

Random anisotropy disorder in superfluid $^3\text{He-A}$ in aerogel

G. E. Volovik¹⁾

Low Temperature Laboratory, Helsinki University of Technology, FIN-02015 HUT, Finland

Landau Institute for Theoretical Physics RAS, 119334 Moscow, Russia

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The anisotropic superfluid $^3\text{He-A}$ in aerogel provides an interesting example of a system with continuous symmetry in the presence of random anisotropy disorder. Recent NMR experiments allow us to discuss two regimes of the orientational disorder, which have different NMR properties. One of them, the (*s*)-state, is identified as the pure Larkin-Imry-Ma state. The structure of another state, the (*f*)-state, is not very clear: probably it is the Larkin-Imry-Ma state contaminated by the network of the topological defects pinned by aerogel.

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1. Introduction. Behavior of systems with continuous symmetry in the presence of random anisotropy disorder is the subject of numerous theoretical and experimental investigations. This is because of the surprising observation made by Larkin [1] and Imry and Ma [2] that even a weak disorder may destroy the long-range translational or orientational order. Recent example is provided by the nematic liquid crystals in random porous medium, in which the order parameter – the unit vector \hat{n} – interacts with the quenched random anisotropy disorder (see e.g. Ref. [3] and references therein). Though the problem of violation of the long-range order by quenched disorder is more than 30 years old, still there is no complete understanding (see e.g. Refs. [3–6] and references therein), especially concerning the role of topological defects.

In the anisotropic phase A of superfluid ^3He , the Larkin-Imry-Ma effect is even more interesting. In this superfluid the order parameter contains two Goldstone vector fields: (1) the unit vector \hat{I} characterizes the spontaneous anisotropy of the orbital and superfluid properties of the system; and (2) the unit vector \hat{d} characterizes the spontaneous anisotropy of the spin (magnetic) degrees of freedom. In aerogel, the quenched random anisotropy disorder of the silicon strands interacts with the orbital vector \hat{I} , which thus must experience the Larkin-Imry-Ma effect. As for the vector \hat{d} of the spontaneous anisotropy of spins it is assumed that \hat{d} does not interact directly with the quenched disorder, at least in the arrangement when the aerogel strands are covered by layers of ^4He atoms preventing the formation of solid layers of ^3He with large Curie-Weiss magnetiza-

tion. There is a tiny spin-orbit coupling between vectors \hat{d} and \hat{I} due to which the \hat{I} -vector may transfer the disorder to the \hat{d} -field. Superfluid $^3\text{He-A}$ experiences many different types of topological defects (see e.g. [7]), which may be pinned by the disorder.

On recent experiments on the superfluid $^3\text{He-A}$ in aerogel see Refs. [8–13] and references therein. In particular, Refs. [9, 13] describe the transverse NMR experiments, in which the dependence of the frequency shift on the tipping angle β of the precessing magnetization has been measured; and Ref. [13] also reports the observation of the longitudinal NMR. Here we discuss these experiments in terms of the Larkin-Imry-Ma disordered state [1, 2] extended for the description of the superfluid $^3\text{He-A}$ in aerogel [14]. In Sec. 2 the general equations for NMR in $^3\text{He-A}$ are written. In Sec. 3 these equations are applied to the states with disordered \hat{d} and \hat{I} fields. In Sec. 4 and Sec. 5 the models for the two observed states are suggested in terms of the averaged distributions of \hat{d} and \hat{I} fields consistent with observations. Finally in Sec. 6 these states are interpreted in terms of different types of disorder.

2. Larmor precession of $^3\text{He-A}$. In a typical experimental arrangement the spin-orbit (dipole-dipole) energy is smaller than Zeeman energy and thus may be considered as a perturbation. In zero-order approximation when the dipole energy and dissipation are neglected, the spin freely precesses with the Larmor frequency $\omega_L = \gamma H$, where γ is the gyromagnetic ratio of ^3He nuclei. In terms of the Euler angles the precession of magnetization is given by

$$\mathbf{S}(t) = S R_z(-\omega_L t + \alpha) R_y(\beta) \hat{z}. \quad (1)$$

¹⁾e-mail: volovik@boojum.hut.fi

Here $S = \chi H$ is the amplitude of spin induced by magnetic field; the axis \hat{z} is along the magnetic field \mathbf{H} ; matrix $R_y(\beta)$ describes rotation about transverse axis y by angle β , which is the tipping angle of the precessing magnetization; R_z describes rotation about z ; α is the phase of the precessing magnetization. According to the Larmor theorem [15], in the precessing frame the vector $\hat{\mathbf{d}}$ is in turn precessing about \mathbf{S} . Because of the interaction between the spin \mathbf{S} and the order parameter vector $\hat{\mathbf{d}}$, the precession of $\hat{\mathbf{d}}$ occurs in the plane perpendicular to \mathbf{S} , and it is characterized by another phase Φ_d . In the laboratory frame the precession of $\hat{\mathbf{d}}$ is given by

$$\hat{\mathbf{d}}(t) = R_z(-\omega_L t + \alpha) R_y(\beta) R_z(\omega_L t - \alpha + \Phi_d) \hat{\mathbf{x}}, \quad (2)$$

while the orbital vector $\hat{\mathbf{l}}$ is time-independent in this approximation:

$$\hat{\mathbf{l}} = \hat{z} \cos \lambda + \sin \lambda (\hat{x} \cos \Phi_l + \hat{y} \sin \Phi_l). \quad (3)$$

This is the general state of the pure Larmor precession of ${}^3\text{He-A}$, and it contains 5 Goldstone parameters: 2 angles α and β of the magnetization in the precessing frame; angle Φ_d which characterizes the precession of vector $\hat{\mathbf{d}}$; and two angles λ and Φ_l of the orbital vector $\hat{\mathbf{l}}$.

The degeneracy is lifted by spin-orbit (dipole) interaction [16]

$$F_D = -\frac{\chi \Omega_A^2}{2\gamma^2} (\hat{\mathbf{l}} \cdot \hat{\mathbf{d}})^2, \quad (4)$$

where Ω_A is the so-called Leggett frequency. In the bulk homogeneous ${}^3\text{He-A}$, the Leggett frequency coincides with the frequency of the longitudinal NMR, $\omega_{\parallel} = \Omega_A$. In typical experiments one has $\Omega_A \ll \omega_L$, which allows us to use the spin-orbit interaction averaged over the fast Larmor precession [17, 15]:

$$\bar{F}_D = \frac{\chi \Omega_A^2}{2\gamma^2} U, \quad (5)$$

$$U = -\frac{1}{2} \sin^2 \beta + \frac{1}{4} (1 + \cos \beta)^2 \sin^2 \Phi \sin^2 \lambda - \left(\frac{7}{8} \cos^2 \beta + \frac{1}{4} \cos \beta - \frac{1}{8} \right) \sin^2 \lambda, \quad (6)$$

where $\Phi = \Phi_d - \Phi_l$.

The dipole interaction generates the frequency shift of the transverse NMR from the Larmor frequency:

$$\omega_{\perp} - \omega_L = -\frac{\partial \bar{F}_D}{\partial (S \cos \beta)} = -\frac{\Omega_A^2}{2\omega_L} \frac{\partial U}{\partial \cos \beta}. \quad (7)$$

In the bulk ${}^3\text{He-A}$, the minimum of the dipole interaction requires that $\Phi_d = \Phi_l$, and $\sin^2 \lambda = 1$, i.e. the

equilibrium position of $\hat{\mathbf{l}}$ is in the plane perpendicular to \mathbf{H} . However, for the ${}^3\text{He-A}$ confined in aerogel, the interaction with the quenched disorder may essentially modify this spatially homogeneous state by destroying the long-range orientational order due to the Larkin-Imry-Ma effect [14].

3. Two states of ${}^3\text{He-A}$ in aerogel. Experiments reported in Ref. [13] demonstrate two different types of magnetic behavior of the A-like phase in aerogel, denoted as ($f+c$)-state and (c)-state correspondingly. The ($f+c$)-state contains two overlapping lines (f) and (c) in the transverse NMR spectrum (far from and close to the Larmor frequency correspondingly). The frequency shift of the transverse NMR is about 4 times bigger for the (f)-line compared to the (c)-line. The behavior under applied gradient of magnetic field suggests that the ($f+c$)-state consists of two magnetic states concentrated in different parts of the cell. The (c)-state contains only a single (c)-line in the spectrum, and it is obtained after application of the 180 degree pulse while cooling through T_c . The pure (f)-state, i.e. the state with a single (f)-line, has not been observed.

The (c) and ($f+c$)-states have different dependence of the frequency shift $\omega_{\perp} - \omega_L$ on the tilting angle β in the pulsed NMR experiments: $\omega_{\perp} - \omega_L \propto \cos \beta$ in the pure (c)-state; and $\omega_{\perp} - \omega_L \propto (1 + \cos \beta)$ in the ($f+c$)-state. The latter behavior probably characterizes the property of the (f)-line which has the bigger shift and is dominating in the spectrum of the ($f+c$)-state. The $(1 + \cos \beta)$ -law has been also observed in Ref. [9].

The experiments with longitudinal NMR were also reported in Ref. [13]. The longitudinal resonance in the ($f+c$)-state has been observed, however no traces of the longitudinal resonance have been seen in the (c)-state.

Let us discuss this behavior in terms of the disordered states emerging in ${}^3\text{He-A}$ due to the orientational disorder.

In the extreme case of weak disorder, the characteristic Imry-Ma length L of the disordered $\hat{\mathbf{l}}$ -texture is much bigger than the characteristic length scale ξ_D of dipole interaction, $L \gg \xi_D$. In this case the equilibrium values of Φ and λ are dictated by the spin-orbit interaction, $\Phi = \Phi_d - \Phi_l = 0$ and $\sin^2 \lambda = 1$; and the Eq.(7) gives

$$\omega_{\perp} - \omega_L = \frac{\Omega_A^2}{8\omega_L} (1 + 3 \cos \beta). \quad (8)$$

This dependence fits neither the (f)-state $(1 + \cos \beta)$ behavior nor the $\cos \beta$ law in the (c)-state, which indicates that the disorder is not weak compared to the spin-orbit energy.

In the extreme case of strong disorder, when $L \ll \xi_D$, both Φ and $\hat{\mathbf{l}}$ are random:

$$\langle \sin^2 \Phi \rangle - 1/2 = 0, \quad (9)$$

$$\langle \sin^2 \lambda \rangle - 2/3 = 0. \quad (10)$$

In this case it follows from Eq.(7) that the frequency shift is absent:

$$\omega_{\perp} = \omega_L. \quad (11)$$

The Eq.(9) means that Φ_l and Φ_d are dipole unlocked, i.e. they are not locked by the spin-orbit dipole-dipole interaction, which is natural in case of small Imry-Ma length, $L \ll \xi_D$. In principle, there can be three different dipole-unlocked cases: (i) when both Φ_l and Φ_d are random and independent; (ii) when Φ_l is random while Φ_d is fixed; (iii) when Φ_d is random while Φ_l is fixed.

The strong disorder limit is consistent with the observation that the frequency shift of the (c)-line is much smaller than the frequency shift in $^3\text{He-B}$ in aerogel. The observed non-zero value can be explained in terms of the small first order corrections to the strong disorder limit. Let us introduce the parameters

$$a = 1/2 - \langle \sin^2 \Phi \rangle, \quad b = \langle \sin^2 \lambda \rangle - 2/3, \quad (12)$$

which describe the deviation from the strong disorder limit. These parameters are zero in the limit of strong disorder $L^2/\xi_D^2 \rightarrow 0$, and one may expect that in the pure Larkin-Imry-Ma state they are proportional to the small parameter $L^2/\xi_D^2 \ll 1$. The behavior of these two parameters can be essentially different in different realizations of the disordered state, since the vector \hat{l} entering the parameter a interacts with the quenched orientational disorder directly, while Φ_d only interacts with Φ_l via the spin-orbit coupling. That is why we shall try to interpret the two observed magnetic states, the (c)-state and the (f)-state, in terms of different realizations of the textural disorder described by different phenomenological relations between parameters a and b in these two states.

4. Interpretation of (c)-state. The observed $\cos \beta$ -dependence of the transverse NMR frequency shift in the (c)-state [13] can be reproduced if we assume that in the (c)-state the parameters a_c and b_c satisfy the following relation: $a_c \ll b_c$. Then in the main approximation,

$$a_c = 0, \quad b_c \neq 0, \quad (13)$$

the effective potential U in Eq.(6) is

$$U_c = -\frac{3}{4}b_c \cos^2 \beta + \frac{1}{4} \left(b_c + \frac{1}{3} \right). \quad (14)$$

If the parameter b_c does not depend on β , the variation of U_c with respect to $\cos \beta$ gives the required $\cos \beta$ -de-

pendence of the frequency shift of transverse NMR in the (c)-state:

$$\omega_{c\perp} - \omega_L = \frac{3\Omega_A^2}{4\omega_L} b_c \cos \beta. \quad (15)$$

Let us estimate the parameter b_c using the following consideration. The dipole energy which depends on λ violates the complete randomness of the Larkin-Imry-Ma state, and thus perturbs the average value of $\sin^2 \lambda$. The deviation of b_c from zero is given by:

$$b_c \sim \frac{L^2}{\xi_D^2} \left(\cos^2 \beta - \frac{1}{3} \right). \quad (16)$$

In this model the potential and frequency shift become

$$U_c \sim -\frac{L^2}{\xi_D^2} \left(\cos^2 \beta - \frac{1}{3} \right)^2, \quad (17)$$

$$\omega_{c\perp} - \omega_L \sim \frac{\Omega_A^2 L^2}{\omega_L \xi_D^2} \cos \beta \left(\cos^2 \beta - \frac{1}{3} \right). \quad (18)$$

Such β -dependence of the transverse NMR is also anti-symmetric with respect to the transformation $\beta \rightarrow \pi - \beta$ as in the model with the β -independent parameter b_c in Eq.(15); however, as distinct from that model it is inconsistent with the experiment (see Fig.6 in Ref. [13]). Certainly the theory must be refined to estimate the first order corrections to the zero values of the parameters a_c and b_c .

The frequency of the longitudinal NMR in the (c)-state is zero in the local approximation. The correction due to the deviation of Φ from the random behavior, i.e. due to the non-zero value of the parameter a_c in the (c)-state, is:

$$\omega_{c\parallel}^2 = \frac{2}{3} (1 - 2 \langle \sin^2 \Phi \rangle) \Omega_A^2 = \frac{4a_c}{3} \Omega_A^2. \quad (19)$$

In the simplest Imry-Ma model $a_c \sim L^2/\xi_D^2 \ll 1$, and thus the frequency of longitudinal NMR is small as compared with the frequency of the longitudinal resonance in (f+c)-state, discussed in the next section. This is consistent with non-observation of the longitudinal NMR in the (c)-state: under conditions of this experiment the longitudinal resonance cannot be seen if its frequency is much smaller than the frequency of the longitudinal resonance observed in the (f+c)-state [13].

5. Interpretation of (f)-state. The observed $(1 + \cos \beta)$ -dependence of the transverse NMR frequency shift of the (f)-line dominating in the (f+c)-state [13, 9] is reproduced if we assume that for the (f)-line one has

$a_f \gg b_f$. In this case, in the main approximation the (f)-state may be characterized by

$$a_f \neq 0, \quad b_f = 0, \quad (20)$$

and one obtains:

$$\omega_{f\perp} - \omega_L = \frac{a_f \Omega_A^2}{6 \omega_L} (1 + \cos \beta). \quad (21)$$

Let us compare the frequency shift of the (f)-line with that of the (c)-line in Eq. (15) at $\beta = 0$:

$$\frac{\omega_{f\perp} - \omega_L}{\omega_{c\perp} - \omega_L} = \frac{4a_f}{9b_c}. \quad (22)$$

According to experiments [13] this ratio is about 4, and thus one obtains the estimate: $b_c \sim 0.1a_f$. This supports the strong disorder limit, $b_c \ll 1$, for the (c)-state. If the statistic properties of the $\hat{\mathbf{l}}$ -texture in the (f)-state are similar to that in the (c)-state, then one has $b_f \ll a_f$ as suggested in Eq. (20).

The frequency of the longitudinal NMR in such a state is

$$\omega_{f\parallel}^2 = \frac{4a_f}{3} \Omega_A^2, \quad (23)$$

which gives the relation between the transverse and longitudinal NMR frequencies

$$\omega_{f\perp} - \omega_L = \frac{1}{8} \frac{\omega_{f\parallel}^2}{\omega_L} (1 + \cos \beta). \quad (24)$$

This relation is also valid for the Fomin's robust phase [18] (see [19]). However, the frequency of the longitudinal NMR measured in the ($f + c$)-state [13] does not satisfy this relation: the measured value of $\omega_{f\parallel}$ is about 0.65 of the value which follows from the Eq. (24) if one uses the measured $\omega_{f\perp} - \omega_L$. Probably the situation can be improved, if one considers the interaction between the f and c lines in the ($f + c$)-state (see Ref.[13]).

6. Discussion. *1. Interpretation of A-phase states in aerogel.* The observed two magnetic states of $^3\text{He-A}$ in aerogel [13] can be interpreted in the following way. The pure (c)-state is the Larkin-Imry-Ma phase with strong disorder, $L \ll \xi_D$. The ($f + c$)-state can be considered as mixed state with the volume V_c filled by the Larkin-Imry-Ma phase, while the rest of volume $V_f = V - V_c$ consists of the (f)-state. The (f)-state is also random due to the Larkin-Imry-Ma effect, but the spin variable Φ_d and the orbital variable Φ_l are not completely independent in this state. If Φ_d partially follows Φ_l , the difference $\Phi = \Phi_d - \Phi_l$ is not random and the parameter a_f in the (f)-state is not very small, being equal to $1/2$ in the extreme dipole-locked case. Thus,

for the ($f + c$)-state one may assume that $a_f \gg b_f, b_c, a_c$. As a result the (f)-line has essentially larger frequency shift of transverse NMR and essentially larger longitudinal frequency compared to the (c)-line.

Both results are consistent with the experiment: from the transverse NMR it follows that $b_c \sim 0.1a_f$ (see Eq. (22)); and from the lack of the observation of longitudinal NMR in the (c)-state it follows that $a_c \ll a_f$. This confirms the assumption of the strong disorder in the (c)-state, in which the smallness of the parameters b_c and a_c is the result of the randomness of the $\hat{\mathbf{l}}$ -texture on the length scale $L \ll \xi_D$. The $\cos \beta$ -dependence of $\omega_{c\perp} - \omega_L$ in Eq. (15) and $(1 + \cos \beta)$ -dependence of $\omega_{f\perp} - \omega_L$ in Eq. (21) are also consistent with the experiment. The open problem is how to estimate theoretically the phenomenological parameters a_f, b_f, b_c, a_c and find its possible dependence on β .

The 'universal' relation (24) between the longitudinal and transverse NMR frequencies is not satisfied in the experiment, but we cannot expect the exact relation in such a crude model, in which the interaction between the f and c lines in the ($f + c$)-state is ignored (see [13]). Moreover, we use the local approximation, i.e. we do not take into account the fine structure of the NMR line which may contain the satellite peaks due to the bound states of spin waves in the texture of $\hat{\mathbf{l}}$ and $\hat{\mathbf{d}}$ vectors. The tendency is however correct: the smaller is the frequency shift of transverse NMR the smaller is the frequency of longitudinal NMR.

2. Global anisotropy and negative frequency shift.

For further consideration one must take into account that in some aerogel samples the large negative frequency shift has been observed for the A-phase [10–12, 20]. The reason of the negative shift is the deformation of the aerogel sample which leads to the global orientation of the orbital vector $\hat{\mathbf{l}}$ in the large region of the aerogel [20]. The effect of regular uniaxial anisotropy in aerogel has been considered in Refs. [21, 22]. It is important that even a rather small deformation of aerogel may kill the subtle collective Larkin-Imry-Ma effect and lead to the uniform orientation of the $\hat{\mathbf{l}}$ -vector. Using the estimation of the Imry-Ma length in Ref. [14], one can find that the critical stretching of the aerogel required to kill the Larkin-Imry-Ma effect is proportional to $(R/\xi_0)^3$. Here R is the radius of the silica strands and ξ_0 is the superfluid coherence length.

From Eqs. (6) and (7) it follows that the maximum possible negative frequency shift could occur if in some region the regular orientation of $\hat{\mathbf{l}}$ induced by deformation of the aerogel is along the magnetic field (i.e. $\lambda = 0$):

$$\omega_{\perp} - \omega_L = -\Omega_A^2/2\omega_L. \quad (25)$$

Such longitudinal orientation of $\hat{\mathbf{l}}$ is possible because the regular anisotropy caused by the deformation of aerogel is bigger than the random anisotropy, which in turn in the strong disorder limit is bigger than the dipole energy preferring the transverse orientation of $\hat{\mathbf{l}}$.

Comparing the measured magnitude of the negative shift (which cannot be bigger than the maximum possible in Eq. (25)) with the measured positive shift of the (f)-line in the ($f+c$)-state [23, 13] one obtains that the parameter a_f in Eq. (21) must be smaller than 0.25. This is also confirmed by the results of the longitudinal NMR experiments [13], which show that the frequency of the longitudinal NMR in the ($f+c$)-state of $^3\text{He-A}$ is much smaller than the frequency of the longitudinal NMR in $^3\text{He-B}$. The latter is only possible if $a_f \ll 1$ in Eq.(23), i.e. the (f)-state is also in the regime of strong disorder. Thus there is only the partial dipole locking between the spin variable Φ_d and the orbital variable Φ_l in the (f)-state.

3. Possible role of topological defects. It is not very clear what is the origin of the (f)-state. The partial dipole locking is possible if the characteristic size of the $\hat{\mathbf{l}}$ texture in the (f)-state is on the order of or somewhat smaller than ξ_D .

Alternatively, the line (f) could come from the topological defects of the A-phase (vortices, solitons, vortex sheets, etc., see Ref. [7]). The defects could appear during cooling down the sample from the normal (non-superfluid) state and are annealed by application of the 180 degree pulse during this process. Appearance of a large amount of pinned topological defects in $^3\text{He-B}$ in aerogel has been suggested in Ref. [24]. The reason why the topological defects may effect the NMR spectrum in $^3\text{He-A}$ is the following. In the case of the strong disorder limit the texture is random, and the frequency shift is zero, if one neglects the L/ξ_D corrections in the main approximation. The topological defect introduces some kind of order: some correlations are nonzero because of the conserved topological charge of the defect. That is why the frequency shift will be nonzero. It will be small, but still bigger than due to the corrections of order $(L/\xi_D)^2$.

If this interpretation is correct, there are two different realizations of the disordered state in the system with quenched orientational disorder: the network of the pinned topological defects and the pure Larkin-Imry-Ma state. Typically one has the interplay between these two realizations, but the defects can be erased by the proper annealing leaving the pure Larkin-Imry-Ma state.

4. Superfluid properties of A-phase in aerogel. The interesting problem concerns the superfluid density ρ_s in the states with the orientational disorder in the vec-

tor $\hat{\mathbf{l}}$. In Ref. [14] it was suggested that $\rho_s = 0$ in such a state. Whether the superfluid density is zero or not depends on the rigidity of the $\hat{\mathbf{l}}$ -vector. If the $\hat{\mathbf{l}}$ -texture is flexible, then due to the Mermin-Ho relation between the texture and the superfluid velocity, the texture is able to respond to the applied superflow by screening the supercurrent. As a result the superfluid density in the flexible texture could be zero. The experiments on $^3\text{He-A}$ in aerogel demonstrated that $\rho_s \neq 0$ (see e.g. [8] and references therein).

However, most probably these experiments have been done in the ($f+c$)-state. If our interpretation of this state in terms of the topological defects is correct, the non-zero value of superfluid density could be explained in terms of the pinning of the defects which leads to the effective rigidity of the $\hat{\mathbf{l}}$ -texture in the ($f+c$)-state. Whether the superfluid density is finite in the pure Larkin-Imry-Ma state, identified here as the (c)-state, remains an open experimental and theoretical question. The theoretical discussion of the rigidity or quasi-rigidity in such a state can be found in Refs. [25, 6]. In any case, one may expect that the observed two states of $^3\text{He-A}$ in aerogel, (c) and ($f+c$), have different superfluid properties.

Recent Lancaster experiments with vibrating aerogel sample indicate that the sufficiently large superflow produces the state with the regular (non-random) orientation of $\hat{\mathbf{l}}$ in aerogel, and in the oriented state the superfluid density is bigger [26]. This suggests that the orientational disorder does lead to at least partial suppression of the superfluid density.

5. Conclusion. In conclusion, the NMR experiments on the A-like superfluid state in the aerogel indicate two types of behavior. Both of them can be interpreted in terms of the random texture of the orbital vector $\hat{\mathbf{l}}$ of the $^3\text{He-A}$ order parameter. This supports the idea that the superfluid $^3\text{He-A}$ in aerogel exhibits the Larkin-Imry-Ma effect: destruction of the long-range orientational order by random anisotropy produced by the randomly oriented silicon strands of the aerogel. The extended numerical simulations are needed to clarify the role of the topological defects in the Larkin-Imry-Ma state, and to calculate the dependence of the NMR line-shape and superfluid density on concentration and pinning of the topological defects.

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