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PHENOMENOLOGICAL PHASE DIAGRAM OF SUPERFLUID ^{3}He IN A STRETCHED AEROGEL

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Highly anisotropic "nematically ordered" aerogel induces global uniaxial anisotropy in superfluid ^{3}He. The anisotropy lowers symmetry of ^{3}He in the aerogel from spherical to axial. As a result, instead of one transition temperature in a state with an orbital moment $l = 1$ there are two, corresponding to projections $l_z = 0$ and $l_z = \pm 1$. This splitting has a pronounced effect on the phase diagram of superfluid ^{3}He and on the structures of the appearing phases. Possible phase diagrams, obtained phenomenologically on the basis of Landau expansion of the thermodynamic potential in the vicinity of the transition temperature are presented here. The order parameters corresponding to each phase and their temperature dependences are found.

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1. INTRODUCTION

At a triplet Cooper pairing, the transition temperature T_c is degenerate with respect to three projections of spin. In superfluid ^{3}He, where Cooper pairs are formed in a state with the orbital angular momentum $l = 1$, there is additional degeneracy with respect to three projections of the orbital angular momentum. A proper superposition of all components is represented by an order parameter, which is a 3×3 matrix of complex amplitudes $A_{\mu j}$. The spin projections are labeled here by the index μ and the orbital ones, by j. The concrete form of the order parameter is determined by minimizing the corresponding thermodynamic potential with respect to $A_{\mu j}$. In the case of superfluid ^{3}He, depending on pressure, the stable minima correspond to order parameters describing the Anderson–Brinkman–Morel (ABM) or Ballian–Werthammer (BW) phases [1]. In both cases, the form of the order parameter does not change with temperature, and only the overall amplitude Δ increases upon cooling. This is eventually a manifestation of the mentioned degeneracy.

Lowering the spherical symmetry of liquid ^{3}He by external fields or oriented impurities can split the transition to the superfluid state and partly separate elements that under cooling evolve together into the corresponding order parameter. For example, the degeneracy of T_c over spin projections is lifted by a magnetic field H_μ. Its principal effect is described by the Zeeman term in the free energy

$$\Phi_H \sim H_\mu H_\nu A_{\mu j} A^*_{\nu j},$$

which has to be added to the expansion of the free energy in powers of $A_{\mu j}$. As a result, the transition temperature T_c is split into two, such that the transition temperature for $s_z = \pm 1$ is higher than that for $s_z = 0$, and in the magnetic field, the ABM phase, which does not include the $s_z = 0$ component, is formed first.

Similarly, the degeneracy of T_c with respect to the orbital projections is lifted by global orbital anisotropy. Such anisotropy can be induced by a deformed aerogel immersed in superfluid ^{3}He [2]. Aoyama and Ikeda [3] theoretically studied the effect of uniaxial global anisotropy on the phase diagram of superfluid ^{3}He. Their argument was based on a model in which global anisotropy is induced by the averaged effect of anisotropic scattering of quasiparticles by oriented impurities. They predicted, in particular, that a uniaxial stretch of aerogel just below the transition temperature would stabilize the polar phase, which on cooling to lower temperatures undergoes a continuous transition to the distorted ABM phase and eventually the BW phase is formed via a first-order transition. These predictions were tested in experiments with the "nematically ordered" aerogel [4], which can be regarded

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as being infinitely stretched. The phase diagram found experimentally confirms the predicted sequence of the phase transitions, but other, even qualitative features of the two phase diagrams are different.

In this paper, possible phase diagrams of superfluid ^{3}He in a stretched aerogel are considered phenomenologically. It is shown that depending on the values of phenomenological parameters characterizing this system, different routes of development of the order parameter upon cooling from the transition temperature are possible. The orbital anisotropy is formally described by an additional term in the thermodynamic potential

$$\Phi_\kappa \sim \kappa_{jl} A_{\mu j} A^*_{\mu l},$$

where κ_{jl} is a real symmetric tensor, which can be taken traceless. It is assumed to be uniform (i. e., independent of the coordinate). Random local anisotropy is neglected. This approximation is well justified in the present context, when only structures of order parameters of possible phases are considered. On the other hand, random anisotropy can strongly affect the orientation of order parameters of the distorted ABM and of the axi-planar phases, giving rise to a randomly nonuniform Larkin–Imry–Ma (LIM) state. In the case of a stretched aerogel, it is a two-dimensional LIM state, as discussed in Refs. [5, 6]. In comparing with experiment, in particular, with NMR data, a corresponding averaging over orientations of the order parameter has to be made.

With the global anisotropy taken into account, the standard expansion of the thermodynamic potential in powers of $A_{\mu j}$ is

$$\begin{aligned}\Phi_s = \Phi_n + N_{eff}[(\tau\delta_{jl} + \kappa_{jl})A_{\mu j}A^*_{\mu l} + \\ + \frac{1}{2}(\beta_1 A_{\mu j}A_{\mu j}A^*_{\nu l}A^*_{\nu l} + \beta_2 A_{\mu j}A^*_{\mu j}A_{\nu l}A^*_{\nu l} + \\ + \beta_3 A_{\mu j}A_{\nu j}A^*_{\mu l}A^*_{\nu l} + \beta_4 A_{\mu j}A^*_{\nu j}A_{\nu l}A^*_{\mu l} + \\ + \beta_5 A_{\mu j}A^*_{\nu j}A_{\mu l}A^*_{\nu l})]. \quad (1)\end{aligned}$$

Here,

$$\tau = \frac{T - T_c}{T_c}$$

is the dimensionless temperature, T_c is the transition temperature, defined such that it includes all global isotropic shifts from the corresponding temperature of bulk ^{3}He. The overall coefficient N_{eff} has the dimension of density of states. Phenomenological coefficients $\beta_1, \ldots, \beta_5$ depend on the pressure and the properties of the aerogel. When anisotropy is uniaxial, in proper axes

$$\kappa_{xx} = \kappa_{yy} = -\kappa, \quad \kappa_{zz} = 2\kappa.$$

In contrast to a magnetic field, which always favors $s_z = \pm 1$ projections, a uniaxial deformation of the aerogel, depending on the sign of κ, favors either the $l_z = \pm 1$ or $l_z = 0$ projection. For a compressed aerogel, $\kappa > 0$ and the states with $l_z = \pm 1$ have a higher transition temperature, while for a stretched aerogel, $\kappa < 0$, a state with $l_z = 0$ is favored.

Stabilization of the polar phase by a stretched aerogel within this approach follows immediately from the explicit form of the second-order terms in the expression for thermodynamic potential (1):

$$(\tau + 2\kappa)A_{\mu z}A^*_{\mu z} + (\tau - \kappa)(A_{\mu x}A^*_{\mu x} + A_{\mu y}A^*_{\mu y}).$$

For negative κ, the highest transition temperature is $\tau = -2\kappa$. For realistic values of the coefficients β, in particular, if $\beta_{15} < 0$ (here and in what follows, conventional shorthand notation for sums of β coefficients is used, e. g., $\beta_1 + \beta_5 = \beta_{15}$, etc.) below $\tau = -2\kappa$ the superfluid polar phase is favored [7]. Its order parameter can be written as

$$A^0_{\mu j} = \Delta_0 \exp(i\varphi)\, d_\mu m_j,$$

where d_μ is a real spin vector and m_j is a unit vector in z-direction.

The polar phase is stable within the interval of temperatures $\tau \sim \kappa$. On further cooling, the suppressed angular momentum projections $l_z = \pm 1$ come into effect; they change the order parameter symmetry and further phase transitions can occur. While stabilization of the polar phase practically depends only on the sign of κ, its stability interval and the sequence of further transitions also depend on the values of the coefficients $\beta_1, \ldots, \beta_5$. To avoid the discussion of nonrealistic situations, we have to restrict the region of admitted values of βs. Within the BCS theory, their values are proportional to one combination of the parameters, $\beta_0 = 7\zeta(3)/8\pi^2 T_c^2$:

$$\beta_1, \ldots, \beta_5 = \beta_0(-1/2, 1, 1, 1, -1).$$

This set of values of the βs is referred to as the weak coupling limit [1]. In the definition of β_0 $\zeta(3)$ is the Riemann zeta function. The observed thermodynamic properties of bulk superfluid ^{3}He in the vicinity of T_c can be fitted by the $\beta_1, \ldots, \beta_5$ that deviate from their weak coupling values by 10–20 % [8]. The deviations are smaller at low pressures. For ^{3}He in an aerogel, the situation is less certain. Impurities give rise to corrections to the β coefficients of the order of ξ_0/λ, where ξ_0 is the correlation length of superfluid ^{3}He and λ is the mean free path. This ratio is of the order of 1/10. In what follows, we assume that deviations of the βs

for superfluid ^{3}He in a nematically ordered aerogel from their-weak coupling values are also of the order of 1/10, at least at low pressures.

There is another reason for restricting the present discussion to a region of low pressures (for example, below 10 bar). The diameters of strands in a nematically ordered aerogel, estimated as $d \sim 10$ nm [4], are bigger than in silica aerogels and can be comparable with the correlation length of superfluid ^{3}He, which at pressures above 20 bar is about 20 nm. When $d \sim \xi_0$, perturbation of the order parameter in a vicinity of a strand is of the order of unity. Well below T_c, the condensate varies over a distance $\sim \xi_0$, which is smaller than the average distance $\xi_a \approx 200$ nm. In this situation, the condensate is essentially nonuniform and the average order parameter does not properly characterize the state of ^{3}He. The uniform approximation works better at low pressures and in the vicinity of T_c in the region where the Ginzburg and Landau (GL) coherence length $\xi(T)$ exceeds not only the diameter of a strand but also ξ_a. In this GL region $d \ll \xi_a \ll \xi(T)$, the average order parameter $A_{\mu j}$ is a suitable characteristic of the state of superfluid ^{3}He.

Preliminary results of phenomenological analysis of the phase diagram of superfluid ^{3}He in a nematically ordered aerogel were published before [7]. A principal suggestion in that paper was to regard the extra line (ESP2) in the experimentally found phase diagram as evidence of the possible stability (or meta-stability) of the axi-planar phase. Further experiments and their analysis [9] have shown that this suggestion is not correct. Nevertheless, there remains the question of possible stability of the axi-planar phase in an anisotropic environment. It was shown in [7] that the axi-planar phase is a possible minimum of the thermodynamic potential at weak-coupling values of the phenomenological coefficients β. The weak-coupling limit corresponds to a singular point in the space of parameters β and the solution presented in Ref. [7] is only one of many possibilities. In what follows, the question of stability of the axi-planar phase is discussed with the account of the mentioned singularity, and the conditions determining the possibility of its existence are found.

2. FURTHER PHASES

To find further possible phase transitions, we represent the order parameter as

$$A_{\mu j} = A^0_{\mu j} + a_{\mu j},$$

where

$$A^0_{\mu j} = \Delta_0 \exp(i\varphi)\, d_\mu m_j$$

is the order parameter of the polar phase and $a_{\mu j}$ is a small increment, and expand the change of the thermodynamic potential

$$\bar{\Phi} \equiv \frac{\Phi_s - \Phi_n}{N_{eff}}$$

in powers of $a_{\mu j}$, separating terms of different orders:

$$\bar{\Phi} = \bar{\Phi}_0 + \bar{\Phi}_2 + \bar{\Phi}_4. \tag{2}$$

The zeroth-order term in $a_{\mu j}$,

$$\bar{\Phi}_0 = \Delta_0^2(\tau + 2\kappa) + \frac{1}{2}\beta_{12345}\Delta_0^4, \tag{3}$$

represents the energy gain of the polar phase with respect to the normal phase and determines the temperature dependence of the amplitude Δ_0:

$$\Delta_0^2 = -\frac{\tau + 2\kappa}{\beta_{12345}}.$$

The explicit form of the second-order term depends on the choice of gauge for $A^0_{\mu j}$. It is convenient to take $\varphi = 0$, such that $A^0_{\mu j}$ is a real matrix. The expected transition is due to the occurrence of two of angular momentum projections $l_z = \pm 1$ previously suppressed by the anisotropy. This means that only the components of $a_{\mu j}$ transverse to m_j are essential, and the condition $a_{\mu j} m_j = 0$ has to be imposed. With this simplification, we have

$$\begin{aligned}\bar{\Phi}_2 &= (\tau - \kappa) a_{\mu j} a^*_{\mu j} + \\ &+ \frac{1}{2}\Delta_0^2\{\beta_1(a_{\mu j}a_{\mu j} + a^*_{\mu j}a^*_{\mu j}) + 2\beta_2 a_{\mu j}a^*_{\mu j} + \\ &+ \beta_3 d_\mu d_\nu (a_{\mu j}a_{\nu j} + a^*_{\mu j}a^*_{\nu j}) + 2\beta_{45} d_\mu d_\nu a_{\mu j} a^*_{\nu j}\}. \end{aligned} \tag{4}$$

All experimentally observed transitions occurs at $|\tau| \sim \sim 0.1$, where Eq. (1) is still a good approximation for $\bar{\Phi}$. For this reason, corrections of the order of τ, originating from the higher-order terms in the expansion of $\bar{\Phi}$ over $A_{\mu j}$, are neglected and terms of the fourth order in $a_{\mu j}$ have the same form as the analogous terms in Eq. (1) with the substitution of $a_{\mu j}$ instead of $A_{\mu j}$:

$$\bar{\Phi}_4 = \frac{1}{2}\sum_s \beta_s I_s(a_{\mu j}, a^*_{\nu l}). \tag{5}$$

The polar phase preserves its stability (or meta-stability) until $\bar{\Phi}_2$ is positive definite with respect to $a_{\mu j}$. The coefficients in $\bar{\Phi}_2$ are determined by the more symmetric (polar) phase. In particular, d_μ and m_j

are the respective symmetry axes in spin and orbital spaces, and therefore the components of $a_{\mu j}$ parallel and perpendicular to d_μ are essentially different. Of possible orientations in the orbital space only the components perpendicular to m_j are essential, as was explained above. Because of the fixed gauge of the polar phase, $\bar{\Phi}_2$ is not gauge invariant with respect to $a_{\mu j}$. On the other hand $\bar{\Phi}_2$ is T-invariant and the transition to T-even,

$$2b_{\mu j} = a_{\mu j} + a^*_{\mu j},$$

and T-odd,

$$2ic_{\mu j} = a_{\mu j} - a^*_{\mu j},$$

combinations of $a_{\mu j}$ diagonalizes $\bar{\Phi}_2$:

$$\bar{\Phi}_2 = \Lambda_1(\tau,\kappa)(\delta_{\mu\nu} - d_\mu d_\nu)b_{\mu j}b_{\nu j} + [\Lambda_1(\tau,\kappa) + (\beta_{45} - 2\beta_1 - \beta_3)\Delta_0^2]d_\mu d_\nu c_{\mu j}c_{\nu j} + \Lambda_2(\tau,\kappa) \times d_\mu d_\nu b_{\mu j}b_{\nu j} + [\Lambda_2(\tau,\kappa) - (\beta_{45} + 2\beta_1 + \beta_3)\Delta_0^2] \times (\delta_{\mu\nu} - d_\mu d_\nu)c_{\mu j}c_{\nu j}. \quad (6)$$

Here,

$$\Lambda_1 = \tau - \kappa + \beta_{12}\Delta_0^2$$

and

$$\Lambda_2 = \tau - \kappa + \beta_{12345}\Delta_0^2.$$

Hence, for each j, there are four different variables: real and imaginary parts of the components of $a_{\mu j}$ parallel and orthogonal to d_μ. In principle, there may be four different continuous transitions from the polar phase into a less symmetric phase. Each of the transitions occurs when the coefficient in front of the corresponding second-order term changes sign, e. g., the component of the real part of $a_{\mu j}$ that is perpendicular to d_μ can occur at τ determined by the condition

$$\Lambda_1(\tau,\kappa) = 0.$$

Most important is the transition with the highest of the four values of τ.

When the coefficients β_s have their weak-coupling values, the combinations

$$\beta_{45} \equiv \varepsilon = 0$$

and

$$2\beta_1 + \beta_3 \equiv \nu = 0.$$

Instead of four different transition temperatures, there are two in this limit, and they are determined by the conditions

$$\Lambda_1(\tau_1,\kappa) = 0$$

and

$$\Lambda_2(\tau_2,\kappa) = 0,$$

both being doubly degenerate. The degeneracy is a manifestation of the "hidden symmetry", which exists in the weak-coupling limit for the equal spin pairing (ESP) states. For such states, the quantization axis of spin can be chosen such that the condensate of Cooper pairs contains only pairs with spin projections ± 1. The BCS Hamiltonian in that case does not couple condensates with different spin projections and these two condensates can be treated as independent. In particular, they can have different orientations of the orbital parts of the order parameter and different complex phases. Strong-coupling corrections give rise to a coupling between the two condensates and lift the corresponding degeneracy. The "hidden symmetry" was previously discussed in the context of classifying collective modes in the ABM phase [1, 10].

The condition $\Lambda_2(\tau,\kappa) = 0$ leads to $\kappa = 0$. At a finite κ, there is no transition, resulting in occurrence of $b_{\mu j}$ parallel to $d_{\mu j}$ and $c_{\mu j}$ perpendicular to $d_{\mu j}$. Another condition $\Lambda_1(\tau,\kappa) = 0$ has the solution

$$\tau \equiv \tau_B = \kappa(1 + 3\beta_{12}/\beta_{345}).$$

In the weak-coupling limit below this τ, the incremental order parameter is a linear combination of $b_{\mu j}$ perpendicular to $d_{\mu j}$ and $c_{\mu j}$ parallel to $d_{\mu j}$. When the strong-coupling corrections are restored, they lift this degeneracy such that τ_B relates only to the perpendicular $b_{\mu j}$, while the parallel components $c_{\mu j}$ can occur below

$$\tau_A = \frac{\kappa(3\beta_{245} - \beta_{13})}{2\beta_{13}}.$$

The difference between the two transition temperatures can be expressed in terms of the parameters ε and ν introduced above:

$$\tau_A - \tau_B = \frac{3\kappa\beta_{12345}}{2\beta_{13}\beta_{345}}(\varepsilon - \nu).$$

Depending on the sign of the difference $\varepsilon - \nu$, one or another type of the order parameter is favored. The strong-coupling corrections "transfer" the orbital anisotropy in the spin space. If d_μ is taken as a quantization axis, the components parallel to d_μ correspond to $s_z = 0$ and the perpendicular ones, to $s_z = \pm 1$. At $\varepsilon > \nu$, the favored incremental order parameter is a combination of projections $s_z = \pm 1$, and hence the full order parameter is

$$A^B_{\mu j} = \Delta_0 d_\mu m_j + \Delta_2 e_\mu l_j + \Delta_3 f_\mu n_j, \quad (7)$$

where n_j and l_j are two mutually orthogonal vectors that together with m_j form a basis in the orbital space, and Δ_2 and Δ_3 are real amplitudes. Minimizing $\bar{\Phi}$ with respect to the amplitudes Δ_0, Δ_2, and Δ_3 yields

$$\Delta_2^2 = \Delta_3^2 = -\frac{\tau - \tau_B}{3\beta_{12} + \beta_{345}}, \quad \Delta_0^2 = -\frac{3\kappa}{\beta_{345}} + \Delta_2^2$$

with the energy gain

$$\Phi_B - \Phi_p = -\frac{\beta_{345}}{\beta_{12345}(3\beta_{12} + \beta_{345})}(\tau - \tau_B)^2.$$

This is the order parameter of the distorted BW phase; it is specified by two amplitudes Δ_0 and Δ_2 with different temperature dependences [3, 7].

For conditions of the experiments in [4] on cooling from the polar phase, the distorted ABM phase occurs first. This means that the opposite inequality $\varepsilon < \nu$ is satisfied. In this case, the increment $c_{\mu j}$ is parallel to d_μ; it corresponds to $s_z = 0$, and the resulting order parameter is

$$A^A_{\mu j} = \Delta_0 d_\mu m_j + i\Delta_1 d_\mu n_j. \tag{8}$$

Here, Δ_1 is a real amplitude. The temperature dependences of Δ_0 and Δ_1 are found by minimizing $\bar{\Phi}$:

$$\Delta_1^2 = -\frac{\tau - \tau_A}{2\beta_{245}}, \quad \Delta_0^2 = -\frac{3\kappa}{2\beta_{13}} - \frac{\tau - \tau_A}{2\beta_{245}}.$$

In comparison with the polar phase, the new phase has a lower thermodynamic potential. The gain is

$$\Phi_A - \Phi_p = -\frac{\beta_{13}}{2\beta_{245}\beta_{12345}}(\tau - \tau_A)^2,$$

see [7]. The distorted BW phase in Eq. (8) in conditions of the experiments in [4] is reached via a first-order transition. The temperature τ_B preserves its meaning of the upper limit temperature for existence of this phase and determines temperature dependences of the amplitudes Δ_2^2 and Δ_3^2.

3. AXI-PLANAR PHASE

In the distorted ABM phase, d_μ is still a symmetry axis in spin space. Further lowering this symmetry via a continuous phase transition is possible at cooling when the suppressed perpendicular projection $b_{\mu j}$ comes into effect. Analysis of the stability of $A^A_{\mu j}$ with respect to $b_{\mu j}$ along the lines of the preceding section with the order parameter of the form

$$A_{\mu j} = A^A_{\mu j} + b_{\mu j}$$

yields

$$\bar{\Phi}_2 = [\tau - \kappa + \beta_{12}\Delta_0^2 + (\beta_{234} - \beta_{15})\Delta_1^2](b_{\mu j}n_j)(b_{\mu i}n_i) + \\ + [\tau - \kappa + \beta_{12}\Delta_0^2 + (\beta_1 - \beta_2)\Delta_1^2](b_{\mu j}l_j)(b_{\mu i}l_i) \tag{9}$$

instead of Eq. (6). The combination in front of $(b_{\mu j}n_j)(b_{\mu i}n_i)$ is positive, but that in front of $(b_{\mu j}l_j)(b_{\mu i}l_i)$ changes sign at $\tau = \tau_{AP}$, where

$$\tau_{AP} = -\frac{\kappa}{2} + \frac{3\kappa}{2}\frac{\beta_{245}}{\beta_{13}}\frac{\nu}{\varepsilon}. \tag{10}$$

It is the temperature of a continuous transition in the phase with finite $(b_{\mu j}l_j)(b_{\mu i}l_i)$. Its dependence on the parameters ν and ε in the limit as $\nu \to 0$ and $\varepsilon \to 0$ is singular. It follows from

$$\tau_{AP} - \tau_A = \frac{3\kappa}{2}\frac{\beta_{245}}{\beta_{13}}\left(\frac{\nu}{\varepsilon} - 1\right)$$

that $\tau_{AP} < \tau_A$ if $\nu > \varepsilon$. The transition temperature τ_{AP} is not far from τ_A if $\varepsilon \approx \nu$. At $\varepsilon \approx 3\kappa\nu$, τ_{AP} moves to lower temperatures, well beyond the limit of applicability of expansion (1). Below τ_{AP}, the order parameter is that of the axi-planar phase:

$$A_{\mu j} = \Delta_0 d_\mu m_j + i\Delta_1 d_\mu n_j + \Delta_2 e_\mu l_j. \tag{11}$$

Minimizing the thermodynamic potential with respect to Δ_0, Δ_1, Δ_2 gives the following equations for the amplitudes:

$$[\beta_{12345}\Delta_0^2 + (\beta_{245} - \beta_{13})\Delta_1^2 + \\ + \beta_{12}\Delta_2^2 + (\tau + 2\kappa)]\Delta_0 = 0, \tag{12}$$

$$[(\beta_{245} - \beta_{13})\Delta_0^2 + \beta_{12345}\Delta_1^2 + \\ + (\beta_2 - \beta_1)\Delta_2^2 + (\tau - \kappa)]\Delta_1 = 0, \tag{13}$$

$$[\beta_{12}\Delta_0^2 + (\beta_2 - \beta_1)\Delta_1^2 + \beta_{12345}\Delta_2^2 + (\tau - \kappa)]\Delta_2 = 0. \tag{14}$$

In the case $\nu > \varepsilon$, solutions of these equations reproduce the sequence of phase transitions under cooling from $\tau = -2\kappa$. At $\tau_A < \tau < -2\kappa$, a stable solution is $\Delta_1 = 0$, $\Delta_2 = 0$, and from Eq. (12),

$$\Delta_0^2 = -\frac{\tau + 2\kappa}{\beta_{12345}}.$$

At $\tau = \tau_A$, Δ_1 starts to increase, indicating a second-order phase transition in the distorted ABM phase, as discussed at the end of the preceding section. If τ_{AP} is within the applicability limits of the Ginzburg–Landau expansion, the distorted ABM phase remains stable in the interval $\tau_{AP} < \tau < \tau_A$. Below τ_{AP}, Δ_2 becomes finite and Eqs. (12)–(14) admit a solution corresponding to the axi-planar phase:

$$\Delta_2^2 = \frac{\varepsilon\beta_{13}(\tau_{AP} - \tau)}{(\nu + \varepsilon)\beta_2\beta_3 + \varepsilon\nu\beta_{23} + \varepsilon^2\beta_{13}}, \tag{15}$$

$$\Delta_1^2 = -\frac{3\kappa}{4\beta_{13}}\,\frac{\nu-\varepsilon}{\varepsilon} + \frac{(\beta_3+\varepsilon)\nu}{2\beta_{13}\varepsilon}\Delta_2^2, \qquad (16)$$

$$\Delta_0^2 = \frac{\nu+\varepsilon}{2\beta_{13}\varepsilon}\left(-\frac{3\kappa}{2} + \beta_3\Delta_2^2\right). \qquad (17)$$

It follows from Eq. (15) that the axi-planar phase can exist only if $\varepsilon > 0$. Experimentally, it can be detected by CW NMR method. As was discussed in Ref. [7], when the magnetic field is perpendicular to the anisotropy axis, the transverse NMR shift is zero for the distorted ABM phase, but it is finite and proportional to Δ_2^2 in the axi-planar phase. The transition could also be detected by a jump in the specific heat at cooling of the distorted ABM phase.

Both Δ_2^2 and τ_{AP} are very sensitive to the values of the parameters ν and ε, which for superfluid ^{3}He in a nematically ordered aerogel are poorly known. Currently, it is difficult to even estimate their values. In the published NMR data [4], the distorted ABM phase remains meta-stable until $\approx 0.7T_c$, when it jumps to the low-temperature phase, which is identified as a distorted BW phase. There is no indication of a continuous transition of the distorted ABM in the axi-planar phase. A more detailed theoretical discussion of the properties of the axi-planar phase could be appropriate if such indications would be available.

4. DISCUSSION

A nematically ordered aerogel turned out to be an efficient tool for "decomposing" the order parameter of superfluid ^{3}He into its constituents. Phenomenological analysis shows that, in principle, there are more possible phase diagrams than observed in real ^{3}He and found in the model calculation in [3]. The first superfluid phase realized immediately below T_c is always the polar phase. Its symmetry is higher than that of the most stable phases at low temperature, ABM and BW. This means that these phases can be reached in steps, via one or more phase transitions. The concrete route depends on particular values of the phenomenological parameters β. Of importance are two combinations of these parameters,

$$\varepsilon = \beta_4 + \beta_5$$

and

$$\nu = 2\beta_1 + \beta_3.$$

If $\varepsilon > \nu$, the distorted BW phase can form at cooling via a second-order transition directly from the polar phase. Such a scenario is admitted by symmetry, but not realized either in experiment [4] or in microscopic calculations [3]. The situation where $\varepsilon < \nu$ corresponds to the observed sequence of phase transitions on cooling: a distorted ABM phase forms from the polar phase via a continuous transition. If $\varepsilon > 0$ and $\varepsilon \approx \nu$, a further continuous transition, to the axi-planar phase, is possible. This transition is also admitted by symmetry, but is not observed. Unfortunately, there are no efficient tools for tuning the parameters β, and therefore not all possible phase diagrams can be realized in real ^{3}He, but the phenomenological description can be used as a framework for a systematic description of experimental data for the realized scenario in the vicinity of the transition of ^{3}He into the superfluid state.

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