

Anisotropic 2D Larkin–Imry–Ma state in polar distorted ABM phase of ^3He in “nematically ordered” aerogel

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Submitted 21 October 2014

We present results of experiments in superfluid phases of ^3He confined in aerogel which strands are nearly parallel to one another. High temperature superfluid phases of ^3He in this aerogel (ESP1 and ESP2) are biaxial chiral phases and have polar distorted ABM order parameter which orbital part forms 2D Larkin–Imry–Ma state. We demonstrate that this state can be anisotropic if the aerogel is squeezed in direction transverse to the strands. Values of this anisotropy in ESP1 and ESP2 phases are different, what leads to different NMR properties.

DOI: 10.7868/S0370274X14220135

1. Introduction. A so-called “nematically ordered” (N-) aerogel differs from standard silica aerogels by a high value of a global anisotropy. This aerogel consists of $\text{Al}_2\text{O}_3\cdot\text{H}_2\text{O}$ strands which are nearly parallel to one another [1], i.e. it may be considered as aerogel with infinite stretching anisotropy. Investigations of superfluid ^3He confined in N-aerogel are especially interesting because according to a theory [2] such a strong anisotropy may make a superfluid polar phase more favorable than Anderson–Brinkman–Morel (ABM) phase which corresponds to A phase of bulk ^3He and to A-like phase of ^3He in isotropic or weakly anisotropic silica aerogels [3–5]. A superfluid phase diagram of ^3He in N-aerogel is different from the case of ^3He in silica aerogel with similar porosity [6]. The superfluid transition temperature (T_{ca}) is slightly (by 3–6%) suppressed in comparison with the transition temperature (T_c) of bulk ^3He . Depending on prehistory, pressure and temperature, three superfluid phases are observed: two Equal Spin Pairing phases (ESP1 or ESP2) and Low Temperature phase (LTP). The ESP1 phase appears on cooling from the normal state. On further cooling the first order transition into the LTP takes place. Due to inhomogeneities of the aerogel, this transition occurs in a wide temperature range ($\sim 0.05 T_c$). On warming from the LTP the back transition into the ESP phase is observed. At high pressures ($P \geq 10$ bar) the NMR frequency shift in this phase, called ESP2 phase, is greater than in the ESP1 phase at the same conditions.

The LTP has a polar distorted Balian–Werthamer (BW) order parameter [7]. As for ESP phases, their NMR properties point out that they both have ABM order parameter with a strong polar distortion [6]. This distortion is larger at low pressures and at higher temperatures. It was also found that the order parameter orbital vector \mathbf{l} of the distorted ABM phase in N-aerogel is in a spatially inhomogeneous Larkin–Imry–Ma (LIM) state similar to that predicted in [8] and observed in A-like phase of ^3He in silica aerogel [4, 9]. In N-aerogel we get the two-dimensional LIM state because the aligned strands orient \mathbf{l} normal to their axis.

In this paper we present results of nuclear magnetic resonance (NMR) studies of liquid ^3He confined in N-aerogel which was slightly squeezed in direction transverse to the strands. In particular, these experiments allow us to explain the difference between properties of ESP1 and ESP2 phases.

2. Theory. Transverse NMR frequency shift can be found from the following equation [10]:

$$\Delta\omega = -\frac{g}{\chi H} \frac{\partial \bar{U}_D}{\partial \cos \beta}, \quad (1)$$

where g is the gyromagnetic ratio, χ – the spin susceptibility, H – the external magnetic field, β – the tipping angle of the magnetization, and \bar{U}_D – the density of the dipole energy, averaged over a fast spin precession. For the LIM state the dipole energy should also be averaged over the space (see e.g. [4, 11]). The order parameter of the ABM phase with polar distortion is:

$$A_{jk} = \Delta_0 e^{i\phi} d_j (am_k + ibn_k), \quad (2)$$

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where Δ_0 is the gap parameter, \mathbf{d} is the unit spin vector, \mathbf{m} and \mathbf{n} are mutually orthogonal unit vectors in the orbital space, and $a^2 + b^2 = 1$. For the ABM phase $a = b$, for polar distorted ABM phase $a^2 > b^2$, and for polar phase $a = 1, b = 0$. Similarly to pure ABM phase, the distorted ABM phase is a chiral phase and we can introduce the orbital vector $\mathbf{l} = \mathbf{m} \times \mathbf{n}$ which orientation defines two Weyl points in the momentum space: the energy gap of this phase equals 0 along \mathbf{l} and equals $\sqrt{2}a\Delta_0$ and $\sqrt{2}b\Delta_0$ along \mathbf{m} and \mathbf{n} . Note that the polar phase is not chiral and its gap has line of zeroes in the plane normal to \mathbf{m} . The dipole energy density for the order parameter (2) is:

$$U_D = \frac{6}{5}g_D [a^2(\mathbf{d}\mathbf{m})^2 + b^2(\mathbf{d}\mathbf{n})^2], \quad (3)$$

where $g_D = g_D(T)$ is the dipole constant. In weak coupling limit g_D can be expressed in terms of the Leggett frequency of the pure ABM phase Ω_A [12]:

$$g_D = \frac{2}{3 - 4a^2b^2} g_D^A = \frac{2}{3 - 4a^2b^2} \left(\frac{5}{6} \frac{\chi}{g^2} \Omega_A^2 \right), \quad (4)$$

where g_D^A is the dipole constant of the ABM phase. Strong coupling corrections to (4) do not exceed $\pm 5\%$ [13], therefore we do not consider them below.

Following [11, 14], we use two coordinate frames: an orbital frame $(\hat{\xi}, \hat{\eta}, \hat{\zeta})$ bound to the aerogel sample and a spin frame $(\hat{x}, \hat{y}, \hat{z})$. We choose $\mathbf{H} = H\hat{z}$ and fix $\hat{\zeta}$ -axis along aerogel strands. Then strands of N-aerogel orient $\mathbf{m} \parallel \hat{\zeta}$ and $\mathbf{l} \perp \hat{\zeta}$ [2]. In the isotropic 2D LIM state vectors \mathbf{l} and \mathbf{n} are randomly distributed in $\hat{\xi} - \hat{\eta}$ plane and $\langle l_\xi^2 \rangle = \langle l_\eta^2 \rangle = \langle n_\xi^2 \rangle = \langle n_\eta^2 \rangle = 1/2$, where angle brackets mean the space averaging. We introduce the angle $\lambda = \lambda(\mathbf{r})$ which defines the orientation of specific \mathbf{l} and the corresponding \mathbf{n} : $l_\xi = -n_\eta = \cos \lambda$ and $l_\eta = n_\xi = \sin \lambda$. For uniaxially anisotropic in $\hat{\xi} - \hat{\eta}$ plane 2D LIM state we fix the $\hat{\zeta}$ -axis along the direction corresponding to the maximum value of $\langle l_\xi^2 \rangle$. Consequently $1 > \langle l_\xi^2 \rangle > 1/2 > \langle l_\eta^2 \rangle$ and we assume that the distribution of $l_\xi = l_\xi(\lambda)$ is symmetric.

An orientation of \mathbf{H} with respect to the aerogel (Fig. 1) is described by angles of rotation of the orbital frame: μ (rotation around $\hat{\xi}$) and φ (rotation around $\hat{\zeta}$). Then we get:

$$\begin{aligned} m_x &= 0, & m_y &= -\sin \mu, & m_z &= \cos \mu, \\ n_x &= \sin(\varphi + \lambda), & n_y &= -\cos \mu \cos(\varphi + \lambda), \\ n_z &= -\sin \mu \cos(\varphi + \lambda). \end{aligned} \quad (5)$$

Motions of \mathbf{d} in the spin frame are described by Euler angles (α, β, γ) , where α corresponds to the phase

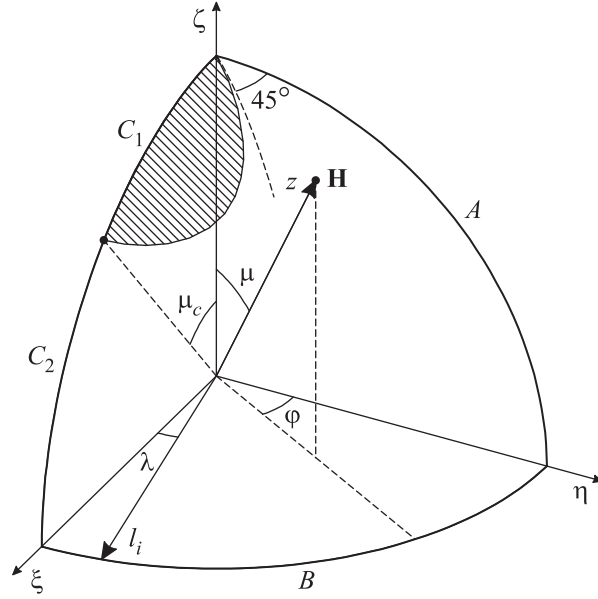


Fig. 1. Orientation of \mathbf{H} with respect to N-aerogel axes

of spin precession and β is the tipping angle. After an averaging over the fast spin precession we obtain:

$$\begin{aligned} \bar{d}_x^2 &= \frac{1}{4} \langle \cos^2 \Phi \rangle (1 + \cos \beta)^2 + \frac{1}{8} (1 - \cos \beta)^2, \\ \bar{d}_y^2 &= \frac{1}{4} \langle \sin^2 \Phi \rangle (1 + \cos \beta)^2 + \frac{1}{8} (1 - \cos \beta)^2, \\ \bar{d}_z^2 &= \frac{1}{2} \sin^2 \beta, & \overline{d_x d_z} &= \overline{d_y d_z} = 0, \\ \overline{d_x d_y} &= -\frac{1}{8} \langle \sin 2\Phi \rangle (1 + \cos \beta)^2, \end{aligned} \quad (6)$$

where $\Phi = \alpha + \gamma$ is a slow variable. Then the dipole energy (3) averaged over the space is:

$$\begin{aligned} \bar{U}_D &= \frac{6}{5}g_D \left[a^2(\bar{d}_y^2 m_y^2 + \bar{d}_z^2 m_z^2) + \right. \\ &\left. + b^2(\bar{d}_x^2 \langle n_x^2 \rangle + \bar{d}_y^2 \langle n_y^2 \rangle + \bar{d}_z^2 \langle n_z^2 \rangle + 2\overline{d_x d_y} \langle n_x n_y \rangle) \right], \quad (7) \end{aligned}$$

where $\langle n_x^2 \rangle = \sin^2 \varphi \langle \cos^2 \lambda \rangle + \cos^2 \varphi \langle \sin^2 \lambda \rangle$, $\langle n_y^2 \rangle = \cos^2 \mu (\cos^2 \varphi \langle \cos^2 \lambda \rangle + \sin^2 \varphi \langle \sin^2 \lambda \rangle)$, $\langle n_z^2 \rangle = \sin^2 \mu (\cos^2 \varphi \langle \cos^2 \lambda \rangle + \sin^2 \varphi \langle \sin^2 \lambda \rangle)$, and $\langle n_x n_y \rangle = (2 \langle \sin^2 \lambda \rangle - 1) \cos \mu \sin \varphi \cos \varphi$. The angle Φ may be spatially homogeneous (the spin nematic state, SN) or random (the spin glass state, SG) [4]. The SN state is more favorable and corresponds to the homogeneous spatial distribution of \mathbf{d} , but the SG state may be created e.g. in pulse NMR experiments after an application of large tipping pulses. In the isotropic SG state $\langle \sin^2 \Phi \rangle = \langle \cos^2 \Phi \rangle = 1/2$ and $\langle \sin 2\Phi \rangle = 0$ while in the SN state Φ is determined by minimization of (7).

The result of the minimization is shown in Fig. 1 where the shaded area corresponds to orientations of \mathbf{H} with $\sin^2 \Phi = 1$ while for other orientations the minimum of (7) corresponds to $\sin^2 \Phi = 0$. The border of the shaded area satisfies to the following condition:

$$b^2 [(\langle l_\xi^2 \rangle \cos^2 \varphi + \langle l_\eta^2 \rangle \sin^2 \varphi) \cos^2 \mu - \langle l_\xi^2 \rangle \sin^2 \varphi - \langle l_\eta^2 \rangle \cos^2 \varphi] + a^2 \sin^2 \mu = 0. \quad (8)$$

In particular, if $\varphi = 90^\circ$ then $\sin^2 \Phi = 1$ for $\mu < \mu_c$ and $\sin^2 \Phi = 0$ for $\mu > \mu_c$, where

$$\sin \mu_c = \frac{b^2(1 - 2\langle l_\eta^2 \rangle)}{1 - b^2 - b^2\langle l_\eta^2 \rangle}. \quad (9)$$

The critical angle μ_c corresponds to an orientational transition: in the equilibrium SN state $\mathbf{d} \perp \hat{\eta}$ for $\mu < \mu_c$, while $\mathbf{d} \parallel \hat{\eta}$ for $\mu > \mu_c$.

The NMR frequency shift from the Larmor value can be obtained from (1) and (7):

$$\begin{aligned} \Delta\omega = & \frac{1}{4}K \left\{ (a^2 m_y^2 - b^2 \langle n_x^2 \rangle + b^2 \langle n_y^2 \rangle) \times \right. \\ & \times [1 - 2\sin^2 \Phi(1 + \cos \beta)] + \\ & \left. + [4 - 5a^2 m_y^2 - b^2(7\langle n_x^2 \rangle + 5\langle n_y^2 \rangle)] \cos \beta \right\}, \quad (10) \end{aligned}$$

where

$$K = \frac{2}{3 - 4a^2 b^2} \frac{\Omega_A^2}{\omega}$$

and $\omega = gH$. Let consider 4 cases: $\varphi = 0$, $0 < \mu < 90^\circ$ (the case A); $\mu = 90^\circ$, $0 < \varphi < 90^\circ$ (B); $\varphi = 90^\circ$, $0 < \mu < \mu_c$ (C_1); $\varphi = 90^\circ$, $\mu_c < \mu < 90^\circ$ (C_2). In Fig. 1 these orientations of \mathbf{H} correspond to arcs marked A, B, C_1 , and C_2 . Then for the case of continuous wave (CW) NMR ($\cos \beta \approx 1$) we get:

$$\begin{aligned} A: \quad \Delta\omega &= K(D \sin^2 \mu + E \cos^2 \mu), \\ B: \quad \Delta\omega &= KD(1 - 2\sin^2 \varphi), \\ C_1: \quad \Delta\omega &= KE \cos 2\mu, \\ C_2: \quad \Delta\omega &= K(E \cos^2 \mu - D), \end{aligned} \quad (11)$$

where $D = b^2(1 - 2\langle l_\eta^2 \rangle) \geq 0$ and $E = 1 - b^2 - b^2\langle l_\eta^2 \rangle > 0$. The dependence of $\Delta\omega$ on μ for $\varphi = 90^\circ$ is shown in Fig. 2. This dependence is fully determined by 2 values of the frequency shift: $\Delta\omega_\xi = -KD(\mathbf{H} \parallel \hat{\xi})$ and $\Delta\omega_\zeta = KE(\mathbf{H} \parallel \hat{\zeta})$ so that $\sin^2 \mu_c = -\Delta\omega_\xi / \Delta\omega_\zeta$.

In the isotropic 2D LIM state $\langle l_\eta^2 \rangle = 1/2$ (i.e. $D = 0$) and for $\mu = 90^\circ$ (the case B) $\Delta\omega = 0$ in agreement with [6]. If the 2D LIM state is anisotropic and $\langle l_\eta^2 \rangle < 1/2$, then for $\mu = 90^\circ$ the shift equals 0 for $\varphi = 45^\circ$. For other values of φ the shift is 0 only for pure polar phase ($b = 0$). In pure ABM phase or in the ABM phase with

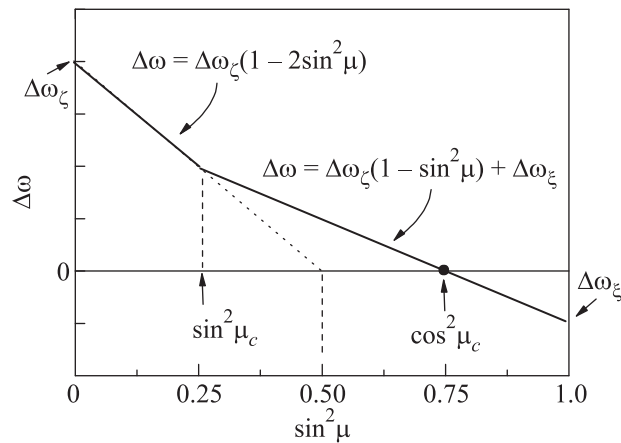


Fig. 2. CW NMR frequency shift versus μ for $\varphi = 90^\circ$ as follows from (11)

polar distortion the shift is positive (if $\varphi < 45^\circ$) or negative (if $\varphi > 45^\circ$).

3. Experimental setup. The experimental chamber used in the present work is similar to the chamber described in [6]. The chamber has two cells with N-aerogel samples. The samples (named below as 1 and 2) have a form of a cuboid with characteristic sizes of 4 mm. Initially the samples had an overall density $\sim 30 \text{ mg/cm}^2$ (sample 1) and $\sim 8 \text{ mg/cm}^2$ (sample 2), but were squeezed by $\sim 10\%$ and $\sim 5\%$ correspondingly along the direction transverse to the aerogel strands. In order to match Fig. 1 we choose the direction of the squeezing as $\hat{\eta}$ -axis, because (see next section) in the anisotropic 2D LIM state of the distorted ABM phase the maximum of $\langle l_\xi^2 \rangle$ corresponds to this direction of the squeezing.

We were able to rotate \mathbf{H} by any angle μ in $\hat{\zeta}$ – $\hat{\eta}$ plane (for the sample 1) or in $\hat{\zeta}$ – $\hat{\xi}$ plane (for the sample 2). Additional gradient coils were used to compensate the external magnetic field inhomogeneity. Experiments were performed in magnetic fields from 104 Oe up to 425 Oe (NMR frequencies were from 340 kHz up to 1.38 MHz) and at pressures from s.v.p. up to 29.3 bar. The necessary temperatures were obtained by a nuclear demagnetization cryostat and were measured by a quartz tuning fork, calibrated by Leggett frequency measurements in bulk ^3He –B. In order to avoid a paramagnetic signal from surface solid ^3He , the samples were preplated by ~ 2.5 atomic monolayers of ^4He .

A superfluid phase diagram of ^3He in the sample 1 was found to be almost the same as the diagram presented in [6]. For the sample 2 the diagram is slightly different (the superfluid transition temperatures are by 2–3% higher).

4. Experiments with the sample 1. Temperature dependencies of CW NMR frequency shifts in the ESP1 phase ($\Delta\omega_\zeta$ and $\Delta\omega_\eta$) for $\mathbf{H} \parallel \hat{\zeta}$ and for $\mathbf{H} \parallel \hat{\eta}$ are shown in Fig. 3. The superfluid transition temperature

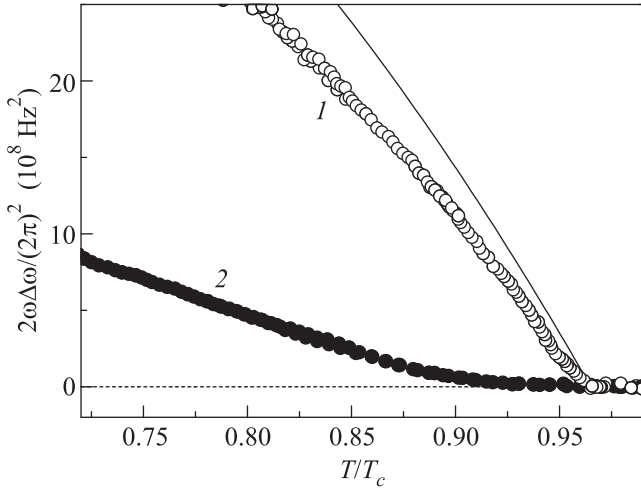


Fig. 3. CW NMR frequency shift versus temperature in sample 1: 1 – $\mathbf{H} \parallel \hat{\zeta}$; 2 – $\mathbf{H} \parallel \hat{\eta}$. Solid line shows Ω_A^2 rescaled from Ω_{A0}^2 in accordance with (12) for $k = 1$. $P = 14.2$ bar, $T_{ca} \approx 0.965 T_c$

of ^3He in this sample at the given pressure (14.2 bar) is $\sim 0.965 T_c$ as it can be seen from appearance of the NMR shift for $\mathbf{H} \parallel \hat{\zeta}$. As follows from (11) $\Delta\omega_\zeta = KE$ and $\Delta\omega_\eta = KD$. In the experiment, we obtain $\Delta\omega_\eta = 0$ down to $\sim 0.93 T_c$ but on further cooling the positive shift appears. It means that for $T < 0.93 T_c$ both b^2 and $1 - 2\langle l_\eta^2 \rangle$ are nonzero and positive, i.e. we get the distorted ABM phase and the squeezing of the sample along $\hat{\eta}$ results in preferable orientation of vectors \mathbf{l} along the $\hat{\xi}$ -axis.

If value of Ω_A is known, then we can find b^2 and $\langle l_\eta^2 \rangle$ from the measured values of $\Delta\omega_\zeta$ and $\Delta\omega_\eta$. The problem is that $\Omega_A \propto \Delta_0$ is known only for the bulk ^3He (below we denote this value by Ω_{A0}). The value of Ω_A in N-aerogel should be smaller due to the suppression of T_c and corresponding decrease of the gap. It is known that in silica aerogels the gap suppression is larger than the suppression of T_c in agreement with the “inhomogeneous isotropic scattering model” [15]. For example, for $T_{ca} = 0.965 T_c$ the gap and Ω_A is suppressed by $\sim 9\%$ [16]. Another model called the “homogeneous isotropic scattering model” [17] predicts that the gap suppression is proportional to T_{ca}/T_c . Both these models, however, can not be directly applied to ^3He in N-aerogel due to its strong anisotropy. Therefore we can only assume that

the suppression of Ω_A in N-aerogel is proportional to T_{ca}/T_c or larger, i.e.:

$$\Omega_A(T/T_{ca}) = k \frac{T_{ca}}{T_c} \Omega_{A0}(T/T_c), \quad (12)$$

where $k \leq 1$. In Fig. 4 we present values of b^2 and $\langle l_\eta^2 \rangle$

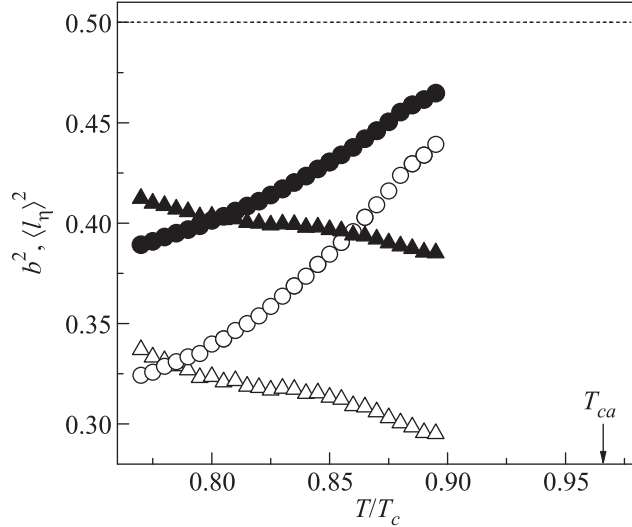


Fig. 4. $\langle l_\eta^2 \rangle$ (circles) and b^2 (triangles) calculated from data in Fig. 3 for $k = 1$ (filled symbols) and for $k = 0.9$ (open symbols)

calculated from the data shown in Fig. 3 in the assumption that $k = 1$ or 0.9 . It can be seen that $\langle l_\eta^2 \rangle$ grows on warming and tends to $1/2$ for both values of k , while b^2 decreases but can not be extrapolated to 0 for $T < T_{ca}$. Thus we conclude that the anisotropy of the 2D LIM state decreases on warming and $\langle l_\eta^2 \rangle$ becomes equal (or close) to $1/2$ at $T > 0.93 T_{ca}$ resulting in $\Delta\omega_\eta = 0$. The polar distortion grows on warming but it is unlikely that we get pure polar phase in a reasonably large temperature range near T_{ca} for these values of k . At lower pressures ($P \leq 9$ bar) we have obtained similar dependencies as shown in Figs. 2 and 3, but if $k \leq 0.9$ the value of b^2 can be extrapolated to 0 at $T < T_{ca}$, so the existence of the pure polar phase near T_{ca} can not be excluded, but only if $k \leq 0.9$.

5. Identification of the ESP2 phase (Experiments with the sample 2). The sample 2 was oriented so that \mathbf{H} could be rotated in $\hat{\zeta}$ - $\hat{\xi}$ plane. Correspondingly, at low temperatures the anisotropy of the 2D LIM state should result in a negative CW NMR frequency shift for the transverse orientation of the field ($\mathbf{H} \parallel \hat{\xi}$) as follows from C_2 in (11) for $\mu = 90^\circ$ and $D > 0$. Examples of temperature dependencies of the shift in ESP phases for both transverse and parallel orientations of \mathbf{H} are shown in Fig. 5. As it was expected,

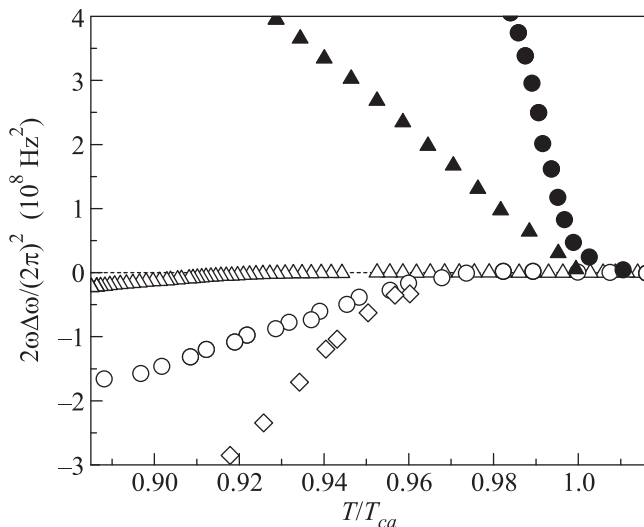


Fig. 5. CW NMR frequency shift versus temperature for sample 2. Open symbols – $\mathbf{H} \parallel \hat{\xi}$; filled symbols – $\mathbf{H} \parallel \hat{\zeta}$; $P = 12.3$ bar ($T_{ca} \approx 0.98 T_c$): circles – ESP1 phase; diamonds – ESP2 phase; $P = 2.9$ bar ($T_{ca} \approx 0.95 T_c$): triangles – ESP1 phase

at low temperatures the shift is negative in both ESP phases for $\mathbf{H} \parallel \hat{\zeta}$. In this case the absolute value of the shift in the ESP2 phase is larger than in the ESP1 phase.

The difference in the NMR shift in ESP1 and ESP2 phases may be explained either by different values of the polar distortion (i.e. of b^2) or by different values of the anisotropy of the 2D LIM state (i.e. of $\langle I_\eta^2 \rangle$). In both cases the dependence of $\Delta\omega = \Delta\omega(\mu)$ for the sample 2 should correspond to the dependence shown in Fig. 2.

As follows from (9) $\mu_c > 0$ only if $\Delta\omega_\xi \neq 0$. Therefore at $P = 12.3$ bar we have chosen the temperature $T \approx 0.85 T_{ca}$ where the absolute value of $\Delta\omega_\xi$ is large enough, but the smeared transition into the LTP just starts. In order to get the ESP2 phase at this temperature the sample was warmed up from the LTP above the point of full transition to the ESP2 phase ($\sim 0.93 T_{ca}$) and then was cooled down. The obtained dependencies of the CW NMR shift on μ are shown in Fig. 6 where solid lines are drawn using only the corresponding values of $\Delta\omega_\xi$ and $\Delta\omega_\zeta$. It can be seen that the data are well described by the theory. Further analysis shows that the difference between the ESP phases can not be attributed to the difference of magnitudes of the polar distortion, but can be explained in assumption that the anisotropy of the 2D LIM state in these phases is different. The data in Fig. 6 allow to calculate b^2 and $\langle I_\eta^2 \rangle$ for a given k in Eq. (12). If $k = 1$ then $b^2 = 0.42$ and $\langle I_\eta^2 \rangle = 0.43$ for the ESP1 phase, and $b^2 = 0.43$ and $\langle I_\eta^2 \rangle = 0.33$ for ESP2 phase. For $k = 0.9$ we get $b^2 = 0.35$ and $\langle I_\eta^2 \rangle = 0.39$ for the ESP1 phase, and $b^2 = 0.36$ and

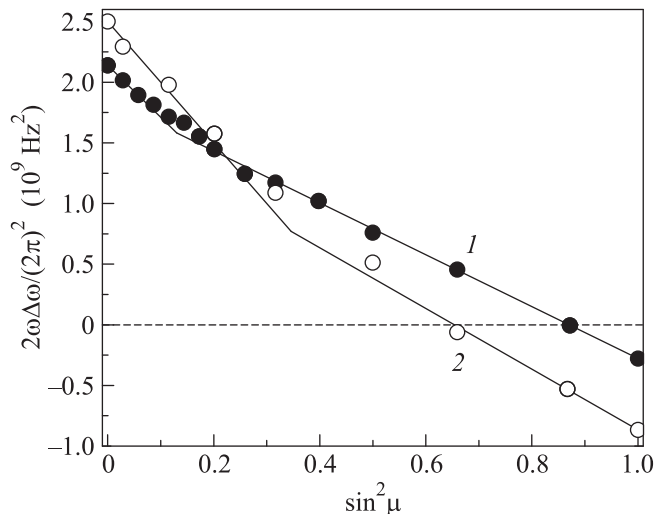


Fig. 6. CW NMR frequency shift versus μ for sample 2. Solid lines – theoretical dependencies C_1 and C_2 in (11); 1 – ESP1 phase; 2 – ESP2 phase; $P = 12.3$ bar, $T \approx 0.85 T_{ca}$, $H = 117$ Oe

$\langle I_\eta^2 \rangle = 0.24$ for ESP2 phase. Note that for both values of k we get nearly equal values of b^2 in both ESP states, while the anisotropy of the 2D LIM state in the ESP2 phase is always greater than in the ESP1 phase. This difference in the anisotropy may be due to the ESP2 phase is formed on warming from the LTP, which order parameter corresponds to a spatially homogeneous polar distorted BW phase. It is the first order transition, i.e. the phase boundary moves through the sample providing the orienting effect on \mathbf{l} and resulting in stabilization of more anisotropic metastable 2D LIM state. Worthy to mark that similar history dependent orientational effect on \mathbf{l} was observed in ABM phase of ^3He in silica aerogel [9].

The ESP phases observed in [6] have been obtained in N-aerogel sample which has not been squeezed in transverse direction. In this case the 2D LIM state of the ESP1 phase should be isotropic, but the above mentioned orienting effect should remain resulting in the anisotropic 2D LIM state in the ESP2 phase.

6. Orientation of orbital vector in N-aerogel.

The influence of aerogel deformation on a spatial distribution of \mathbf{l} in ABM-like phase of ^3He is a complex problem and depends on how \mathbf{l} -orienting centers are transformed during deformation. Different types of aerogel have different microscopic structures. This can result in a different response of the \mathbf{l} -field to the deformation. For example, silica aerogels used in [18, 19] orient \mathbf{l} along the axis of stretching and normal to the axis of squeezing. On the other hand, N-aerogel (i.e. the infinitely

stretched array of cylindric strands) orients \mathbf{l} normal to the strands, i.e. *normal* to the stretching.

There are three theoretical models describing the influence of the aerogel deformation on \mathbf{l} in the ABM phase of ^3He [8, 20, 21]. Theoretical studies of liquid crystals in strained aerogel (see e.g. [22]) are also conceptually related to this subject. The model [8] seems to be the most consistent with N-aerogel. This model considers the aerogel as a system of randomly oriented cylinders. The deformation changes their angular distribution and \mathbf{l} tends to align along the axis of squeezing and normal to the axis of stretching. N-aerogel corresponds to the infinite stretching and the model predicts the 2D LIM state in the $\hat{\xi}$ - $\hat{\eta}$ plane in agreement with experiments [6] and with this work.

As it was shown above the squeezing of the N-aerogel in the $\hat{\xi}$ - $\hat{\eta}$ plane results in preferable orientation of \mathbf{l} along the direction *normal* to the squeezing direction. At first glance, this disagrees with [8]. However there is no contradiction here: the point is that the squeezing in the $\hat{\xi}$ - $\hat{\eta}$ plane does not change orientations of the strands, i.e. the orienting effect in frames of the model [8] is absent. However, this type of deformation changes spatial correlations of strands. If these correlations are anisotropic in $\hat{\xi}$ - $\hat{\eta}$ plane then the orienting force is along the direction normal to the squeezing. This is illustrated by Fig. 7 where the result of the squeezing of N-aerogel

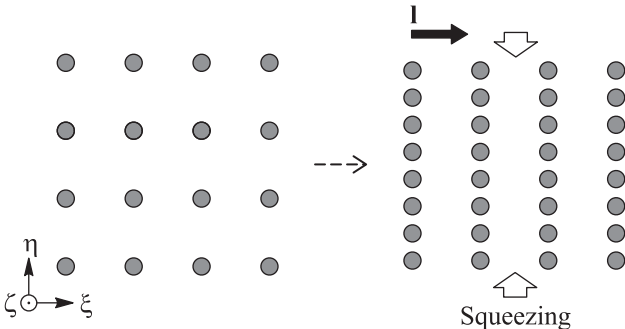


Fig. 7. Squeezing of N-aerogel which strands form the 2D square lattice

is shown for the simplest case of a 2D square lattice of the strands in $\hat{\xi}$ - $\hat{\eta}$ plane. It can be seen that a strong squeezing results in the formation of “wall-like” structures. These “walls” should orient \mathbf{l} normal to the surface, i.e. *normal* to the squeezing direction. We think that this effect remains not only for the square lattice but also for any locally anisotropic spatial correlations. The similar phenomenon may cause the orienting effect in a 3D lattice of ball-like \mathbf{l} -orienting centers if their spatial correlations are locally anisotropic. We assume that it may explain orienting effects observed in [18, 19].

7. Conclusions. 1. The observed NMR properties of ABM phase with polar distortion of ^3He in “nematically ordered” aerogel agree with the developed theoretical model. This allows us to explain the difference in NMR properties of the ESP phases: the 2D LIM state of vector \mathbf{l} in the ESP2 phase is more anisotropic than in the ESP1 phase.

2. We have shown that the squeezing of N-aerogel along the direction normal to the strands results in the anisotropic 2D LIM state in the ESP phases so that the preferable orientation of \mathbf{l} -field is normal to the squeezing. The explanation of this effect is suggested. The anisotropy of the 2D LIM state decreases on warming and may disappear below T_{ca} .

3. The order parameter orientational transition have been observed. The transition occurs when the angle between \mathbf{H} and the axis of the anisotropy reaches the critical value.

This work was supported in part by RFBR (grants # 13-02-00674 and 14-42-03045) and by RAS Program “Quantum mesoscopic and disordered structures”. We are grateful to I.A. Fomin and G.E. Volovik for useful comments.

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