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Enhancement of Superconductivity in Disordered Films by Parallel Magnetic Field

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We show that the superconducting transition temperature $T_c(H)$ of a very thin highly disordered film with strong spin-orbital scattering can be increased by parallel magnetic field H. This effect is due to polarization of magnetic impurity spins which reduces the full exchange scattering rate of electrons; the largest effect is predicted for spin-$\frac{1}{2}$ impurities. Moreover, for some range of magnetic impurity concentrations the phenomenon of *superconductivity induced by magnetic field* is predicted: superconducting transition temperature $T_c(H)$ is found to be nonzero in the range of magnetic fields $0 < H^* \le H \le H_c$.

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The problem of superconducting alloys with magnetic impurities was addressed long ago by Abrikosov and Gor'kov (AG) [1]. They have shown that superconductivity (SC) is suppressed due to exchange scattering (ES) of electrons on magnetic impurities, the transition temperature T determined from the equation (hereafter, we employ units, in which $\hbar = 1$):

$$\ln \frac{T_{c0}}{T} = \pi T \sum_{\varepsilon} \left(\frac{1}{|\varepsilon|} - \frac{1}{|\varepsilon| + \nu_S} \right). \qquad (1)$$

Here $\varepsilon = 2\pi T(m + 1/2)$ is the fermionic Matsubara frequency (m is integer), T_{c0} is the transition temperature of clean sample, and $\nu_S = 2\pi N_F n_S J^2 S(S+1)$ is the ES rate of electrons on magnetic impurities (N_F is the normal metal density of states per single spin state, n_S is the concentration of magnetic impurities, J is the exchange coupling constant, and S is the impurity spin length). The solution of (1) yields the function $T = T_{AG}(\nu_S)$. There exists a critical point at which the transition temperature is suppressed down to zero, the critical scattering rate being $\nu_S^* = \pi/(2e^C)\, T_{c0} = 0.882\, T_{c0}$, where $C = 0.577$ is the Euler constant. The critical concentration, corresponding to ν_S^*, is further denoted by n_S^*. We emphasize that ν_S is the *full* ES rate, i.e. the sum of the spin flip scattering rate $2\pi N_F n_S J^2(\langle S_x^2\rangle + \langle S_y^2\rangle) = 2/3\, \nu_S$ and the rate of scattering without spin flip $2\pi N_F n_S J^2 \langle S_z^2\rangle = 1/3\, \nu_S$.

The AG's results were derived for unpolarized magnetic impurity spins. In this Letter we investigate how the polarization of impurity spins affects the ES mechanism of SC suppression. We show that polarization of magnetic impurity spins by external magnetic field reduces the full ES rate $\Gamma(\varepsilon)$. It reaches its minimal value $\nu_\infty = \nu_S S/(S+1) < \nu_S$ at the infinite field, when the impurity spins are completely polarized and spin flip processes have frozen out. This reduction is due to quantum fluctuations of impurity spins, thus it is strongest for $S = 1/2$ and vanishes in the limit $S \gg 1$.

If ES was *the only* mechanism of SC suppression in nonzero magnetic field $h = \mu_B H$, the transition temperature $T_c^\circ(h)$ would always be higher than $T_c(h = 0) = T_{AG}(\nu_S)$, determined by AG's result (1). $T_c^\circ(h)$ is a growing function, approaching the value $T_\infty = T_{AG}(\nu_\infty)$ at very high fields $h \to \infty$. The transition temperature increase $T_c^\circ(h) - T_c(0)$ comparable to $T_\infty - T_c(0)$ is attained at the field range $h \gtrsim T_c^\circ(h)$. However, apart from ES, there are other mechanisms of SC suppression by magnetic field, namely, paramagnetic effect (PE) and orbital effect (OE). Thus, to observe an increase $T_c(h) > T_c(0)$ of the actual transition temperature, PE and OE should be small compared to ES in the field range $h \sim T_c(h)$. Strong reduction of PE is achieved in presence of high spin-orbital scattering rate $\nu_{\mathrm{so}} \gg T_{c0}$ [2 – 4]. OE is suppressed for a thin-film (thickness d shorter than the magnetic length $l_H = \sqrt{c/eH}$) with parallel orientation of external magnetic field [5].

In this Letter we show that the increase in the transition temperature can be observed if two quite stringent conditions on the smallness of PE and OE are met. First, the spin-orbit scattering rate ν_{so} must be sufficiently high:

$$\nu_{\mathrm{so}}/\nu_S \gg \zeta^2. \qquad (2)$$

Here $\zeta = n_S|J|S/\nu_S = (2\pi N_F|J|(S+1))^{-1} \gg 1$ is the inverse Born parameter for the exchange scattering. Second, elastic scattering rate ν and thickness of the film d must satisfy the condition

$$1 \lesssim (p_F d)^2 \ll \nu / T_{c0} \,, \qquad (3)$$

where p_F is Fermi momentum.

We distinguish between two different regimes depending on the value of ES rate ν_S. If $\nu_S < \nu_S^*$, i.e. there exist a finite transition temperature $T_c(0) = T_{AG}(\nu_S)$ at zero field, then, provided that the conditions (2), (3) on PE and OE are met, the increase $T_c(h) > T_c(0)$ of the transition temperature in some range of h is expected (see solid line in Fig.1). The growth of $T_c(h)$ at

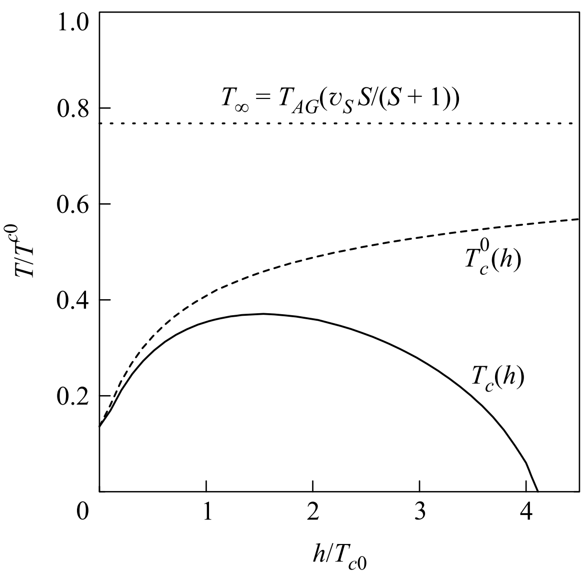

Fig.1. Enhancement of superconductivity by magnetic field. $T_c(h)$ is the transition temperature as a function of magnetic field with PE and OE taken into account (solid line). The area under $T_c(h)$ curve corresponds to superconducting state. $T_c^\circ(h)$ is the transition temperature with PE and OE disregarded (dashed line), $T_\infty = T_c^\circ(\infty)$ (dotted line). The parameters used: $\nu_S = 0.85\, T_{c0} < \nu_S^*$, $S = 1/2$, $J < 0$ (ferromagnetic exchange), $\zeta = 5$, $\nu_{\text{so}} = 10^3\, T_{c0}$, $\nu = 10^4\, T_{c0}$, $p_F d = 30$, $T_c(0) = 0.135\, T_{c0}$, $T_\infty = 0.768\, T_{c0}$

$h \lesssim T_c(h)$ is due to the reduction of the full ES rate. At higher fields PE and OE inevitably prevail, leading to complete suppression of SC at some critical field h_c. The most favorable regime for the observation of $T_c(h)$ increase is when SC is significantly suppressed at zero field, i.e. ν_S is close to (but smaller than) ν_S^*. In this case large ratio $(T_c^{\max} - T_c(0))/T_c(0)$ is expected (see Fig.1).

The most exotic situation occurs when $\nu_S > \nu_S^* > \nu_\infty = \nu_S S/(S+1)$. Then at $h = 0$ superconductivity is totaly suppressed. Disregarding PE and OE, one obtains a finite transition temperature $T_\infty = T_{AG}(\nu_\infty)$ at very high fields (indicated by dotted line in Fig.2). If the conditions (2), (3) are satisfied, superconductivity does

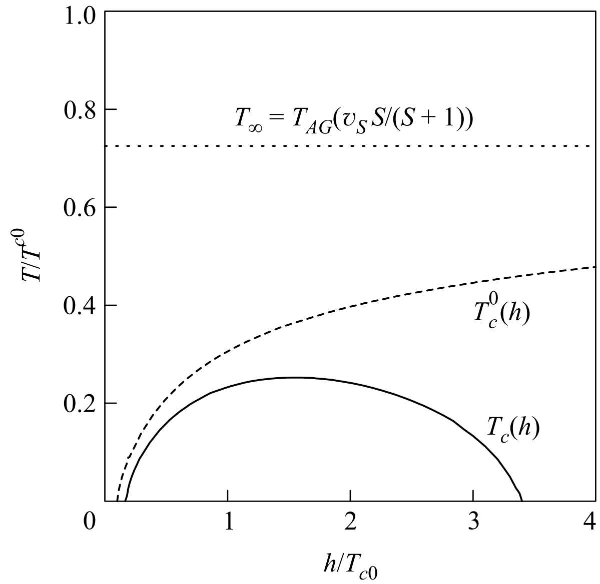

Fig.2. Magnetic-field-induced superconductivity. $T_c(h)$ is the transition temperature as a function of magnetic field with PE and OE taken into account (solid line). The area under $T_c(h)$ curve corresponds to superconducting state. $T_c^\circ(h)$ is the transition temperature with PE and OE disregarded (dashed line); $T_\infty = T_c^\circ(\infty)$ (dotted line). The parameters used: $\nu_S = 1.0\, T_{c0} > \nu_S^*$, $S = 1/2$, $J < 0$ (ferromagnetic exchange), $\zeta = 5$, $\nu_{\text{so}} = 10^3\, T_{c0}$, $\nu = 10^4\, T_{c0}$, $p_F d = 30$, $h^* = 0.17\, T_{c0}$, $T_\infty = 0.725\, T_{c0}$

not exist below some critical field h^*, but it appears at higher fields $h \geq h^*$. A nonzero transition temperature $T_c(h)$ (solid line in Fig.2) exists in a range of fields starting from h^* and terminating at some higher critical field h_c, when PE and OE dominate over ES. Such behavior is possible in the range of concentrations $n_S^* < n_S < n_S^{**}$, where n_S^{**} is smaller than $n_S^*(S+1)/S$ and is determined by the parameters involved in PE and OE. The better the conditions (2), (3) are satisfied, the closer is n_S^{**} to $n_S(S+1)/S$. The most favorable situation for the experimental observation of "magnetic-field-induced superconductivity" is realized when n_S is only slightly larger than n_S^*. In this case h^* is sufficiently small and the curve $T_c(h)$ produces a quite steep growth at the fields h just above h^* (see Fig.2). Two specific examples of $T_c(h)$ behaviour are presented in Figs.1 and 2 for $S = 1/2$, for the following set of parameters: $J < 0$ (ferromagnetic exchange), $\zeta = 5$, $\nu_{\text{so}} = 10^3\, T_{c0}$, $\nu = 10^4\, T_{c0}$, $p_F d = 30$. The similar set of parameters corresponds, for example, to the 3nm-thick PtSi film studied in [6, 7].

Below we briefly outline the method used to derive the announced results, details of our calculations will be presented in a separate publication.

The starting point of our problem is the following Hamiltonian:

$$\mathcal{H} = \mathcal{H}_{BCS} + \mathcal{H}_S + \mathcal{H}_{eS} + \mathcal{H}_{eU}.$$

Here,

$$\mathcal{H}_{BCS} = \int \left\{ \psi^\dagger_\alpha \left(\frac{1}{2m}(\mathbf{p} - e/c\mathbf{A})^2 - \varepsilon_F \right) \psi_\alpha + \right.$$
$$\left. + \frac{\lambda}{2} \psi^\dagger_\alpha \psi^\dagger_\beta \psi_\beta \psi_\alpha - \psi^\dagger_\alpha \sigma^z_{\alpha\beta} h \psi_\beta \right\} d\mathbf{r}$$

is the BCS Hamiltonian which includes the orbital and paramagnetic effects of external magnetic field on conduction electrons;

$$\mathcal{H}_{eS} = \int \left\{ \psi^\dagger_\alpha \sum_a (u_S \delta_{\alpha\beta} + J(\mathbf{S}_a, \sigma_{\alpha\beta}))\delta(\mathbf{r} - \mathbf{R}_a)\psi_\beta \right\} d\mathbf{r}$$

describes the interaction with magnetic impurities and $\mathcal{H}_S = -\sum_a \omega_S S^z_a$ is the Hamiltonian of impurity spins in external magnetic field ($\omega_S = g_S h = 2h$ is their Zeeman splitting). Finally,

$$\mathcal{H}_{eU} = \int \left\{ \psi^\dagger_\alpha(\mathbf{r}) \sum_b v_{\alpha\beta}(\mathbf{r} - \mathbf{R}_b, \mathbf{r}' - \mathbf{R}_b)\psi_\beta(\mathbf{r}') \right\} d\mathbf{r} d\mathbf{r}'$$

describes the scattering of electrons on non-magnetic impurities, which includes both potential and spin-orbit parts. Here $v_{\alpha\beta}(\mathbf{r}, \mathbf{r}')$ is the Born amplitude in coordinate representation; since we work in momentum space, we only need its Fourier tranform $v_{\alpha\beta}(\mathbf{p}, \mathbf{p}') = u_0 \delta_{\alpha\beta} + + i\, v_{\rm so}/p_F^2([\mathbf{p}, \mathbf{p}'], \sigma_{\alpha\beta})$. Magnetic and non-magnetic impurities are uniformely distributed over the sample volume with concentrations n_S and n respectively.

We solve the problem using the standard diagrammatic technique for BCS theory and disordered metals [1, 8] and employing the following approximations: i) $p_F l \gg 1$, where $l = v_F/\nu$ is the mean free path for potential scattering; ii) Born approximation for impurity scattering; iii) "dirty limit", i.e. $\nu \gg \nu_{so} \gg T_c$.

The equation for the transition temperature T can be obtained in the form

$$\ln \frac{T_{c0}}{T} = \pi T \sum_\varepsilon \left(\frac{1}{|\varepsilon|} - C_0(\varepsilon) \right), \tag{4}$$

where $C_0(\varepsilon) = 1/2(C^{\uparrow\uparrow}_{\downarrow\downarrow} - C^{\uparrow\downarrow}_{\downarrow\uparrow} + C^{\downarrow\downarrow}_{\uparrow\uparrow} - C^{\downarrow\uparrow}_{\uparrow\downarrow})$ is the singlet Cooperon component. In the approximation $p_F l \gg 1$ the Cooperon is given by an infinite sum of ladder-type diagrams, each "ladder step" containing an impurity line and the product of two disorder-averaged normal state Green functions. The expression for the components of such Green function with electron spin directed along ($\uparrow$) the external field h and in the opposite direction ($\downarrow$) reads:

$$G^{-1}_{\uparrow,\downarrow}(\varepsilon, \mathbf{p}) = i\,\varepsilon - \xi \pm h' +$$

$$+ \frac{i}{2}(\nu + \nu_{\rm so} + \Gamma\varepsilon))\mathrm{sgn}\,\varepsilon \pm i\tilde{\nu}_S \mathrm{sgn}\,\varepsilon.$$

Here, $\nu = 2\pi N_F(n_S u_S^2 + n u_0^2)$ is the potential scattering rate, $\nu_{\rm so} = 2\pi N_F n v_{\rm so}^2 2/3$ is the spin-orbit scattering rate, $\tilde{\nu}_S = 2\pi N_F n_S u_S J\langle S_z\rangle$ is the interference contribution between potential and exchange scattering on magnetic impurities (however, this term is irrelevant and falls out of the final result), $h' = h - n_S J\langle S_z\rangle$ is the effective magnetic field acting on electron spins comprised of the external field h and exchange field of polarized impurities $-n_S J\langle S_z\rangle$. Hereafter $\langle\ldots\rangle$ stands for thermodynamic average over the states of an isolated impurity spin, subjected to external magnetic field h: $\langle\hat{A}\rangle = 1/Z\sum_{m=-S}^{S} A_{mm} e^{m\omega_S/T}$, $Z = \sum_{m=-S}^{S} e^{m\omega_S/T}$. Thus,

$$\langle S_z\rangle = (S + \frac{1}{2}) \coth\left[(S + \frac{1}{2})\frac{\omega_S}{T}\right] - \frac{1}{2}\coth\frac{\omega_S}{2T}.$$

Further, $\Gamma(\varepsilon) = \nu_z + \Gamma_{\rm sf}(\varepsilon)$ is the full ES rate due to exchange interaction of electrons with polarized magnetic impurities. It is given by the sum of the rate of scattering without spin flip $\nu_z = \nu_S\langle S_z^2\rangle/S(S+1)$ and the spin flip scattering rate

$$\Gamma_{\rm sf}(\varepsilon) = \nu_S \frac{\langle S_\perp^2\rangle}{S(S+1)} - \delta\Gamma(\varepsilon), \tag{5}$$

where

$$\delta\Gamma(\varepsilon) = \nu_S \frac{\langle S_z\rangle}{S(S+1)} T \sum_{|\omega|>|\varepsilon|} \frac{2\omega_S}{\omega^2 + \omega_S^2}. \tag{6}$$

Here $\omega = 2\pi T n$ is the bosonic Matsubara frequency (n is integer) and $S_\perp^2 = S_x^2 + S_y^2$.

We now discuss properties of the full exchange scattering rate $\Gamma(\varepsilon) = \nu_z + \Gamma_{\rm sf}(\varepsilon)$ and then use the knowledge of this function while determining $T_c(h)$. For $|\varepsilon| \gg \omega_S$ at any ratio ω_S/T we have $\Gamma_{\rm sf}(\varepsilon) \approx \nu_S\langle S_\perp^2\rangle/S(S+1)$ and $\Gamma(\varepsilon) \approx \nu_S$. At zero field $\Gamma_{\rm sf}(\varepsilon) = 2/3\,\nu_S$, $\nu_z = 1/3\,\nu_S$, and $\Gamma(\varepsilon) = \nu_S$ for any ε. The full ES rate $\Gamma(\varepsilon) \approx \nu_S$ for electrons with energies $|\varepsilon| \gg \omega_S$ is not modified by magnetic field, although $\Gamma_{\rm sf}(\varepsilon)$ and ν_z do depend on h.

Consider the limit of strong polarization $\omega_S \gg T$. In this case one can replace in (6) the sum over ω by the integral and obtain

$$\Gamma_{\rm sf}(\varepsilon) = \nu_S \frac{1}{S+1}\frac{2}{\pi}\arctan\frac{|\varepsilon|}{\omega_S} \quad \text{and} \quad \nu_z = \nu_S S/(S+1). \tag{7}$$

For electron energies $|\varepsilon| \ll \omega_S$ less than Zeeman splitting $\Gamma_{\rm sf}(\varepsilon) \approx \nu_S \frac{1}{S+1}\frac{2}{\pi}\frac{|\varepsilon|}{\omega_S} \ll \nu_S$ reflecting the fact that spin flip processes freeze out for strongly polarized spins. Hence, the full ES rate $\Gamma(\varepsilon) \approx \nu_z = \nu_S S/(S+1) < \nu_S$ in a wide range of energies $|\varepsilon| \lesssim \omega_S$. At very strong field $\Gamma_{\rm sf}(\varepsilon) \to 0$ and $\Gamma(\varepsilon) = \nu_\infty = \nu_S S/(S+1)$ for all ε. Expressing $\Gamma(\varepsilon)$ in the form $\Gamma(\varepsilon) = \nu_S - \delta\Gamma(\varepsilon)$ we see that the full ES rate in nonzero field is always less than ν_S, with $\delta\Gamma(\varepsilon, \omega_S)$ for a fixed ε being a growing function of ω_S with limiting values $\delta\Gamma(\varepsilon, 0) = 0$, $\delta\Gamma(\varepsilon, \infty) = \nu_S/(S+1)$.

The Cooperon can be shown to obey the following equation for $C_0(\varepsilon)$

$$\left(|\varepsilon| + \Gamma(\varepsilon) + \frac{1}{2}(\hat{L}_0 - \Gamma_{\rm sf}(\varepsilon)) + \frac{3h'^2}{2\nu_{\rm so}} + \gamma_{\rm orb}\right) C_0 = 1. \tag{8}$$

Here

$$\gamma_{\rm orb} = \frac{1}{2} D\left(\frac{2e}{c}H\right)^2 \frac{d^2}{12} = \frac{2}{9}(p_F d)^2 \frac{h^2}{\nu}$$

is the dephasing rate corresponding to OE of magnetic field ($D = \frac{1}{3}v_F l$ is the diffusion constant) and the operator $\hat{L}_0$ acts as

$$\hat{L}_0 C_0(\varepsilon) = \nu_S \frac{\langle S_z\rangle}{S(S+1)} T \sum_\omega \frac{2\omega_S}{\omega^2 + \omega_S^2} C_0(\varepsilon - \omega).$$

At zero field $h = 0$ it is straightforward to check that $\hat{L}_0 - \Gamma_{\rm sf}(\varepsilon) = 0$ and $\Gamma(\varepsilon) = \nu_S$. Therefore the solution to (8) is $C_0(\varepsilon) = 1/(|\varepsilon| + \nu_S)$ and one recovers the AG's result (1) for transition temperature.

Enhancement of T_c by parallel field. We start our analysis from the case $\nu_S < \nu_S^*$, when a nonzero transition temperature $T_c(0) = T_{AG}(\nu_S)$ exists at zero field. First we study the equation

$$\left(|\varepsilon| + \nu_S - \delta\Gamma(\varepsilon) + \frac{1}{2}(\hat{L}_0 - \Gamma_{\rm sf}(\varepsilon))\right) C_0 = 1 \tag{9}$$

leaving in (8) the terms related to ES only and neglecting PE and OE. In the limit $h \to \infty$ we get: $\hat{L}_0 \to 0$, $\Gamma_{\rm sf}(\varepsilon) \to 0$, $\Gamma(\varepsilon) \to \nu_\infty$, and $C_0(\varepsilon) = 1/(|\varepsilon| + \nu_\infty)$. Thus in the strong-field limit and in the absense of PE and OE the transition temperature would be $T_\infty = T_{AG}(\nu_\infty)$ (indicated by dotted line in Fig.1), which is higher than the zero field value $T_{AG}(\nu_S)$ since $\nu_\infty < \nu_S$. For an arbitrary field solving Eqs. (4),(9) together numerically, one obtains the transition temperature curve $T_c^\circ(h)$ with PE and OE disregarded (dashed line in Fig.1). Formally, the enhancement of transition temperature compared to AG's zero field result $T_{AG}(\nu_S)$ is due to the term $-\delta\Gamma(\varepsilon)$ in (9) those effect is always stronger than the (opposite-sign) effect from the term operator $1/2(\hat{L}_0 - \Gamma_{\rm sf}(\varepsilon))$ in the same equation.

We are now in position to derive the conditions (2) and (3) for the strengths of paramagnetic and orbital effects compatible with observation of an increase of the actual transition temperature $T_c(h)$. Indeed, the terms in (8) related to PE and OE must be sufficiently smaller than the terms responsible for ES in the relevant fields $h \sim T_{c0}$: $[h'(h \sim T_{c0})]^2/\nu_{\rm so} \ll \nu_S$ and $\gamma_{\rm orb}(h \sim T_{c0}) \ll \nu_S$. Since we are interested in $\nu_S \sim T_{c0}$ the latter condition immediately leads to (3). Due to Born approximation ($\zeta \gg 1$) for $h \sim T_{c0}$ and $T \lesssim T_{c0}$ the exchange field $n_S J\langle S_z\rangle$ dominates over h in the effective field h' and is of the order of its maximal value $n_S JS$. Therefore, estimating $h' \sim n_S JS$, we obtain (2). Thus, provided the conditions (2) and (3) are satisfied, one observes an increase in the transition temperature $T_c(h)$ (solid line in Fig.1).

Superconductivity induced by magnetic field. Now we turn to the case $\nu_S > \nu_S^* > \nu_\infty$ or, expressed in terms of magnetic impurity concentrations, $n_S^* < n_S < < n_S^*(S+1)/S$. First we study the Eqs.(4),(9) neglecting PE and OE. Since $\nu_S > \nu_S^*$, the SC is totally suppressed at $h = 0$, but at infinite field one obtains a finite transition temperature $T_\infty = T_{AG}(\nu_\infty)$ (indicated by dotted line in Fig.2), because $\nu_\infty < \nu_S^*$. This leads to the existence of critical field $h_\circ^*$, below which SC does not exist at any temperature, but appears in greater fields $h \geq h_\circ^*$. The field $h_\circ^*$ is determined from the equation

$$\int_0^\infty d\varepsilon\, (C_0(\varepsilon, h) - 1/(\varepsilon + \nu_S^*)) = 0, \tag{10}$$

where $C_0(\varepsilon, h)$ is the solution to (9) in zero temperature limit, and depends on only one parameter ν_S. The transition temperature $T_c^\circ(h)$ in the absense of PE and OE (dashed line in Fig.2) is a growing function of h, starting from the zero value $T_c^\circ(h_\circ^*) = 0$ at $h_\circ^*$ and tending to T_∞ as $h \to \infty$. The critical field $h_\circ^*$ as a function of n_S has the following limiting values: $h_\circ^* \to 0$ as $n_S \to n_S^* + 0$, $h_\circ^* \to \infty$ as $n_S \to n_S^*(S+1)/S - 0$; and $h_\circ^* \sim T_{c0}$ when n_S is close neither to n_S^* nor to $n_S^*(S+1)/S$.

For magnetic impurity concentrations n_S not very close to $n_S^*(S+1)/S$, the field $h_\circ^* \lesssim T_{c0}$. Then, provided the conditions (2) and (3) are met, the described behavior of transition temperature in the fields $h \sim h_\circ^*$ survives under the action of orbital and paramagnetic effects. PE and OE slightly change $h_\circ^*$, making the actual critical field h^* greater than $h_\circ^*$. The actual transition temperature curve $T_c(h)$ (solid line in Fig.2) is close to $T_c^\circ(h)$ at fields $h \sim h^*$ and deviates sufficiently only at higher fields when PE and OE dominate over ES. We found the critical field $\omega_S^* = g_S h^*$ analytically (with logarithmic accuracy in ω_S^*/ν_S^*) for the case when n_S is slightly greater than critical n_S^*, i.e. $\delta\nu_S = \nu_S - \nu_S^* \ll \nu_S^*$:

$$\frac{\omega_S^*}{\nu_S^*} \ln \frac{\nu_S^*}{\omega_S^*} = \pi(S+1)\left[\frac{\delta\nu_S}{\nu_S^*} + \frac{3\,(n_S J S)^2}{2\,\nu_{\text{so}}\nu_S^*}\right]. \qquad (11)$$

If $h_\circ^* \gg T_{c0}$, i.e. n_S is close to $n_S^*(S+1)/S$, accounting for PE and OE, even with conditions (2) and (3) fulfilled, destroys SC in such a high field. Thus, for such n_S SC is totally suppressed at any field. This yields that the regime of "magnetic-field-induced SC" actually exists in a more narrow (than in the absense of PE and OE) range of concentrations $n_S^* < n_S < n_S^{**}$, where n_S^{**} is smaller than $n_S^*(S+1)/S$ and is determined by the values of parameters involved in PE and OE.

In conclusion, we have predicted the mechanism of superconductivity enchancement in thin films by external parallel magnetic field. The effect is due to the polarization of magnetic impurity spins, which reduces the full rate of electron exchange scattering. In some range of magnetic impurity concentrations the phenomenon of *magnetic-field-induced superconductivity* is predicted. The predicted effect is expected to be observable in very thin disordered superconductive films containing heavy metals leading to high spin-orbital scattering rate. We expect that similar effect may exist in superconductive-ferromagnet thin-film bilayers with spontaneous magnetization parallel to the surface.

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