \( k_B T \ll |U| \), there exist only LPs which are not broken by the magnetic field and the system is equivalent to that of a hard-core Bose gas on a lattice. The thin dash–double-dotted line in fig. 2(a), (b) inside the NO state marks the region where \( \Omega \) has two minima (below the curve): lower at \( \Delta = 0 \) and higher at \( \Delta \neq 0 \). It means that there can exist a metastable superconducting state.

An interesting aspect of the analysis is the influence of the Hartree term on the phase diagrams. First, the presence of the Hartree term leads to the re-entrant transition (RT) in the weak coupling limit (fig. 2(a)), which is not observed in the phase diagrams without the Hartree term. We also find a narrow region around MF TCP in which \( \rho_s < 0 \) although \( \Omega^{SC} < \Omega^{NO} \) (fig. 2(a) inset), in the phase diagram on the BCS side with the Hartree term. If RT exists, it becomes dynamically unstable because \( \rho_s < 0 \).

In addition, the Hartree term causes an increase in the Chandrasekhar-Clogston limit [10], [22].

iii) \( r \neq 1, h \neq 0, T = 0 \). The BCS-LP crossover diagrams in the presence of a Zeeman magnetic field for \( r \neq 1 \) exhibit a novel behaviour. As opposed to the \( r = 1 \) case, for strong attraction, SC\(_M\) occurs at \( T = 0 \) (fig. 3). In general, these types of solutions (Sarma-type with \( \Delta(h) \) appearance) do not exist. Hence, the SC\(_M\) phase is the superfluid state of coexisting LPs (hard-core bosons) and single-species fermions, with the latter responsible for finite polarization (magnetization) and the gapless excitations characteristic for this state of Bose-Fermi mixture.

The structure of the ground-state diagrams in fig. 3 is different from those shown in fig. 1, where one has only a first-order phase transition from the pure SC\(_0\) to the NO phase in the \( \mu-h \) plane. In addition, there exist critical values of \(|U|\) (red points in the diagrams), for which the SC\(_M\) state becomes stable, instead of PS. However, one should mention that there is a critical value of \( r \), for which SC\(_M\) is stable. Figure 3(a) shows the ground-state \((r-E_b/E_F)\) phase diagram for fixed \( n = 0.01 \) and \( h = 0 \). The SC\(_M\) state does not appear stable up to \( r \approx 5 \) in the diagram with the Hartree term and also up to \( r \approx 4.2 \) in the diagram without the Hartree term. The presence of such a term restricts the range of occurrence of SC\(_M\), except for a very dilute limit.

The diagrams \( h-|U| \) for higher filling (\( n = 0.2 \)) and fixed \( r \) are presented in fig. 3(b), (c). For higher \( n \), the region of SC\(_M\) is narrowing. The SC\(_M\) state is unstable even at \( r = 5 \), in the diagram without the Hartree term. The transition from SC\(_M\) to FP can be accomplished in two ways: through PS-III (SC\(_M\) + FP) or through a 2nd-order
Fig. 3: (Colour on-line) Ground-state phase diagrams for the $d=2$ square lattice: (a) mass imbalance vs. $E_b/E_F$ (where $E_F$ is the lattice Fermi energy of unpolarized, non-interacting fermions with hopping $t$) for $n=0.01$ and $h=0$ with and without Hartree term (inset). (b), (c) magnetic field vs. attractive interaction for two values of $r$ and $n=0.2$, without Hartree term. SC: magnetized superconducting state, PP: partially polarized state, FP: fully polarized state, PS-I: (SC$_0$ + PP), PS-II: (SC$_0$ + PP), PS-III: (SC$_M$ + PP). Red points: $k_{B}T_{SC}^{M}$, blue points: tricritical points, green points: the BCS-LP crossover points in the SC$_0$ phase ($r = 1$). The dotted red and the solid green lines are the 2nd-order transition lines.

Fig. 4: (Colour on-line) (a) Ground-state phase diagram $r$ vs. $|U|$ for $n = 0.1$, $h = 0$, (b-c) temperature vs. $|U|$ phase diagrams, at fixed $n = 0.1$, $h = 0$ for the square lattice. (b) $r = 4$, (c) $r = 7$ (inset: $P(T_{KT}^{c})$ vs. $|U|$). The thick dashed-dotted line (red color) is the KT transition line. The thick solid line in (c) denotes transition from pairing without coherence region to FP within the BCS approximation, the dotted line (red color) is the second-order transition line from SC$_0$ to SC$_M$ state at $T=0$ (a) or from qSC$_0$ to the gapless region (yellow color) at $T \neq 0$ (c). PS: phase separation, qSC$_0$: qSC without polarization, qSC$_M$: qSC in the presence of polarization (a spin-polarized KT superfluid). Red point at $T=0$ (see (c)): QCP for Lifshitz transition.

phase transition. The character of this transition changes with increasing $|U|$. In the very strong coupling limit, PS is more stable than the SC$_M$ phase. Hence, we also find two TCP in these diagrams (fig. 3(c)). However, in the very dilute limit ($n \to 0$) there is only one TCP in the $h$-$|U|$ diagram. Therefore, we can distinguish the following sequences of transitions: SC$_0$ $\to$ PP (FP) or SC$_0$ $\to$ SC$_M$ $\to$ FP.

In fact, SC$_0$ $\to$ SC$_M$ is a topological quantum phase transition (Lifshitz type). Across this transition there is a cusp in the order parameter and the chemical potential vs. magnetic field plots. There is also a change in the electronic structure. In the SC$_0$ phase, there is no FS, but in the SC$_M$ state, there is one FS for excess fermions. It is worth mentioning that the value of $|U|$ for which $\mu$ reaches the lower band edge does not depend on the mass imbalance in the SC$_0$ state. The BP-2 phase in $d=2$ is unstable, even for large mass ratios. If $r \neq 1$, the symmetry with respect to $h=0$ is broken. However, this symmetry is restored upon replacement $r \to r^{-1}$.

iv) $r \neq 1$, $h=0$, $T \neq 0$. Figure 4 shows the $(r - |U|)$ ground-state diagram (a) and also the $T$-$|U|$ phase diagrams, for fixed $h=0$, $r=4$ (b) and $r=7$ (c). At $T=0$, $r \neq 1$, we have the following sequences of transitions with increasing $|U|$ (see fig. 4(a)): FP $\to$ SC$_M$ $\to$ SC$_0$ (for higher values of $r$) or FP (PP) $\to$ PS $\to$ SC$_0$ (for lower values of $r$). Because of the occurrence of the SC$_M$ state at $T=0$ for higher $r$, this phase persists to non-zero temperatures (as shown in fig. 4(c), $r = 7$). However, if SC$_M$ is unstable at $T=0$ for lower $r$ (fig. 4(b)), the gapless region can still occur at some temperatures (with one FS in the strong coupling). The system is a quasiperiodic conductor below $T_{KT}^{c}$. Apart from unpolarized qSC$_0$ state, qSC$_M$ occurs which can be termed a spin-polarized KT superfluid (fig. 4(c)). Above $T_{KT}^{c}$ we have an extended region of incoherent pairs which is bounded from above by the pair breaking temperature. In the strong coupling limit, $T_{KT}^{c}$ does not depend on magnetic field, but it depends on mass imbalance and takes the form: $k_BT_{KT}^{c} = \frac{2\pi}{(1+r)^2} \frac{r}{|U|} n(2-n)$ ($r>0, \ n<n_c$). In
that limit only LPs exist and the system is equivalent to that of hard-core Bose gas on a lattice, described by the Hamiltonian (4). Because of the existence of the SC$_d$ state at $T=0$ in the $r \neq 1$ case, the range of $P$ for occurrence of a spin-polarized KT superfluid is much larger than for $r = 1$ (see inset in fig. 4(c)).

Conclusions. – We have investigated the evolution from the weak coupling to the strong coupling limit of tightly bound local pairs with increasing attraction for $d = 2$, within the spin-polarized AHM. If the number of particles is fixed and $n \neq 1$, one obtains two critical Zeeman magnetic fields, which limit PS of the SC$_0$ (or SC$_M$) and the NO states. The occurrence of the BP-1 phase depends on the lattice structure, i.e., if $r = 1$, SC$_M$ is unstable for $d = 2$ but it can be realized for $d = 3$ lattices [9]. However, in the AHM the very existence of the BP-1 phase is restricted to low fillings. The Hartree term, usually promoting ferromagnetism in the Stoner model $(U > 0)$, here $(U < 0)$ strongly competes with superconductivity. Thus, such a term restricts the SC$_M$ state to lower densities. However, the mass imbalance can change this behaviour even for $d = 2$ due to spin polarization stemming from the kinetic energy term. In this way, SC$_M$ can be realized in $d = 2$ for the intermediate and strong coupling regimes, but there is a critical value of the mass ratio for which SC$_d$ is stable, at finite fixed $n$. In other words, the combination of mass and population imbalance can stabilize the BP-1 phase in 2D, on the BEC side of crossover. We also determined the critical value of $n$ above which SC and CO can form PS state at $h \approx 0$ and $r \neq 1$. We have found that the BP-2 state is unstable in the whole range of parameters, in the $d = 2$ one-band spin-polarized AHM. Nevertheless, one can suppose that the Liu-Wilczek (BP-2) phase can be realized within the two-band model. The TCPs were found in the $h$--$|U|$ diagrams at $r \neq 1$ and $T = 0$. We have also extended the analysis of the crossover to finite temperatures in $d = 2$ by invoking the KT scenario. The KT transition temperatures are definitely lower than the ones determined in the BCS scheme. Moreover, spin polarization has a strong destroying influence on the KT superfluid state at $r = 1$ and allows this phase in the weak coupling regime, in agreement with the results for the continuum case [21]. In the strong coupling limit, $T_{KT}^c$ does not depend on magnetic field (below $h_{c1}$) and for $k_BT \ll |U|$ only unbroken LPs exist, which can form unpolarized qSC below $T_{KT}^c$ or stay phase disordered. However, if $r \neq 1$, a spin-polarized KT superfluid state can be stable even in the intermediate and strong coupling region. In this work we have not considered non-homogeneous FFLO states which are possible in weak to intermediate attraction range, albeit much more susceptible to phase fluctuations at finite $T$ in 2D system [23].

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REFERENCES

[1] ZWIEREIN M. W., ABO-SHAER J., SCHIROTZEK A., SCHUNCK C. and KETTERLE W., Nature, 435 (2005) 1047; ZWIEREIN M. W., SCHIROTZEK A., SCHUNCK C. H. and KETTERLE W., Science, 311 (2006) 492; PARTRIDGE G. B., LI W., KAMAR R. I., LIAO Y. and HULET R. G., Science, 311 (2006) 503; KETTERLE W., SHIN Y., SCHIROTZEK A. and SCHUNCK C. H., J. Phys.: Condens. Matter, 21 (2009) 164206.

[2] FULPE P. and FERRELL R. A., Phys. Rev., 135 (1964) A550; LARKIN A. I. and OVCHEINIKOV Y. N., Zh. Eksp. Teor. Fiz., 47 (1964) 1136.

[3] SARMA G., J. Phys. Chem. Solids, 24 (1963) 1029.

[4] SHEEHY D. E. and RADZHOVSKY L., Phys. Rev. Lett., 96 (2006) 060401.

[5] LIU W. V. and WILCZEK F., Phys. Rev. Lett., 90 (2003) 047002.

[6] LIU W. V., WILCZEK F. and ZOLLER P., Phys. Rev. A, 70 (2004) 033603.

[7] ISKIN M. and SÁ DE MELO C. A. R., Phys. Rev. Lett., 97 (2006) 100404; Phys. Rev. A, 76 (2007) 013601.

[8] PARISH M. M., MARCHETTI F. M., LAMACRAFT A. and SIMONS B. D., Phys. Rev. Lett., 98 (2007) 160402; Nat. Phys., 3 (2007) 124.

[9] KUJAWA-CICHY A., Acta Phys. Pol. A, 118 (2010) 423; KUJAWA-CICHY A. and MICNAS R., in preparation.

[10] KUJAWA A. and MICNAS R., Acta Phys. Pol. A, 114 (2008) 43; 115, (2009) 138.

[11] ROBASZKIEWICZ S., MICNAS R. and CHAO K. A., Phys. Rev. B, 23 (1981) 1447.

[12] MICNAS R., RANNINGER J. and ROBASZKIEWICZ S., Rev. Mod. Phys., 62 (1990) 113.

[13] CAZALILHA M. A., HO A. F. and GIAMARCHI T., Phys. Rev. Lett., 95 (2005) 226402 and references therein.

[14] MICNAS R. and ROBASZKIEWICZ S., Phys. Rev. B, 45 (1992) 9990; MICNAS R., ROBASZKIEWICZ S. and KOSTYRKO T., Phys. Rev. B, 52 (1995) 6863.

[15] DAO T.-L., GEORGES A. and CAPONE M., Phys. Rev. B, 76 (2007) 104517.

[16] YUNOKI S., Phys. Rev. B, 65 (2002) 092402.

[17] LEGGETT A. J., J. Phys. (Paris), Colloq., 41 (1980) C7-19.

[18] HE L. and ZHUANG P., Phys. Rev. A, 78 (2008) 033613.

[19] DU J.-J., CHEN C. and LIANG J.-J., Phys. Rev. A, 80 (2009) 023601.

[20] TEMPERE J., WOUTERS M. and DEVREESSE J. T., Phys. Rev. B, 75 (2007) 184526.

[21] For an analysis of the KT transition in 2D continuum Fermi gas with the s-wave pairing and population imbalance see: TEMPERE J., KLIMIN S. N. and DEVREESSE J. T., Phys. Rev. A, 79 (2009) 053637.

[22] CLOGSTON A. M., Phys. Rev. Lett., 9 (1962) 266.

[23] SHIMAHARA H., J. Phys. Soc. Jpn., 67 (1998) 1872; OHASHI Y., J. Phys. Soc. Jpn., 71 (2002) 2625; MATSUDA Y. and SHIMAHARA H., J. Phys. Soc. Jpn., 76 (2007) 051005.