



HAL
open science

Influence of exchange scattering on superfluid ^3He states in nematic aerogel

V. P. Mineev

► **To cite this version:**

V. P. Mineev. Influence of exchange scattering on superfluid ^3He states in nematic aerogel. *Physical Review B: Condensed Matter and Materials Physics (1998-2015)*, 2018, 98, pp.014501. 10.1103/PhysRevB.98.014501 . hal-02022866

HAL Id: hal-02022866

<https://hal.science/hal-02022866>

Submitted on 12 Apr 2019

HAL is a multi-disciplinary open access archive for the deposit and dissemination of scientific research documents, whether they are published or not. The documents may come from teaching and research institutions in France or abroad, or from public or private research centers.

L'archive ouverte pluridisciplinaire **HAL**, est destinée au dépôt et à la diffusion de documents scientifiques de niveau recherche, publiés ou non, émanant des établissements d'enseignement et de recherche français ou étrangers, des laboratoires publics ou privés.

Influence of Exchange Scattering on Superfluid ^3He states in Nematic Aerogel

V.P.Mineev^{1,2}

¹*Univ. Grenoble Alpes, CEA, INAC, PHELIQS, GT, F-38000 Grenoble, France*

²*Landau Institute for Theoretical Physics, 142432 Chernogolovka, Russia*

(Dated: November 9, 2018)

The superfluid state in bulk liquid ^3He is realized in the form of A or B phases. Uniaxially anisotropic aerogel (nafen) stabilizes transition from the normal to the polar superfluid state which on further cooling transitions to the axipolar orbital glass state (Phys. Rev. Lett. **115**, 165304 (2015)). This is the case in nafen aerogel preplated by several atomic layers of ^4He . When pure liquid ^3He fills the same nafen aerogel a solid-like layer of ^3He atoms coats the aerogel structure. The polar state is not formed anymore and a phase transition occurs directly to the axipolar phase (Phys. Rev. Lett. **120**, 075301 (2018)). The substitution of ^4He by ^3He atoms at the aerogel surface changes the potential and adds the exchange scattering of quasiparticles on the aerogel strands. A calculation shows that both of these effects can decrease the degree of anisotropy of scattering and suppress the polar phase formation. The derived anisotropy of the spin diffusion coefficient in globally anisotropic aerogel is determined by the same parameter which controls the polar state emergence which allows one to check the effect of anisotropy change for different types of covering.

I. INTRODUCTION.

The superfluid state of liquid ^3He is formed by means of the Cooper pairing with spin and orbital angular momentum equal to 1. In isotropic space the phase transition depending on pressure occurs in either the A or B superfluid phase¹. Investigation of superfluid phases in high porosity aerogel allows one to study the influence of impurities on superfluidity with nontrivial pairing^{2,3}. It has been found that similar to bulk ^3He two superfluid A-like and B-like phases exist in ^3He in aerogel⁴. However, both the superfluid fraction and the temperature at which the superfluid is manifested are suppressed from their bulk values². The interesting possibility is opened in globally anisotropic aerogel of lifting the degeneracy between the different superfluid phases with p -pairing. It was shown⁵ that in the case of easy-axis anisotropy a new superfluid phase of ^3He , the polar phase, is stabilized below the transition temperature. It was also predicted⁵ that on further cooling a second-order transition into a polar-distorted A phase should occur. Indeed, quite recently, the first observation of the polar phase was reported⁶. This phase appears in ^3He confined in new type of "nematically ordered" aerogel called "nafen" with a nearly parallel arrangement of strands which play the role of ordered impurities. It was shown that in nafen the transition to the superfluid state always occurs to the polar phase and the region of its existence increases with density of strands. In another type of nematically ordered but less dense and much less anisotropic aerogel called "Obninsk aerogel"⁶ the superfluid state is always formed in the orbital glassy A-like state.

To avoid a paramagnetic signal from surface solid ^3He , the nafen samples in the measurements⁶ were preplated by ~ 2.5 ^4He monolayers. Then the new experiment series was performed with the same samples filled by pure ^3He ⁷. In this case the temperature of the superfluid transition is suppressed more strongly and this ef-

fect increases with strands density such that in the most dense nafen the superfluid transition was not detected down to the lowest attained temperatures. The superfluid transition occurs directly to the polar-distorted A-phase without the formation of an intermediate region of polar-state. The small addition of ^3He in the surface ^4He layer, corresponding to 0.1 monolayer, also completely kills superfluidity at 29,3 bar in the most dense nafen, and in the less dense aerogel noticeably suppresses the critical temperature. In this case also, the transition occurs directly to the distorted A-state. Thus, the situation looks as if the ^3He covering suppresses anisotropy necessary for the existence of the polar phase.

There was pointed out in Ref.7 : "The observed phenomena cannot be explained by a change of the scattering specularity because they are observed also at high pressures where the scattering should be diffusive regardless of the presence or absence of solid ^3He ". This statement is based on previous studies (see references 5-9 in the paper⁷) of the degree of specularity of ^3He quasiparticles scattering on metallic surfaces with different coverings. The corresponding information for liquid ^3He filling nafen aerogel is absent. But, generally speaking, the substitution of ^4He by ^3He atoms at aerogel surface changes the potential and adds the exchange scattering of quasiparticles on the aerogel strands.

I study this problem taking into account both the potential and the exchange scattering of quasiparticles of liquid ^3He on ^3He atoms localized at the strands surface. In Section II I show that ^3He covering changes the intensity and the anisotropy of scattering. In the Section III the presented derivation of spin diffusion current shows that the anisotropy of the spin diffusion coefficient in globally anisotropic aerogel is expressed through the same parameter which determines the polar state emergence. Thus, being measured, the anisotropy decrease of spin diffusion in nafen filled by pure ^3He can serve as a direct indication of suppression of the temperature

interval of polar state existence.

II. SUPERFLUID ^3He IN UNIAXIALLY ANISOTROPIC AEROGEL WITH MAGNETIC SCATTERING.

The order parameter of superfluid phases of ^3He is given¹ by the complex 3×3 matrix $A_{\alpha i}$, where α and i are the indices numerating the Cooper pair wave function projections on spin and orbital axes respectively. All the phases with different order parameters $A_{\alpha i}$ have the same critical temperature. The degeneracy is lifted by the fourth-order terms with respect to $A_{\alpha i}$ in the Landau expansion of the free energy density. The most energetically profitable are the B-phase with the order parameter $A_{\alpha i}^B = \Delta R_{\alpha i} e^{i\varphi}$, where $R_{\alpha i}$ is a rotation matrix, and (in the high pressures region) the A-phase with the order parameter

$$A_{\alpha i}^A = \Delta V_{\alpha} (m_i + in_i), \quad (1)$$

where \mathbf{V} is the unit spin vector and \mathbf{m} and \mathbf{n} are the orthogonal unit vectors such that $\mathbf{m} \times \mathbf{n} = \mathbf{l}$ is the unit vector directed along the Cooper pair angular momentum.

The different pairing states of superfluid ^3He in a random medium with global uniaxial anisotropy can be compared by making use the second-order terms in the Landau free energy density. They consist of an isotropic part, common to all the superfluid phases with p -pairing, and the anisotropic part

$$\begin{aligned} F^{(2)} &= F_i^{(2)} + F_a^{(2)} \\ &= \alpha_0 \left(\frac{T}{T_c} - 1 \right) A_{\alpha i} A_{\alpha i}^* + \eta_{ij} A_{\alpha i} A_{\alpha j}^*, \end{aligned} \quad (2)$$

where $T_c = T_c(P)$ is the transition temperature in the superfluid state suppressed with respect of transition temperature in the bulk liquid $T_{c0}(P)$. The medium uniaxial anisotropy with anisotropy axis parallel to \hat{z} direction coincident in our case with the average direction of aerogel strands is given by the traceless tensor

$$\eta_{ij} = \eta \begin{pmatrix} 1 & 0 & 0 \\ 0 & 1 & 0 \\ 0 & 0 & -2 \end{pmatrix}. \quad (3)$$

In the absence of global anisotropy ($\eta = 0$) all p -wave phases have the same critical temperature. At positive $\eta > 0$ the polar state with the order parameter of the form

$$A_{\alpha i} = a V_{\alpha} z_i, \quad (4)$$

where V_{α} is the unit spin vector, has the lowest energy of anisotropy⁸,

$$F_a = -2\eta |a|^2. \quad (5)$$

Hence, it has the highest critical temperature T_{c1} of transition from the normal state. At some lower temperature T_{c2} the polar state passes to the more energetically profitable distorted A-state⁹ with the order parameter

$$A_{\alpha i} = V_{\alpha} [a \hat{z}_i + ib(\hat{x}_i \cos \varphi + \hat{y}_i \sin \varphi)] \quad (6)$$

intermediate between the polar state at $b = 0$ and the A-state at $a = b$. This state has the Cooper pair angular momentum $\hat{\mathbf{l}} = -\hat{x} \sin \varphi(\mathbf{r}) + \hat{y} \cos \varphi(\mathbf{r})$ lying in the basal plane and locally ordered ($\varphi(\mathbf{r}) = \text{const}$) on the scales L exceeding the coherence length ξ_0 but smaller than the dipole length ξ_d and randomly distributed on scales larger than L . The pure polar state exists in the temperature interval roughly determined by the energy of anisotropy difference between of the polar and the distorted A-states⁹,

$$T_{c1} - T_{c2} \approx \frac{\eta}{\alpha_0} T_c. \quad (7)$$

Hence, at small η parameter the temperature interval of the polar state existence is small and hardly observable.

The quasiparticle interaction with the nafen strands is modelled by the interaction with the randomly distributed impurities including the globally anisotropic potential and the globally anisotropic exchange part,

$$\begin{aligned} H_{int} &= \sum_i \int d^3r \psi_{\alpha}^{\dagger}(\mathbf{r}) [u(\mathbf{r} - \mathbf{r}_i) \delta_{\alpha\beta} \\ &\quad + J(\mathbf{r} - \mathbf{r}_i) \sigma_{\alpha\beta} \mathbf{S}] \psi_{\beta}(\mathbf{r}), \end{aligned} \quad (8)$$

where \mathbf{S} is the spin of the impurity and σ are the ^3He quasiparticles spin matrices. The exchange scattering in an isotropic aerogel has been considered by Sauls and Sharma¹⁰ and by Baramidze and Kharadze¹¹. They have shown that if the scattering amplitude on impurities includes an exchange part then the critical temperatures splitting of A_1 and A_2 transitions under an external field H decreases in comparison with the impurity free case:

$$T_{A_1} - T_{A_2} = (\gamma_0 - \gamma_{imp}) H.$$

The effect arises due to an interference between scalar and exchange scattering such that

$$\gamma_{imp} \propto uJ$$

is proportional to the product of the corresponding amplitudes. In the NMR experiments^{6,7} the field is small and this effect is negligible, but one needs to consider an influence of the globally anisotropic scattering on critical temperature.

To find the critical temperature of superfluid transition in globally anisotropic aerogel one must calculate the second-order terms in the Landau free energy density

$$F^{(2)} = \frac{1}{3} \left\{ \frac{1}{g} \delta_{ij} - T \sum_{\omega} \int \frac{d^3 p}{(2\pi)^3} \hat{p}_i \Gamma_j(\mathbf{p}, \omega,) G(\mathbf{p}, \omega) G(-\mathbf{p}, -\omega) \right\} A_{\mu i}^* A_{\mu j}, \quad (9)$$

where g is the constant of p -wave triplet pairing. Here,

$$G(\mathbf{p}, \omega) = \frac{1}{i\omega - \xi_{\mathbf{p}} - \Sigma_{\mathbf{p}}(\omega)} \quad (10)$$

is the normal state quasiparticle Green function and $\Gamma_j^{\mu\nu}(\mathbf{p}, \omega)$ is the vertex part. $\xi_{\mathbf{p}} = \varepsilon_{\mathbf{p}} - \mu$ is the quasiparticles energy counted from the chemical potential, and $\omega = \pi T(2n+1)$ is the fermion Matsubara frequency. The Planck constant \hbar was everywhere put equal to 1. The self-energy part is given by the equation

$$\Sigma_{\mathbf{p}}(\omega) = \int \frac{d^3 p'}{(2\pi)^3} U_{\mathbf{p}-\mathbf{p}'}^2 G(\mathbf{p}', \omega). \quad (11)$$

Here, according to Abrikosov and Gor'kov¹², the "impurity line" $U_{\mathbf{p}-\mathbf{p}'}^2$ arises after averaging over impurity positions and also over the orientation of the spins of all impurity atoms, $\langle S_i S_k \rangle = \frac{1}{3} S(S+1) \delta_{ik}$, where in our particular case $S = 1/2$. Then taking into account $\sigma_{\alpha\gamma}^i \sigma_{\gamma\alpha}^i = \frac{3}{4}$ we obtain

$$\begin{aligned} U_{\mathbf{p}}^2 &= n_i [u_{\mathbf{p}}^2 + \langle S_i \sigma_{\alpha\gamma}^i S_k \sigma_{\gamma\alpha}^k \rangle J_{\mathbf{p}}^2] \\ &= n_i \left[u_{\mathbf{p}}^2 + \frac{1}{4} S(S+1) J_{\mathbf{p}}^2 \right], \end{aligned} \quad (12)$$

where n_i is impurity concentration and $u(\mathbf{p})$ and $J(\mathbf{p})$ are the Fourier transforms of the amplitudes of potential and exchange scattering from Eq.(8). According to assumption about global anisotropy they depend on the momentum direction such that

$$n_i u_{\mathbf{p}}^2 = \frac{1}{2\pi N_0 \tau_p} \left\{ 1 - \delta_p \left[\hat{p}_z^2 - \frac{1}{2} (\hat{p}_x^2 + \hat{p}_y^2) \right] \right\}, \quad (13)$$

$$n_i J_{\mathbf{p}}^2 = \frac{1}{2\pi N_0 \tau_{ex}} \left\{ 1 - \delta_{ex} \left[\hat{p}_z^2 - \frac{1}{2} (\hat{p}_x^2 + \hat{p}_y^2) \right] \right\}, \quad (14)$$

where N_0 is the density of states per one spin projection, \hat{p}_i are the projections of momentum unit vector $\frac{\mathbf{p}}{|\mathbf{p}|}$ on the $i = (x, y, z)$ coordinate axis, τ_p , and τ_{ex} are the isotropic parts of mean free time of potential and exchange scattering and δ_p and δ_{ex} are the corresponding degree of anisotropy. The anisotropic part of $U_{\mathbf{p}}^2$ is taken with the sign opposite to that in Ref.5 and chosen such that $\int \frac{d\Omega}{4\pi} U_{\mathbf{p}}^2$ is independent of the anisotropic part of scattering.

So, the self energy obtained from Eqs.(11)-(14) is

$$\Sigma_{\mathbf{p}}(\omega) = -\frac{i}{2\tau} \left\{ 1 - \delta \left[\hat{p}_z^2 - \frac{1}{2} (\hat{p}_x^2 + \hat{p}_y^2) \right] \right\} \text{sign } \omega. \quad (15)$$

Along with the isotropic term it includes a term describing the global uniaxial anisotropy. Each of these terms consists of two independent parts: the potential part and the exchange one determined in the following way

$$\frac{1}{\tau} = \frac{1}{\tau_p} + \frac{1}{\tau_{ex}}, \quad \frac{\delta}{\tau} = \frac{\delta_p}{\tau_p} + \frac{\delta_{ex}}{\tau_{ex}}. \quad (16)$$

The vertex part must be found from the integral equation

$$\Gamma_j(\mathbf{p}, \omega) = \hat{p}_j + n \int \frac{d^3 p'}{(2\pi)^3} \left[u_{\mathbf{p}-\mathbf{p}'}^2 + \frac{1}{3} S(S+1) (g^\dagger)_{\alpha\beta}^\mu \sigma_{\lambda\alpha}^i \sigma_{\rho\beta}^i g_{\lambda\rho}^\mu J_{\mathbf{p}-\mathbf{p}'}^2 \right] \Gamma_j(\mathbf{p}', \omega) G(\mathbf{p}', \omega) G(-\mathbf{p}', -\omega). \quad (17)$$

It is known¹² that for the case of singlet superconductivity the exchange part of scattering in this equation is given by

$$\frac{1}{3} S(S+1) g_{\alpha\beta}^t \sigma_{\lambda\alpha}^i \sigma_{\rho\beta}^i g_{\lambda\rho} J_{\mathbf{q}}^2 = -\frac{1}{4} S(S+1) J_{\mathbf{q}}^2, \quad (18)$$

where the matrix $\hat{g} = \begin{pmatrix} 0 & 1 \\ -1 & 0 \end{pmatrix}$, and the superscript t indicates transposition. As result, there are two different

"scattering time" originating from the self-energy and the vertex¹². The corresponding combination for the triplet pairing is

$$\frac{1}{3} S(S+1) (g^\dagger)_{\alpha\beta}^\mu \sigma_{\lambda\alpha}^i \sigma_{\rho\beta}^i g_{\lambda\rho}^\nu J_{\mathbf{q}}^2 = \frac{1}{4} S(S+1) J_{\mathbf{q}}^2 \delta_{\mu\nu}, \quad (19)$$

where $g_{\lambda\rho}^\nu = (-\sigma_{\lambda\rho}^z, i\delta_{\lambda\rho}, \sigma_{\lambda\rho}^x)$, such that the "scattering time" originating from the self-energy and the vertex are

equal to each other. Thus, the Eq.(17) is

$$\Gamma_j(\mathbf{p}, \omega) = \hat{p}_j + \int \frac{d^3 p'}{(2\pi)^3} U_{\mathbf{p}-\mathbf{p}'}^2 \Gamma_j(\mathbf{p}', \omega) G(\mathbf{p}', \omega) G(-\mathbf{p}', -\omega) \quad (20)$$

and its solution has the form

$$\Gamma_j^{\mu\nu}(\omega, \mathbf{p}) = \left\{ \hat{p}_j + \Gamma_\omega \left[\hat{p}_z \hat{z}_j - \frac{1}{2} (\hat{p}_x \hat{x}_j + \hat{p}_y \hat{y}_j) \right] \right\}, \quad (21)$$

where for the $\delta \ll 1$

$$\Gamma_\omega = \frac{\delta}{3\tau \left| \omega + \frac{1}{2\tau} \text{sign } \omega \right|} + \mathcal{O}(\delta^2). \quad (22)$$

Substitution of the vertex $\Gamma_j^{\mu\nu}(\omega, \mathbf{p})$ and the Green function $G(\mathbf{p}, \omega)$ into Eq.(9) yields

$$F_2 = \alpha A_{\mu i}^* A_{\mu i} - 2\eta \left[A_{\mu z}^* A_{\mu z} - \frac{1}{2} (A_{\mu x}^* A_{\mu x} + A_{\mu y}^* A_{\mu y}) \right] \quad (23)$$

where

$$\alpha = \frac{N_0}{3} \left[\ln \frac{T}{T_{co}} + \psi \left(\frac{1}{2} + \frac{1}{4\pi T\tau} \right) - \psi \left(\frac{1}{2} \right) - \frac{1}{5} \frac{\delta}{4\pi T\tau} \psi^{(1)} \left(\frac{1}{2} + \frac{1}{4\pi T\tau} \right) \right], \quad (24)$$

$$\eta = \frac{8N_0}{45} \frac{\delta}{4\pi T\tau} \psi^{(1)} \left(\frac{1}{2} + \frac{1}{4\pi T\tau} \right). \quad (25)$$

Here, $\psi(z)$, and $\psi^{(1)}(z)$ are the digamma function and its first derivative.

At $\delta > 0$ the critical temperature of the phase transition to the superfluid state is maximal for the polar phase Eq.(4), and is determined by the equation

$$\alpha - 2\eta = 0. \quad (26)$$

In the limit of weak scattering $4\pi T\tau_c \gg 1$ the transition to the polar state occurs at

$$T_{c1} = T_{co} - \frac{\pi}{8\tau} + \frac{11\pi}{60\tau} \delta. \quad (27)$$

It is worthwhile to recall that at small degrees of anisotropy the temperature interval of the polar state existence is small and hardly observable.

According to Eq.(16) the degree of global anisotropy δ is determined by two independent terms originating from the potential and the exchange scattering. The latter can in principle suppress the anisotropy. However, the quasi-particle self-energy and the vertex part due to exchange scattering have the same structure as for pure potential scattering. Hence, the anisotropy suppression can also originate from the potential scattering which is different for the covering of aerogel stands by a solid ^3He layer.

The change in anisotropy of scattering for different types of covering must also reveal itself in the changes of spin diffusion anisotropy. In the next section I derive the normal liquid ^3He spin diffusion current flowing through the media filled by the randomly distributed impurities with globally anisotropic potential and exchange scattering.

III. SPIN CURRENT

The spin current in neutral Fermi liquid is calculated^{13,14} as the response to the gradient of angle of rotation of the spin space $\boldsymbol{\omega}_i = \nabla_i \boldsymbol{\theta}$,

$$\mathbf{j}_i = -\frac{\delta H}{\delta \boldsymbol{\omega}_i}, \quad (28)$$

where

$$H = \frac{1}{2m} \int d^3 r (D_i^{\alpha\lambda} \psi_\lambda)^\dagger D_i^{\alpha\mu} \psi_\mu + H_{int}, \quad (29)$$

$$D_i^{\alpha\beta} = -i\delta_{\alpha\beta} \nabla_i + \frac{1}{2} \boldsymbol{\sigma}_{\alpha\beta} \boldsymbol{\omega}_i, \quad (30)$$

and H_{int} includes the Fermi liquid interaction and the interaction with impurities, Eq.(8).

At low temperatures the collisions between the Fermi liquid quasiparticles induce negligibly small correction to the spin diffusion due to the scattering on aerogel strands. On the other hand, we are mainly interested in the spin current anisotropy in the anisotropic media and will ignore the temperature dependence of exchange amplitude of scattering due to the Kondo effect¹⁵. In this case one can work with the field theory technique for $T=0$. The response of the gauge field $\boldsymbol{\omega}_i$ is calculated in complete analogy with response to the usual vector potential A_i in the calculation of electric current in an isotropic metal with randomly distributed impurities performed in¹⁶. The spin current at finite wave vector \mathbf{k} and external frequency ω is

$$\mathbf{j}_i(\mathbf{k}, \omega) = \frac{i}{4m} Tr \int_{-\infty}^{+\infty} \frac{d\varepsilon}{2\pi} \int \frac{d^3 p}{(2\pi)^3} p_i \boldsymbol{\sigma}_{\alpha\beta} (\boldsymbol{\sigma}_{\beta\alpha} \boldsymbol{\omega}_j) \Pi_j - \frac{1}{4} n \boldsymbol{\omega}_i, \quad (31)$$

where n is the number of liquid ^3He atoms in the unit volume, function Π_j is determined by the equation

$$\Pi_j(p, p-k) = G(\mathbf{p}, \varepsilon) G(\mathbf{p}-\mathbf{k}, \varepsilon-\omega) \times \left[p_j + \int \frac{d^3 p'}{(2\pi)^3} U^2(\mathbf{p}-\mathbf{p}') \Pi_j(p', p'-k) \right], \quad (32)$$

$p = (\mathbf{p}, \varepsilon)$, $k = (\mathbf{k}, \omega)$,

$$G(\mathbf{p}, \varepsilon) = \frac{1}{\varepsilon - \xi_{\mathbf{p}} - \Sigma_{\mathbf{p}}(\varepsilon)}, \quad (33)$$

$U^2(\mathbf{p})$ is determined by Eq.(12), and $\Sigma_{\mathbf{p}}(\varepsilon)$ is given by Eq.(15). The vertex correction does not introduce changes in the spin structure of Eq.(31) due to the identity $\sigma_{\alpha\lambda}^i \sigma_{\lambda\mu}^p \sigma_{\rho\alpha}^p \sigma_{\mu\rho}^j = \sigma_{\alpha\beta}^i \sigma_{\beta\alpha}^j$.

At $\mathbf{k} = 0$, $\omega = 0$ the first term in the current expression (31) cancels out the second "diamagnetic" term. We are interested in calculating the current at $\mathbf{k} = 0, \omega \neq 0$. In this case,

$$\mathbf{j}_i = \frac{i}{4m} Tr \int_0^\omega \frac{d\varepsilon}{2\pi} \int \frac{d^3 p}{(2\pi)^3} p_i \boldsymbol{\sigma}_{\alpha\beta} (\boldsymbol{\sigma}_{\beta\alpha} \boldsymbol{\omega}_j) \Pi_j(\mathbf{k} = 0) \quad (34)$$

and the solution of Eq.(32) in a linear approximation with respect to δ and at $\omega\tau \ll 1$ is

$$\begin{aligned} \Pi_j(\mathbf{k} = 0) &= G(\mathbf{p}, \varepsilon)G(\mathbf{p}, \varepsilon - \omega) \\ &\times \left\{ p_j + \frac{2}{3}\delta \left[\hat{p}_z \hat{z}_j - \frac{1}{2}(\hat{p}_x \hat{x}_j + \hat{p}_y \hat{y}_j) \right] \right\}. \end{aligned} \quad (35)$$

Substituting this into Eq.(34), we obtain

$$\mathbf{j}_i = \frac{1}{6} \left\{ \delta_{ij} + \frac{16}{15}\delta \left[\hat{z}_i \hat{z}_j - \frac{1}{2}(\hat{x}_i \hat{x}_j + \hat{y}_i \hat{y}_j) \right] \right\} i\omega\tau N_0 v_F^2 \boldsymbol{\omega}_j. \quad (36)$$

Here v_F is the Fermi velocity. Making use of the Larmor theorem

$$\boldsymbol{\gamma}\mathbf{H} = \frac{\partial\boldsymbol{\theta}}{\partial t} = -i\omega\boldsymbol{\theta}, \quad (37)$$

where $\boldsymbol{\gamma} = 2\boldsymbol{\mu}$ is the gyromagnetic ratio, and $\boldsymbol{\mu}$ is the magnetic moment of ${}^3\text{He}$ atoms, one can rewrite the expression for current as

$$\mathbf{j}_i = -\frac{1}{3} \left\{ \delta_{ij} + \frac{16}{15}\delta \left[\hat{z}_i \hat{z}_j - \frac{1}{2}(\hat{x}_i \hat{x}_j + \hat{y}_i \hat{y}_j) \right] \right\} \tau N_0 v_F^2 \boldsymbol{\mu} \nabla_j \mathbf{H}. \quad (38)$$

To rewrite the spin current as the magnetic diffusion current one should multiply both sides of this equation by $2\boldsymbol{\mu}$ to obtain

$$\mathbf{j}_i^M = -\frac{1}{3} \left\{ \delta_{ij} + \frac{16}{15}\delta \left[\hat{z}_i \hat{z}_j - \frac{1}{2}(\hat{x}_i \hat{x}_j + \hat{y}_i \hat{y}_j) \right] \right\} \tau v_F^2 \nabla_j \mathbf{M}, \quad (39)$$

where the Fermi-liquid magnetization is $\mathbf{M} = 2\mu^2 N_0 \mathbf{H}$. Thus, the spin diffusion currents along the direction of nafen strands and in the direction perpendicular to them are

$$\mathbf{j}_z^M = -\frac{1}{3} \left\{ 1 + \frac{16}{15}\delta \right\} \tau v_F^2 \nabla_z \mathbf{M}, \quad (40)$$

$$\mathbf{j}_x^M = -\frac{1}{3} \left\{ 1 - \frac{8}{15}\delta \right\} \tau v_F^2 \nabla_x \mathbf{M} \quad (41)$$

respectively. One can demonstrate that a similar calculation taking into account the Fermi liquid renormalization adds in these formulas the pre-factor $(1 + F_0^a)(1 + F_1^a/3)$.

Thus, the anisotropy of the spin diffusion coefficient is expressed through the same parameter of anisotropy δ that determines the temperature interval of the polar state existence.

IV. CONCLUSION

It was shown that the degree of global anisotropy responsible for polar state stability is determined by two mechanisms, originating from the potential and the exchange scattering. The suppression of anisotropy narrows the temperature interval of the polar state existence, making it hardly observable. The anisotropy decrease can be controlled by the measurements of spin diffusion

because the difference in the spin diffusion coefficients in directions parallel and perpendicular to nafen strands is found to be proportional to the same parameter that determines the polar state emergence.

The authors of the paper⁷ have pointed out the dominate role of the exchange scattering in the anisotropy suppression (see the citation of Ref.7 in the Introduction). However, according to the presented results, the exchange interaction yields the quasiparticle self-energy and the vertex part of the same structure as for pure potential scattering. Hence, although the exchange mechanism works only in the case of coating of aerogel strands by a solid ${}^3\text{He}$ layer it is possible that the main role in the anisotropy decrease is played by the change of potential scattering with aerogel strands covered by ${}^3\text{He}$ instead of ${}^4\text{He}$. The problem of choosing between the two mechanisms of anisotropy suppression will be addressed in future investigations.

Being mainly interested in the role of anisotropy of exchange scattering, I neglect throughout this paper the possible temperature dependence of the amplitude of exchange scattering due to the Kondo effect. The logarithmic increase of positive ion mobility starting at $T=50$ mK up to the superfluid transition temperature (see the paper¹⁵ and references therein) means that the exchange scattering has a ferromagnetic character, in agreement with the notion that ${}^3\text{He}$ is an almost ferromagnetic Fermi liquid. Thus, the Kondo effect weakens the mag-

nitude of the pair breaking by magnetic scattering.

Acknowledgments

I am indebted to V.Dmitriev for helpful and stimulating discussions.

-
- ¹ D.Vollhardt and P.Wölfle, *The Superfluid Phases of Helium-3* (Taylor and Francis, London,1990).
- ² J.V.Porto and J.M.Parpia, Phys. Rev. Lett. **74**, 4667 (1995).
- ³ D. T. Sprague, T. M. Haard, J.B. Kycia, M. R. Rand, Y. Lee, P. J. Hamot, and W. P. Halperin, Phys. Rev. Lett. **75**, 661 (1995).
- ⁴ B. I. Barker, Y. Lee, L. Polukhina, D. D. Osheroff, L.W. Hrubesh and J. F. Poco, Phys. Rev. Lett. **85**, 2148 (2000).
- ⁵ K. Aoyama and R. Ikeda, Phys. Rev. B **73**, 060504 (2006).
- ⁶ V. V. Dmitriev, A. A. Senin, A. A. Soldatov, and A. N. Yudin, Phys. Rev. Lett. **115**, 165304 (2015).
- ⁷ V. V. Dmitriev, A. A. Soldatov, and A. N. Yudin, Phys. Rev. Lett. **120**, 075301 (2018).
- ⁸ V. P. Mineev, Low Temp. Phys. **39**, 818 (2013).
- ⁹ V. P. Mineev, J. Low Temp.Phys. **177**, 48 (2014).
- ¹⁰ J.A.Sauls and P.Sharma, Phys.Rev. B **68**, 224502 (2003).
- ¹¹ G.A.Baramidze and G.A.Kharadze, preprint arXiv:cond-mat/0307612.
- ¹² A.A.Abrikosov and L.P.Gor'kov, Zh. Exp.Teor. Fiz. **39**, 1781 (1960) [Sov. Phys. JETP **12**, 1243 (1961)].
- ¹³ V.P.Mineev and Yu.G.Makhlin, Zh. Eksp.Teor. Fiz. **102**, 876 (1992) [Sov. Phys. JETP **75**, 479 (1992)].
- ¹⁴ V.P.Mineev and G.E.Volovik, J. Low Temp.Phys. **89**, 823 (1992).
- ¹⁵ V.M.Edelstein, Zh. Eksp. Teor. Fiz. bf 85, 543 (1983) [So. Phys. JETP **58**, 317 (1983)].
- ¹⁶ A.A.Abrikosov and L.P.Gor'kov, Zh. Exp.Teor. Fiz. **35**, 1558 (1958) Sov. Phys. JETP **8**, 1090 (1959).