Phase transition of superfluid ^{3}He-B to an inhomogeneous state induced by an electric field

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The combined effect of electric and magnetic fields on the orientation of the anisotropy axis of superfluid B-phase liquid ^{3}He located in a gap between two parallel plates is considered. The configuration is analyzed for the case in which **E** and **H** are parallel and directed perpendicularly to the plane surface. It is shown that the superfluid B-phase undergoes a phase transition to an inhomogeneous state at a certain critical electric field intensity E_c. The E_c depends on the gap width, the magnetic field intensity, and the temperature.

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In the investigation of the properties of the superfluid state of liquid ^{3}He, a significant position is occupied by the problem of the orienting action exerted on the anisotropy axes of the A and B phases by the external fields by the superflow, and by the solid walls (see, for example, the review of Leggett[1]). Under the competing influences of these factors, the equilibrium state of the superconducting phases turns out generally to be inhomogeneous. In the present work, we consider the combined effect of electric and magnetic fields on the orientation properties of the superfluid B phase of liquid ^{3}He contained in a gap between two parallel plates. The case is analyzed in which the intensities **E** and **H** are mutually parallel and directed normal to the plane of the plates.

The superfluid state of a Fermi liquid with triplet Cooper pairing is described by a vector order parameter $\mathbf{\Delta}=\mathbf{\Delta}(T)\mathbf{d}(\mathbf{n})$, and the different phases differ from one another by the character of the dependence of the orientation **d** on the position on the surface of the Fermi sphere (**n** is the normal to this surface). For the B phase, in the absence of external fields, and far away from the walls of the container, $\mathbf{d}(\mathbf{n})=\hat{R}(\nu,\theta)\mathbf{n}$, where $\hat{R}(\nu,\theta)$ denotes the operation of rotation about the ν axis by an angle θ. The directrix ν defines the anisotropy axis of the B phase of liquid ^{3}He.

In the formation of the properties of the superfluid phases of ^{3}He, an important role is played by the dipole-dipole forces acting between the magnetic moments of the ^{3}He nuclei and the electric moments of the ^{3}He atoms ($-\alpha\mathbf{E}$), induced by the external field **E**. This is due to the spontaneous violation of the spin-orbit symmetry. The corresponding contributions to the free energy of the superfluid phases are equal to[1,2]

$$\mathscr{F}_{MD}=g_D\int\{3|\mathbf{n}\mathbf{d}(\mathbf{n})|^2-|\mathbf{d}(\mathbf{n})|^2\}\frac{d\mathbf{n}}{4\pi},\tag{1}$$

$$\mathscr{F}_{ED}=-\varepsilon^2 g_D\int\{3|\mathbf{n}\mathbf{e}|^2|\mathbf{d}(\mathbf{n})|^2-|\mathbf{d}(\mathbf{n})|^2\}\frac{d\mathbf{n}}{4\pi},\tag{2}$$

where the constant g_D characterizes the intensity of the magnetic dipole interaction, $\varepsilon=2\alpha E/\mu$, and $\mathbf{e}=\mathbf{E}/E$ (μ is the magnetic moment of the ^{3}He nucleus). The magnetic-dipole forces lead to a fixing of the equilibrium value of the angle $\theta=\theta_L=\arccos(-\frac{1}{4})$, and in the absence of an external magnetic field the orientation of the directrix ν is nevertheless arbitrary.

When the magnetic field is switched on, as a result of the predominant breaking of the Cooper pairs, the superfluid B phase acquires anisotropic properties, characterized by the anisotropy parameter $\delta(H)=(\Delta_\perp-\Delta_\parallel)/\Delta_\perp\sim H^2$. In this situation, the B phase should be described by the vector $\mathbf{d}(\mathbf{n})$ with components

$$d_\mu(\mathbf{n})\sim(R_{\mu i}-\delta(H)\delta_{\mu z}R_{zi})n_i,\quad \hat{z}\parallel\mathbf{H}.\tag{3}$$

Substituting (3) in (1) and (2), we easily obtain[3] an expression for that part of the dipole-dipole contribution to the B-phase free energy which depends on the angle θ between the directrix ν and the magnetic field **H** (we shall analyze below the case with $\mathbf{E}\parallel\mathbf{H}$):

$$\Delta\mathscr{F}_D(\vartheta)=g_D\delta(H)U(\vartheta)\approx aH^2U(\vartheta),\tag{4}$$

the dimensionless "potential" being

$$U(\vartheta)=-[\cos^2\vartheta+{}^1/_4\varepsilon^2(2\cos^2\vartheta-5\cos^4\vartheta)].\tag{5}$$

Of course, our analysis assumes that the effect of the external magnetic field on the wave function of the condensate predominates over the effect of the dipole-dipole forces, which are taken into account in the free energy only in the lowest order in the corresponding coupling constant ($g_zH^2\gg g_D,\ \varepsilon g_D$).

As indicated by us in[3], at $\varepsilon<1/\sqrt{2}$ the angle $\vartheta=0$ corresponds to a minimum of the free energy, i.e., at equilibrium we have $\nu\parallel\mathbf{H}$, just as in the absence of an electric field. In the region of sufficiently large electric fields, when $\varepsilon>1/\sqrt{2}$, the equilibrium orientation of the directrix ν should differ from the direction of **H** by an angle $\vartheta_0(E)$, where

$$\sin^2\vartheta_0=\frac{4}{5}\left(1-\frac{1}{2\varepsilon^2}\right)=\frac{4}{5}\left(1-\frac{E_{c0}^2}{E^2}\right),\quad E>E_{c0}=\frac{\mu}{2^{3/2}\alpha}.\tag{6}$$

On going through the critical value of the electric field E_c, a phase transition takes place from the state with $\vartheta_0=0$ to the state with $\vartheta_0\neq0$.

These considerations apply to an "open" geometry, when the volume effects due to the action of the external

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fields dominates over the influence of the walls. However, in real situations, one frequently deals with bounded volumes of the superfluid such that the presence of solid walls becomes important. This in turn applies to the B-phase orientation effects considered in the present work. We considered below the behavior of the directrix ν for a superfluid contained in the gap of width $2l$ between two plane-parallel walls and subjected to the combined action of electric and magnetic fields ($\mathbf{E}\parallel\mathbf{H}\parallel\mathbf{s}$, where $\mathbf{s}$ is the normal to the surface of the walls).

Since the directrix ν is assumed in the present situation, near the solid wall, to be oriented along $\mathbf{s}$, it is clear that the inclination of ν away from H under the action of the external electric field is made difficult. This should manifest itself in an increase in the critical value of the electric field ($E_c > E_{c0}$), to a greater extent the narrower the gap between the walls. Along with this, one should expect the phase with $\vartheta \neq 0$ to be inhomogeneous at $E > E_c$. These intuitive considerations are confirmed by the direct calculation given below.

For a complete analysis of the given problem, we must start out from the expression for the free energy $\mathscr{F}$ of the superfluid B phase, with account taken of the terms that describe the spatial inhomogeneity and the boundary effects.[4] In the description of the surface part of $\mathscr{F}$, we keep only the term corresponding to the effect of the magnetic field (we recall that $\mathbf{H}\parallel\mathbf{s}$):

$$\mathscr{F}_H^{(s)} = -d(\mathbf{s}\hat{R}\mathbf{H})^2 = \tilde{d}H^2[(\mathbf{s}\nu)^2 - {}^5/_2(\mathbf{s}\nu)^4]. \tag{7}$$

Directing the z axis perpendicular to the surface of the wall (with the origin of the coordinates at the center of the gap), describing the local orientation of the directrix ν with the help of the polar and azimuthal angles (ϑ and φ), and introducing the characteristic values of the magnetic field H_c and of the length z_c according to the formulas

$$H_c = \left(\frac{16ac}{13\tilde{d}^2}\right)^{1/2}, \quad z_c = \frac{1}{H}\left(\frac{16c}{13a}\right)^{1/2}, \tag{8}$$

(c is the rigidity coefficient, which enters into the gradient part of the free energy[4]), we can write the expression for the free energy in the dimensionless form (cf. Ref. 5):

$$\tilde{\mathscr{F}} = \frac{\mathscr{F}}{\tilde{d}HH_c} = \int_{-\tau}^{\tau} f_B(\vartheta, \dot{\vartheta}, \dot{\varphi})\,dt + f_s(\vartheta)\,|_{\tau+(-\tau)}, \tag{9}$$

where

$$f_B(\vartheta, \dot{\vartheta}, \dot{\varphi}) = -[\cos^2\vartheta + {}^1/_4\varepsilon^2(2\cos^2\vartheta - 5\cos^4\vartheta)]$$
$$+ (1 - {}^3/_{16}\sin^2\vartheta)\dot{\vartheta}^2 - {}^1/_8\sqrt{15}\dot{\vartheta}\dot{\varphi}\sin^3\vartheta + (1 - {}^5/_{16}\sin^2\vartheta)\dot{\varphi}^2\sin^2\vartheta, \tag{10}$$

$$f_s(\vartheta) = h(\cos^2\vartheta - {}^5/_2\cos^4\vartheta), \tag{11}$$

and $\dot{\vartheta} = d\vartheta/dt$, $t = z/z_c$, $\tau = l/z$, and $h = H/H_c$ (as has already been noted, we assume the magnetic field to be strong and do not take into account dipole-dipole effects in the surface part of the free energy).

To determine the equilibrium configuration $\nu = \nu(z)$, we must minimize the functional (9). The angle φ is a cyclic coordinate and it is not difficult to establish the fact that, at equilibrium,

$$\dot{\varphi} = \frac{\sqrt{15}}{16}\frac{\sin\vartheta}{1 - {}^5/_{16}\sin^2\vartheta}\dot{\vartheta}. \tag{12}$$

To find the equilibrium value $\vartheta = \vartheta(z)$, we shall assume that the deviation from the configuration with $\vartheta = 0$ is small and keep only terms of order ϑ^2 and ϑ^4. Then the problem reduces to minimization of the functional

$$\Phi = \frac{1}{2\tau}\int_{-\tau}^{\tau}\left\{\left(1 - \frac{3}{16}\vartheta^2\right)\dot{\vartheta}^2 - \varkappa^2\vartheta^2 + \left(\frac{5}{8} + \frac{23\varkappa^2}{24}\right)\vartheta^4\right\}dt$$
$$+ \frac{2h}{\tau}\left(\vartheta^2 - \frac{23}{24}\vartheta^4\right)\Big|_{\tau+(-\tau)}, \tag{13}$$

where the parameter $\varkappa^2 = 2\varepsilon^2 - 1$. It is clear that the role of the effects at the walls is determined by the ratio $h/\tau = l_c/l$; in this case, the new characteristic length is $l_c = \tilde{d}/a$.

Near the transition to the state with $\vartheta \neq 0$, the solution of the variational problem must be sought in the form $\vartheta = \vartheta_0 \cos\varkappa t$. Substituting this expression in (13) and determining the critical value $\varkappa_c$ with the help of the equation

$$\varkappa_c^2 \operatorname{tg}\varkappa_c\tau/\varkappa_c\tau = 4h/\tau, \tag{14}$$

we find that in the region $\varkappa \approx \varkappa_c$

$$\Phi(\vartheta_0) = -A(\varkappa_c\tau)\varkappa_c(\varkappa - \varkappa_c)\vartheta_0^2 + {}^1/_2 B(\varkappa_c\tau)\vartheta_0^4. \tag{15}$$

Here

$$A(x) = 1 + \sin 2x/2x, \tag{16}$$

$$B(x) = \frac{5}{4}\left[\eta(x) + \frac{23\varkappa_c^2}{15}\left(\eta(x) - \frac{\operatorname{tg}x}{x}\cos^4 x\right) - \frac{3\varkappa_c^2}{80}\left(1 - \frac{\sin 4x}{4x}\right)\right], \tag{17}$$

where

$$\eta(x) = \frac{3}{8} + \frac{1}{2}\frac{\sin 2x}{2x} + \frac{1}{8}\frac{\sin 4x}{4x}. \tag{18}$$

It is now clear that at $\varkappa < \varkappa_c$, i.e., at $E < E_c = E_{c0}(1 + \varkappa_c^2)^{1/2}$, the minimum free energy corresponds to $\vartheta_0 = 0$. Consequently, in this situation, the state of the superfluid B phase is homogeneous ($\vartheta \equiv 0$) and at equilibrium the directrix ν is parallel to H, as in the absence of the electric field. On the other hand, in the case $E > E_c$ ($\varkappa > \varkappa_c$) the equilibrium state turns out to be inhomogeneous ($\vartheta(z) = \vartheta_0 \cos\varkappa t$), and

$$\vartheta_0^2 \approx \frac{A(\varkappa_c\tau)}{B(\varkappa_c\tau)}\varkappa_c(\varkappa - \varkappa_c). \tag{19}$$

Thus, we come to the conclusion that at $E = E_c$ there is a phase transition to a inhomogeneous state with $\vartheta \neq 0$. As was to be expected, the critical value of the electric field is $E_c > E_{c0}$ by virtue of the retarding effect of the walls. Only in the limit $l \gg l_c$ do we have $\varkappa_c \ll 1$ and surface effects become insignificant. Estimate shows that at low temperatures ($T \ll T_c$) $l_c \approx 1$ mm and decreases as the temperature T_c is approached. As to the characteristic value of the magnetic field H_c, it depends weakly on the temperature and amounts to several hundred gausses.

The appearance of a state with $\vartheta \neq 0$ should be immediately reflected in the NMR spectrum of the superfluid ^{3}He-B. When the directrix ν is inclined to the magnetic field, one should expect the appearance of a shift in the frequency of the transverse resonance, accompanied by a corresponding decrease in the frequency of longitudinal NMR.

It should be emphasized that the formation of an inhomogeneous (textured) state of the B phase of ^{3}He with controllable spatial variation of the directrix $\nu(\mathbf{r})$ should be very useful in the experimental investigation of the spin-wave excitations in a superfluid Fermi liquid. As was recently demonstrated by Smith *et al.*,[6] the effect of texture on NMR in ^{3}He-B leads to the excitation of standing spin waves, the observation of which allows us, in particular, to measure the p component of the exchange Fermi-liquid Landau parameter F_1^a. The interaction F_1^a renormalizes the spectrum of longitudinal and transverse spin waves that propagate in the homogeneous B phase of ^{3}He. This question is briefly discussed in the Appendix.

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APPENDIX

In our preceding papers[7,8] we carried out a theoretical analysis of spin-wave excitations in the superfluid B-phase of liquid ^{3}He on the basis of a calculation scheme suggested by Larkin and Migdal.[9] As a result, the linear response of the spin system to a weak inhomogeneous variable magnetic field was calculated and the spectrum of spin waves with longitudinal and transverse polarizations was determined with account taken of the s- and p-components of the Landau-Fermi-liquid exchange interaction. The dynamic magnetic susceptibility tensor that we obtained for ^{3}He-B at $T=0$ is of the following explicit form:

$$\chi_{ij}(\mathbf{q},\omega)=\chi_{BW}\left[\frac{\omega_\parallel^2(\mathbf{q})}{\omega_\parallel^2(\mathbf{q})-\omega^2}e_ie_j+\frac{\omega_\perp^2(\mathbf{q})}{\omega_\perp^2(\mathbf{q})-\omega^2}(\delta_{ij}-e_ie_j)\right], \qquad (A.1)$$

where χ_{BW} denotes Balian-Werthamer (BW) static homogeneous susceptibility of the superfluid state, and the squares of the eigenfrequencies of the longitudinal and transverse branches are given by the formulas

$$\omega_\parallel^2(\mathbf{q})={}^1/_5 v_F^2\mathbf{q}^2(1+{}^2/_3F_0^a)(1+{}^1/_3F_1^a)(1+{}^2/_{15}F_1^a)^{-1}, \qquad (A.2a)$$

$$\omega_\perp^2(\mathbf{q})={}^2/_5 v_F^2\mathbf{q}^2(1+{}^2/_3F_0^a)(1+{}^1/_3F_1^a)(1+{}^1/_{12}F_1^a)(1+{}^2/_{15}F_1^a)^{-1}. \qquad (A.2b)$$

In the expression (A.1), e_1 denotes the component of the unit vector $\mathbf{e}=\hat{R}\mathbf{q}/|\mathbf{q}|$, where $\hat{R}$ is the rotation operator which describes the BW state. We note that our analysis pertains to the macroscopic limit and is valid in the region $(\omega, |\mathbf{q}|v_F)\ll\Delta$.

Introducing the notation

$$X=-\omega[\omega^2-\omega_\parallel^2(\mathbf{q})]^{-1}, \quad \mathscr{R}=-\omega[\omega^2-\omega_\perp^2(\mathbf{q})]^{-1}, \qquad (A.3)$$

we can put the expression (A.1) for χ_{ij} in the following form:

$$\chi_{ij}(\mathbf{q},\omega)=\chi_{BW}[\delta_{ij}+\omega Xe_ie_j+\omega\mathscr{R}(\delta_{ij}-e_ie_j)], \qquad (A.4)$$

which is exactly identical with the basic result of the recently published work of Chervonko.[10] Thus Chervonko duplicated the result obtained previously by us, but took into account F_2^a and F_3^a, in addition to the first two Landau exchange parameters F_0^a and F_1^a. In this connection, it is surprising that Chervonko found it possible to declare our results to be in error. The reason for this is evidently an incorrect understanding of the calculational procedure used by us.

In the macroscopic limit, the equation for the vector vertex $\boldsymbol{\Gamma}(\mathbf{n})$ has the form (details can be found in our previous paper[8]):

$$\Gamma(\mathbf{n})+\int F^a(\mathbf{nn}')\left\{\frac{1+\mathscr{P}}{2}\Gamma(\mathbf{n}')-\mathscr{P}(\mathbf{d}(\mathbf{n}')\Gamma(\mathbf{n}'))\mathbf{d}(\mathbf{n}')\right.$$
$$\left.-\frac{\omega+(\mathbf{qn}')v_F}{2\Delta}[\mathbf{d}(\mathbf{n}')\times\tau(\mathbf{n}')]\right\}\frac{d\mathbf{n}'}{4\pi}=\Gamma^\omega. \qquad (A.5)$$

In this equation we have the "anomalous" vertex $\boldsymbol{\tau}(\mathbf{n})$, which appears in the form of the combination $\boldsymbol{\gamma}(\mathbf{n})=[\mathbf{d}(\mathbf{n})\times\boldsymbol{\tau}(\mathbf{n})]$. It also enters in this form in the expression for the linear response of the spin system to an external magnetic excitation (see Eq. (20) of Ref. 8). The equation for $\boldsymbol{\gamma}(\mathbf{n})$ is

$$\int\left\{\frac{\omega^2-(\mathbf{qn})^2v_F^2}{4\Delta^2}[\mathbf{d}(\mathbf{n})\boldsymbol{\tau}(\mathbf{n})]+\frac{\omega+(\mathbf{qn})v_F}{2\Delta}[\mathbf{d}(\mathbf{n})[\mathbf{d}(\mathbf{n})\times\Gamma(\mathbf{n})]]\right\}\frac{d\mathbf{n}}{4\pi}=0. \qquad (A.6)$$

To solve the set of equations (A.5) and (A.6), we write

$$\boldsymbol{\tau}(\mathbf{n})=[\mathbf{T}(q)\times\mathbf{d}(\mathbf{n})], \qquad (A.7)$$

after which it is not difficult to obtain the relation connecting the quantity $\mathbf{T}(q)$ with the vertex $\boldsymbol{\Gamma}(\mathbf{n})$ (see Eq. (10) of Ref. 8):

$$\left(\omega^2-\frac{1}{5}v_F^2\mathbf{q}^2\right)\mathbf{T}_\parallel+\left(\omega^2-\frac{2}{5}v_F^2\mathbf{q}^2\right)\mathbf{T}_\perp$$
$$=3\Delta\int(\omega+(\mathbf{qn})v_F)[\mathbf{d}(\mathbf{n})[\Gamma(\mathbf{n})\mathbf{d}(\mathbf{n})]]\frac{d\mathbf{n}}{4\pi}. \qquad (A.8)$$

The described procedure permits us, by constructing the vertex functions $\boldsymbol{\Gamma}(\mathbf{n})$ and $\boldsymbol{\gamma}(\mathbf{n})$, to calculate the dynamic susceptibility of the superfluid B phase of ^{3}He, as was accomplished in our work. The choice of $\boldsymbol{\tau}(\mathbf{n})$ in the form (A.7) is used in the solution of the set of equations (A.5) and (A.6) and nothing else is expected of it.

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Magnetic properties of certain terbium alloys with CsCl structure

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The magnetic properties are investigated of solid solutions of the systems $TbCu_{1-x}B_x$ (B = Ag, Zn, Al). When the copper atoms are replaced by zinc and aluminum atoms, the Neel temperature is lowered and the magnetic Curie temperature Θ_p reverses sign at 16 at.% Al and 24 at.% Zn, with $\Theta_p > 0$ for alloys with larger contents of these metals, i.e., the configuration changes from antiferromagnetic to ferromagnetic. The antiferromagnetism is preserved in the solid solution $TbCu_{0.5}Ag_{0.5}$. The results are interpreted on the basis of the RKKY theory.

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To study the role of the conduction electrons in the mechanism of exchange interaction in metallic-conductivity solid solutions having antiferromagnetic or ferromagnetic order, we have investigated the magnetic properties of solid solutions of the system $TbCu_{1-x}B_x$ (where B stands for silver, zinc, or aluminum), where the nonmagnetic copper atoms are replaced by other nonmagnetic silver, zinc, or aluminum atoms, the magneto-active-atom concentration and the lattice parameters remaining constant, the only change being in the number of the conduction electrons per magnetic atom.

The investigated alloys were prepared by arc melting in an atmosphere of pure argon under pressure and were subjected to a homogenizing annealing. An x-ray phase shift analysis of the obtained samples (CuK_α radiation with a nickel filter) has revealed that all the alloys are solid solutions and have a crystal structure of the CsCl type (see Table I). The magnetic properties were investigated with the aid of a pendulum balance in the temperature interval 78–300 K and in magnetic fields of intensity from 1 to 15 kOe.

The results of the investigations have shown that the Neel temperature Θ_N becomes lower with increasing zinc and aluminum content (Figs. 1 and 2), and antiferromagnetic ordering exists in the $TbCu_{1-x}Zn_x$ alloys in the concentration region up to 50 at.% zinc, while in the solid solutions, where the copper atoms are replaced by aluminum atoms, an antiferromagnetic transition is observed in the investigated temperature interval for samples containing 20 at.% aluminum (see the table). In the alloy $TbCu_{0.5}Ag_{0.5}$, where 50 at.% copper is replaced by silver atoms, Θ_N remains practically the same as in the TbCu compound. As to the paramagnetic Curie temperature Θ_p, in the systems $TbCu_{1-x}Zn_x$ and $TbCu_{1-x}Al_x$ at a definite zinc concentration (24 at.%) and aluminum concentration (16 at.%) it reverses sign and becomes positive ($\Theta_p > 0$), while in the alloy containing 50 at.% silver, just as in the TbCu compound, $\Theta_p < 0$ (Fig. 2). The effective magnetic moment μ_{eff} per terbium atom does not depend on the silver, zinc, or aluminum content and corresponds to the moment of the trivalent terbium ion in the ground state $^7F_6$. The table lists the values of Θ_N, Θ_p, and μ_{eff} for all the obtained solid solutions. Using the obtained values of Θ_N and Θ_p and the relations of the molecular-field theory,[1] we have estimated the exchange-interaction parameters J_1 and J_2, which characterize respectively the interaction between the nearest neighbors and the next-to-nearest neighbors. It turned out that $J_1 > 0$ and $J_2 < 0$ for all the investigated alloys (see the table), and the change of these parameters is faster with increasing aluminum content than that of zinc, with J_1 increasing and J_2 decreasing and tending to zero.

This variation of the exchange parameters as a function of the composition explains qualitatively the transition from the type-$\pi\pi 0$ antiferromagnetic configuration,

TABLE I.

Composition	a_0, Å	Θ_N, K	Θ_p, K	μ_{eff}	J_1, K	J_2, K
TbCu	3.480	116	−25	9.72	0.95	−1.1
$TbCu_{0.5}Ag_{0.5}$	3.533	114	−23	9.96	0.94	−1.04
$TbCu_{0.9}Zn_{0.1}$	3.490	112	−18	9.62	0.98	−1.02
$TbCu_{0.8}Zn_{0.2}$	3.501	107	−6	9.59	1.05	−0.88
$TbCu_{0.7}Zn_{0.3}$	3.510	98	13	9.68	1.16	−0.66
$TbCu_{0.6}Zn_{0.4}$	3.519	88	36	9.89	1.3	−0.41
$TbCu_{0.5}Zn_{0.5}$	3.524	76[2]	68	9.82	1.5	−0.04
$TbCu_{0.9}Al_{0.1}$	3.499	107	−12	9.60	1.0	−0.93
$TbCu_{0.8}Al_{0.2}$	3.518	95	11	9.56	1.14	−0.63
$TbCu_{0.7}Al_{0.3}$	3.540	—	42	9.56	—	—
$TbCu_{0.6}Al_{0.4}$	3.547	—	71	9.82	—	—
$TbCu_{0.5}Al_{0.5}$	3.549	—	98	9.39	—	—

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