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Solution of the problem of catastrophic relaxation of homogeneous spin precession in superfluid ^{3}He-B

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The quantitative analysis of the "catastrophic relaxation" of the coherent spin precession in ^{3}He-B is presented. This phenomenon has been observed below the temperature about 0.5 T_c as an abrupt shortening of the induction signal decay. It is explained in terms of the decay instability of homogeneous transverse NMR mode into spin waves of the longitudinal NMR. Recently the cross interaction amplitude between the two modes has been calculated by Sourovtsev and Fomin [9] for the so-called Brinkman-Smith configuration, i.e. for the orientation of the orbital momentum of Cooper pairs along the magnetic field, $\mathbf{L} \parallel \mathbf{H}$. In their treatment, the interaction is caused by the anisotropy of the speed of the spin waves. We found that in the more general case of the non-parallel orientation of $\mathbf{L}$ corresponding to the typical conditions of experiment, the spin-orbital interaction provides the additional interaction between the modes. By analyzing experimental data we are able to distinguish which contribution is dominating in different regimes.

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1. Introduction. For magnetically ordered systems the instability of homogeneous precession is a well known phenomenon. Suhl [1] explained it in terms of parametric instability of the mode of precession with respect to excitations of pairs of spin waves satisfying the condition of resonance:

$$n\omega_L = \omega_s(\mathbf{k}) + \omega_s(-\mathbf{k}), \qquad (1)$$

where ω_L is the precession frequency and n is integer (see also book [2]). Parametric instability (see, e.g. textbook by Landay and Lifshits, [3]) is a particular case of the decay instability, which is very general and well known in physics of nonlinear wave phenomenon, that includes induced light scattering (photon decaying into photon and phonon or spin wave); decay of Langmuir waves into Langmuir and ion-sound waves in non-isothermal plasmas; decay of capillary wave into capillary and gravity wave on the water surface; etc. In quantum solids and liquids the decay instability has been observed in anti-ferromagnetic solid ^{3}He [4], and has been predicted for superfluid liquid ^{3}He-A [5] where it has been observed later [6].

Unique feature of superfluid ^{3}He-B is the phenomenon of self sustained and long-lived precession with the coherent phase across the whole precessing domain, the so-called Homogeneously Precessing Domain (HPD) [7]. What is really amazing is a huge precession angle up to 104° in the presence of lower-frequency spin waves (longitudinal NMR), for which the resonance conditions (1) are definitely satisfied. In general nonlinear media the life time of this kind of excitation (with dimensionless amplitude of order unity) should be about their period. The main question is why the HPD is so long-living excitation? The explanation is unique symmetry of the leading Zeeman interaction which therefore does not contribute to interaction amplitude of the decay processes (1) [that we denote as $V(k)$]. Moreover, even subleading spin-orbit interaction in the Brinkman-Smith configuration (with $\mathbf{L} \parallel \mathbf{H}$) does not contribute to the interaction amplitude $V(k)$ due to the symmetry constraint.

Nevertheless the abrupt instability of HPD and of the homogeneous precession in general, called the catastrophic relaxation, has been observed in ^{3}He-B below $\sim (0.4 - 0.5)T_c$ [8]. The main physical question here is to clarify what is the origin of contribution to interaction amplitude $V(k)$ that is responsible for the catastrophic relaxation in general and in the experimental conditions [8] in particular. First reasonable step in this direction was recently made in [9] where the authors considered simple $\mathbf{L} \parallel \mathbf{H}$ configuration and found a

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non-vanishing contribution to $V(k)$, originated from the dependence of the spin-wave velocity on the direction of propagation. As we will show below this contribution hardly can be considered as the main one in typical experimental conditions [8].

In our Letter we found another contribution to the decay-interaction (1) amplitude $V(k)$ [denoted below as $V_{\rm BLV}(k)$] that allows one to rationalize main features of observed catastrophic relaxation. Our point is that under conditions of the experiment [8], the boundary conditions on the wall of container induce the texture of the order parameter in which the orbital vector **L** deviates from its symmetric orientation along the magnetic field **H** in the most of the container volume. The symmetry of the spin-orbit interaction is violated providing the additional term in the interaction $V(k)$ between the modes, which is dominating in typical experiments with the catastrophic relaxation [8, 10, 11].

2. General precessing states. The homogeneous precession of magnetization in ^{3}He-B has been analyzed by Brinkman and Smith [12] and in a great detail by Fomin [13] for the ideal case of the symmetric orientation of the orbital momentum $\mathbf{L} \parallel \mathbf{H}$. To discuss the real experiments in which the orbital vector **L** is deflected due to the boundary conditions and forms the texture, the computer simulations on the basis of the full set of Leggett-Takagi equations for spin dynamics have been employed, using the program elaborated by Golo [14] for the one dimensional texture. The instability of homogeneous precession was demonstrated: it was found that at the temperatures, corresponding to the experimentally observed catastrophic relaxation, some mode with fixed non-zero wave vector k starts to grow exponentially. It was also found [15] that the magnetic field dependence of the onset of instability is in quantitative agreement with the field dependence of catastrophic relaxation observed in [16]. This demonstrates that the catastrophic relaxation is indeed in the frame of the Leggett-Takagi equations, but for the real understanding of this phenomenon we must identify the principal mechanism of this calculated instability. For that we modified the theory [9] for the general case of arbitrary orientation of the orbital vector **L**.

2.1. *Symmetry of precessing states.* Let us consider the general homogeneous free precession in external magnetic field **H**. In liquid ^{3}He the spin-orbit (dipole-dipole) interaction is weak. If it is neglected, we can apply the powerful Larmor theorem, according to which, in the spin-space coordinate frame rotating with the Larmor frequency the effect of magnetic field on spins of the ^{3}He atoms is completely compensated. Since the magnetic field becomes irrelevant, the symmetry group of the physical laws in the precessing frame is

$$G = SO_3^L \times SO_3^S , \tag{2}$$

where SO_3^L is the group of orbital rotations in the laboratory frame; and SO_3^S is the group of spin rotations in the rotating frame whose elements $\tilde{\mathbf{g}}(\mathbf{t})$ are constructed from the elements **g** of conventional spin rotations in the laboratory frame:

$$\tilde{\mathbf{g}}(t) = \mathbf{O}^{-1}(\hat{\mathbf{z}}, \omega_L t)\ \mathbf{g}\ \mathbf{O}(\hat{\mathbf{z}}, \omega_L t)\,. \tag{3}$$

Here the matrix $O_{\alpha\beta}(\hat{\mathbf{z}}, \omega t)$ describes the transformation from the laboratory frame into the rotating frame – this is the rotation about the magnetic field axis $\hat{\mathbf{z}}$ by angle $\omega_L t$. Now we can find all the degenerate coherent states of the Larmor precession applying the symmetry group G to the simplest equilibrium state of the given superfluid phase: $\mathbf{A} = \mathbf{O}^{-1}\mathbf{R}^{(1)}\mathbf{O}\mathbf{A}^{(0)}(\mathbf{R}^{(2)})^{-1}$, where $\mathbf{R}^{(1)}$ is the arbitrary matrix describing spin rotations in the precessing frame and $\mathbf{R}^{(2)}$ is another arbitrary matrix which describes the orbital rotations in the laboratory frame. In case of ^{3}He-B this state corresponds to the state of Cooper pairs with $L = S = 1$ and the total angular momentum $J = 0$ [17]:

$$A^{(0)}_{\alpha i} = \Delta_B\ \delta_{\alpha i}\,. \tag{4}$$

The action of elements of the group G on this stationary state leads to the following general precession of ^{3}He-B with the Larmor frequency (if the spin-orbit interaction is neglected):

$$A_{\alpha i}(t) = \Delta_B R_{\alpha i}(t), \tag{5}$$

$$R_{\alpha i}(t) = O_{\alpha\beta}(\hat{\mathbf{z}}, -\omega t) R^{(1)}_{\beta\gamma} O_{\gamma\mu}(\hat{\mathbf{z}}, \omega t) (R^{(2)})^{-1}_{\mu i}\,. \tag{6}$$

The matrix $\mathbf{R}^{(1)}$ determines the direction of spin density in the precessing frame:

$$S_\alpha = \chi R^{(1)}_{\alpha\beta} H_\beta\,, \tag{7}$$

where χ is the spin susceptibility of ^{3}He-B. This corresponds to the precession of spin with the tipping angle $\cos\beta_1 = R^{(1)}_{zz}$. The matrix $\mathbf{R}^{(2)}$ determines the direction of orbital momentum density in the laboratory frame:

$$L_i = -R_{\alpha i}(t) S_\alpha(t) = -\chi R^{(2)}_{i\alpha} H_\alpha\,, \tag{8}$$

with the tipping angle $\cos\beta_2 = R^{(2)}_{zz}$.

2.2. *Spin-orbit interaction as perturbation.* The spin-orbit interaction in ^{3}He-B is

$$F_D = \frac{2\chi\Omega_L^2}{15}\left[\mathrm{Tr}\mathbf{R}(t) - \frac{1}{2}\right]^2 = \frac{8\chi\Omega_L^2}{15}\left[\cos\theta(t) + \frac{1}{4}\right]^2, \tag{9}$$

where Ω_L is the so called Leggett frequency – the frequency of the longitudinal NMR; θ is the angle of rotation in the parametrization of the matrix $R_{\alpha i}$ in terms of the angle and axis of rotation [17]; we use the system of units in which the gyromagnetic ratio γ for the ^{3}He atom is 1, hence the magnetic field and the frequency will have the same physical dimension.

In the general state of the Larmor precession (5), the spin-orbit interaction contains the time independent part and rapidly oscillating terms with frequencies ω_L, $2\omega_L$, $3\omega_L$ and $4\omega_L$:

$$F_D(\gamma) = F_0 + \sum_{n=1}^{4} F_n \cos(n\omega_L t)\,. \qquad (10)$$

The time-independent part – the average over fast oscillations – is

$$F_0 = \frac{2}{15}\chi\Omega_L^2[(sl - \frac{1}{2} + \frac{1}{2}(1+s)(1+l)\cos\gamma)^2 + \\ + \frac{1}{8}(1-s)^2(1-l)^2 + (1-s^2)(1-l^2)(1+\cos\gamma)]\,. \quad (11)$$

Here $s = \cos\beta_1$ and $l = \cos\beta_2$ are z projections of unit vectors $\hat{\mathbf{s}} = \mathbf{S}/S$ and $\hat{\mathbf{l}} = -\mathbf{L}/L$; and γ is another free parameter of the general precession. Altogether the free precession is characterized by 5 independent parameters coming from two matrices $\mathbf{R}^{(1)}$ and $\mathbf{R}^{(2)}$ [18]: two angles of spin $\mathbf{S}$, two angles of the orbital momentum $\mathbf{L}$, and the relative rotation of matrices by angle γ. In the case of the non-precessing magnetization, the γ-mode corresponds to the longitudinal NMR mode.

3. Parametric instability of HPD to radiation of γ-mode.

3.1. *Lagrangian for γ mode.* In the simplest description, the dynamics of the γ-mode is determined by the following Lagrangian:

$$\mathcal{L} = -\frac{1}{2}\chi\left(\dot\gamma^2 - c^2(\nabla\gamma)^2\right) + F_D(\gamma)\,. \qquad (12)$$

Here we used the approximation of a single speed of spin waves c, since the effect of the anisotropy of the spin wave velocity has already been discussed in Ref. [9]. In the time-dependent part of F_D we only consider the first harmonic, i.e. according to Eq.(1) we discuss the parametric excitation of two γ-modes with $ck \approx \omega_L/2$. The amplitude of the first harmonic is:

$$F_1 = \frac{4}{15}\chi\Omega_L^2 \sin\beta_1 \sin\beta_2 \cos(\gamma/2) \times \\ \times \left(2sl - 1 + \frac{(1-s)(1-l)}{2} + (1+s)(1+l)\cos\gamma\right)\,. \quad (13)$$

Further we assume that the system is in the minimum of the dipole energy F_0 as a function of γ. The equilibrium value $\gamma = \gamma_0$ is

$$\cos\gamma_0 = -\frac{(2sl-1) + 2(1-s)(1-l)}{(1+s)(1+l)}\,, \qquad (14)$$

which is valid if the right hand side of Eq. (14) does not exceed unity, i.e. when $s + l - 5sl < 2$.

For the discussion of Suhl instability we need the time-dependent term which is quadratic in $\gamma - \gamma_0$. Then the Lagrangian (12) which describes the parametric instability towards decay of Larmor precession to two γ-modes with $kc \approx \omega_L/2$ is (after the shift $\gamma - \gamma_0 \to \gamma$; neglecting Ω_L compared to ω_L; and neglecting the anisotropy of the spin-wave velocity):

$$\mathcal{L} = \frac{1}{2}\chi\left(-\dot\gamma^2 + c^2(\nabla\gamma)^2 + a\Omega_L^2\gamma^2\cos\omega_L t\right)\,, \qquad (15)$$

where, if $s + l - 5sl < 2$, the parameter a is

$$a = \frac{4}{15}\sin\beta_1\sin\beta_2\left[\frac{3(s+l-sl)}{2(1+s)(1+l)}\right]^{1/2} \times \\ \times \left[(1+s)(1+l) + 2(2sl-1) + \frac{35}{8}(1-s)(1-l)\right]\,. \quad (16)$$

3.2. *Parametric instability.* Let us rewrite the Lagrangian (15) in terms of Hamiltonian as function of creation and annihilation operators $b_{\mathbf{k}}$ and $b^*_{\mathbf{k}}$:

$$\gamma_{\mathbf{k}} = \frac{i\,(b_{\mathbf{k}} - b^*_{\mathbf{k}})}{\sqrt{2\chi\omega_s(k)}}, \quad \omega_s^2(k) = c^2k^2, \qquad (17)$$

$$p_{\mathbf{k}} = \chi\dot\gamma_{\mathbf{k}} = \sqrt{\chi\omega_s(k)/2}(b_{\mathbf{k}} + b^*_{\mathbf{k}})\,, \qquad (18)$$

$$\mathcal{H} = \sum_{\mathbf{k}}\omega_s(k)b^*_{\mathbf{k}}b_{\mathbf{k}} + \frac{1}{2}\sum_{\mathbf{k}}V(k)\left(e^{-i\omega_L t}b_{\mathbf{k}}b_{-\mathbf{k}} + \text{c.c.}\right)\,, \qquad (19)$$

$$V(k) = V_{\text{BLV}}(k) \equiv a\Omega_L^2\Big/\omega_s(k)\,. \qquad (20)$$

Here we neglected Ω_L compared to ω_L. The spectrum of the excited mode is $b_{\mathbf{k}}(t) = \tilde b_{\mathbf{k}}\exp(-i\omega_L/2t + \nu_{\mathbf{k}}t)$, where the instability increment $\nu_{\mathbf{k}}$ is

$$\nu_{\mathbf{k}} = \sqrt{V(k)^2 - [\omega_s(k) - \omega_L/2]^2}\,. \qquad (21)$$

Increment $\nu_{\mathbf{k}}$ reaches its maximum $\max\nu_k = V(k)$ for spin waves, satisfying condition (1) of the parametric resonance, $\omega_s(k') = \omega_L/2$. These modes grow exponentially:

$$b_{\mathbf{k}}(t) \propto \exp(V(k')t)\,. \qquad (22)$$

At finite temperatures this growing is damped by dissipation, but at low temperature the dissipation becomes small and catastrophic relaxation occurs. We shall follow Ref. [9] and assume the spin diffusion mechanism of dissipation. In this case the equation for temperature $T_{\rm cat}$ below which the instability of the homogeneous precession towards radiation of spin waves with $\omega_s(k') = ck' = \omega_L/2$ starts to develop is:

$$D(T_{\rm cat}) = 2V(k')c^2/\omega_L^2\,. \qquad (23)$$

Here $D(T)$ is the spin diffusion coefficient, which depends on temperature and decreases with decreasing T.

In our case

$$V_{\rm BLV}(k') = 2a\Omega_L^2/\omega_L\,. \qquad (24)$$

4. Two mechanism of Suhl instability. The characteristic feature of spin-orbit interaction in the isotropic ^{3}He-B is that it is symmetric with respect to the interchange of spin and orbital momenta [18]. This is the reason why Eqs. (11), (13) and (16) are symmetric under transformation $\beta_1 \leftrightarrow \beta_2$ (or $s \leftrightarrow l$). The immediate consequence of this symmetry is that the Brinkman-Smith (BS) precession ($\beta_1 = \beta$, $\beta_2 = 0$) is equivalent to the static state of ^{3}He-B in magnetic field ($\beta_1 = 0$, $\beta_2 = \beta$). Since the static state is stable, the BS precession cannot be destabilized by spin-orbit interaction, which is also seen from Eq.(16): the interaction amplitude $V(k) = 0$ if $\beta_2 = 0$ ($l = 1$).

This Z_2 symmetry of the Larmor precession in ^{3}He-B is violated by the term in the gradient energy which is responsible for the dependence of the spin-wave velocity on the direction of propagation. This term, which we omitted in our consideration, becomes time dependent even in the background of the BS mode. This leads to the Sourovtsev-Fomin (SF) contribution $V_{\rm SF}(k)$ to the interaction amplitude $V(k)$, discussed by [9] for $\mathbf{L}||\mathbf{H}$. SF mechanism can be easily extended to the general precession with arbitrary β_2. For that one must take into account, that the axis of anisotropy of spin-wave propagation is determined by the orbital vector $\mathbf{L}$. This gives the following modification of $V_{\rm SF}(k')$:

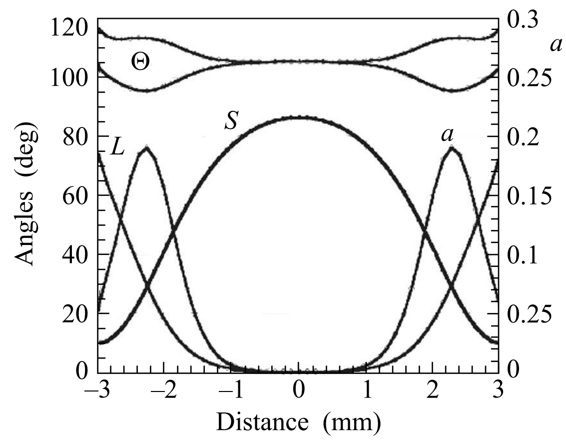
$$V_{\rm SF}(k') \Rightarrow \widetilde{V}_{\rm SF}(k') = \frac{\mu\,\omega_L}{4}\frac{\sin\beta_1|1-2s|}{2s^2-2s+5}|\sin\,2\delta|\,, \quad (25)$$

$$\cos\delta \equiv \hat{\mathbf{k}}\cdot\hat{\mathbf{l}}\,. \qquad (26)$$

Here $\mu = 1 - c_\perp^2/c_\parallel^2$ is the anisotropy of the spin-wave velocity, where $c_\parallel$ ($c_\perp$) is the velocity of spin waves propagating along (perpendicular to) $\hat{\mathbf{l}}$. We assume that both Ω_L^2/ω_L^2 and the spin-wave anisotropy μ are small parameters. In general

$$V(k) = V_{\rm BLV}(k) + \widetilde{V}_{\rm SF}(k)\,. \qquad (27)$$

Comparing these two contributions to $V(k)$, causing the catastrophic relaxation – due to spin-orbit interaction, $V(k) = V_{\rm BLV}(k)$, Eq. (20) and due to spin-wave anisotropy in Eq. (25) – one can see that $\widetilde{V}_{\rm SF}(k') > V_{\rm BLV}(k')$ at high fields, when the Larmor frequency ω_L essentially exceeds the Leggett frequency Ω_L. However, in the typical experiments [10, 11] the spin-orbit amplitude, $V_{\rm BLV}(k)$, appears to be more important. There are two reasons for that: the field is not sufficiently high ($\Omega_L/\omega_L \sim 0.3 - 0.5$); the texture of the orbital $\hat{\mathbf{l}}$ is typically formed due to boundary conditions at the walls of container: $\hat{\mathbf{l}}$ must be oriented along the normal to the wall. The 1D simulation [14] of the texture of the $\hat{\mathbf{s}}$ and $\hat{\mathbf{l}}$ fields (angles β_1 and β_2) under conditions of experiments but in the parallel-plate geometry is shown in Figure. The main part of the deviation from the BS mode develops in about 1 mm size region near the walls, which means that in the cylindrical cell the texture should be very similar. Thus the real mode of precession in the finite vessel is very different from the ideal BS configuration, which is static in the precessing frame. On the contrary, the angle θ of the order parameter $R_{\alpha i}$ is oscillating as shown in the upper part of the Figure. These oscillations lead to parametric instability at low temperature.



The texture in the precessing state in the cell between two parallel walls at $x = 0$ and $x = 6$ mm. Left scale shows angles β_1, β_2 and θ, while the right scale shows the local value of the parameter a in Eq.(16)

Figure also shows the parameter a in Eq. (16) as function of the coordinate x. Since the diffusion damping of spin waves occurs in the whole volume of the cell, the increment $V(k')$ in Eq.(24) must be averaged over the cylindrical cell: $V(k') = 2\bar{a}\Omega_L^2/\omega_L$, with $\bar{a}$= 0.099 for our texture. This leads to the critical value of spin diffusion at which the catastrophic relaxation must occur

$$D(T_{\rm cat}) = 4\bar{a}\frac{\Omega_L^2 c_\|^2}{\omega_L^3} = 0.085\ {\rm cm}^2/{\rm s}\,. \tag{28}$$

For Ω_L = 244 kHz and ω_L = 460 kHz this value corresponds to the temperature of the catastrophic relaxation $T_{\rm cat} = 0.5T_c$, according to [19], in a good agreement with the experimental value $T_{\rm cat} = 0.47T_c$ ([10, 11]).

In Ref. [9] the critical diffusion for optimal configuration for spin wave velocity anisotropy was estimated as about 0.03 cm^2/c. Because the texture of the $\hat{\mathbf{l}}$-vector influences $\widetilde{V}_{\rm SF}(k')$ in Eq.(25), the increment must be even smaller. This indicates that the spin-orbit mechanism of catastrophic relaxation is dominating under the conditions of the experiments [10, 11], $V_{\rm BLV}(k') > \widetilde{V}_{\rm SF}(k')$.

To verify that it is really so, let us analyze the Grenoble experiment [15], in which the texture was destroyed by the RF pulse so that **L** was parallel to **H** even in the surface layer. Under these conditions the catastrophic relaxation was observed at lower temperature of about 0.4 T$_c$. In this configuration one has $\beta_2 = 0$; the spin-orbit term, $V_{\rm BLV}(k)$, is switched off; the spin-wave anisotropy of Ref. [9], leading to $V_{\rm SF}(k)$, becomes the main source of instability. And indeed this $T_{\rm cat}$ well corresponds to the critical spin diffusion $D(T_{\rm cat})$ about 0.03 cm$^2/c$, calculated in Ref. [9]. When the texture is restored, then if the spin-wave anisotropy is the only mechanism of the catastrophic relaxation, the texture should lead to decrease of $V(k)$ and thus to decreasing $T_{\rm cat}$. Instead, $T_{\rm cat}$ increases demonstrating that when the texture appears another mechanism of instability emerges as we discussed here.

Two contributions in Eq. (27) to the increment of parametric instability (21) can be also compared using results of experiments made in Cornell [16]. In these experiments $T_{\rm cat}$ was measured for 31 bar at different ω_L. They found that $T_{\rm cat}$ decreases from 0.39 T$_c$ at ω_L = 600 kHz to 0.24 T$_c$ at ω_L = 3 MHz. This is well described by the spin-orbit mechanism of Suhl instability, which becomes weaker at higher field according to Eq. (22). On the contrary, if the spin-wave anisotropy is the dominating mechanism, $T_{\rm cat}$ would not depend on ω_L (there are processes in the spin-wave anisotropy mechanism whose increment is proportional to Ω_L^2/ω_L^2 [9] as in the case of the spin-orbit mechanism, but they are relatively small, since contain the product of two small parameters, μ and Ω_L^2/ω_L^2).

5. Conclusion. We calculated in a unified way two contributions $V_{\rm BLV}(k)$ and $\widetilde{V}_{\rm SF}(k)$ to the increment of the parametric instability of the Larmor precession in superfluid ^{3}He-B for any angle between **L** and **H**. The term $V_{\rm SF}(k)$, previously discussed in Ref. [9] for simple geometry with **L**||**H**, originates from the anisotropy of spin-wave velocity. We modified this contribution, $V_{\rm SF}(k) \to \widetilde{V}_{\rm SF}(k)$ for the case of the texture of the order parameter, which occurs in real experiments. In addition we found another term, $V_{\rm BLV}(k)$, which originates from the spin-orbit interaction. This second mechanism only takes place when the orbital momentum **L** deviates from its symmetric orientation along the magnetic field, in particular in the presence of texture. Our analytical result for the onset of the parametric instability due to this mechanism is in a good quantitative agreement with experimental results. In particular it gives the correct dependence of the catastrophic relaxation temperature on magnetic field. The spin-wave anisotropy mechanism must dominate either when the texture is destroyed or in the limit of strong magnetic field.

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