Phase diagram for a superfluid Fermi gas in a strong magnetic field

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We calculate the strong-coupling corrections to the Ginzburg–Landau free-energy functional in a strong magnetic field and analyze local and global minima of this functional. We show that in the case of a Fermi gas with repulsion the global minima in strong magnetic fields correspond only to the $A1$ and $A2$ phases discovered by Anderson, Brinkman, and Morel.

1. In our previous papers[1,2] we calculated from first principles the strong-coupling corrections to the Ginzburg–Landau free-energy functional in a superfluid Fermi gas with repulsion. These calculations were performed in the absence of magnetic field up to the third order in gas parameter $\lambda=2ap_F/\pi$, where a is the scattering length and p_F is a Fermi momentum. They yield the following expressions for the standard $\beta_1 ... \beta_5$ coefficients:[3]

$$\begin{aligned}
\beta_1&=|\beta_1^{WC}|\left\{-1+\frac{T_C}{2\varepsilon_F}(-76.1\lambda^2+286.5\lambda^3)\right\},\\
\beta_2&=|\beta_1^{WC}|\left\{2+\frac{T_C}{2\varepsilon_F}(-7.20\lambda^2+126.1\lambda^3)\right\},\\
\beta_3&=|\beta_1^{WC}|\left\{2+\frac{T_C}{2\varepsilon_F}(-6.40\lambda^2-16.30\lambda^3)\right\},\\
\beta_4&=|\beta_1^{WC}|\left\{2+\frac{T_C}{2\varepsilon_F}(-48.4\lambda^2-233.0\lambda^3)\right\},\\
\beta_5&=|\beta_1^{WC}|\left\{-2+\frac{T_C}{2\varepsilon_F}(-110\lambda^2-277.5\lambda^3)\right\},
\end{aligned}\tag{1}$$

where $T_C \simeq \varepsilon_F \exp\{-128/(\pi\lambda)^2\}$ is the super fluid transition temperature of a Fermi gas with repulsion in a triplet p-wave state,[4] ε_F is the Fermi energy, $\beta_1^{WC} = -(N(0)/T_C^2) \times(7\zeta(3)/240\pi^2)$ is the Ginzburg–Landau coefficient β_1 in the weak-coupling approximation,[3] $\zeta(z)$ is the Riemann zeta function ($\zeta(3)=1.202$), and $N(0)=mp_F/\pi^2$ is the density of states at the Fermi surface. The analysis of the phase diagram performed in Ref. 2 with the coefficients β_i from Ref. 5 shows the tendency of the global minima to shift from the B to the A phase as λ increases approaching unity. (Formal application of

Ref. 5 gives the transition from the B to the A phase at $\lambda=1.26$.) All the other phases of a triplet superfluid either lie above the B and A phases on the energy scale or do not even correspond to local minima of the energy.

Note that the results of the exact calculations of $\beta_1 ... \beta_5$ differ rather significantly to third order in λ from the results of the standard $s-p$ approximation but lead to a qualitatively similar phase diagram. Note also that the large values of the numerical coefficients (~ 200) near the λ^3 terms in $\beta_1 ... \beta_5$ are due to the tensor structure of the order parameter and to the large number of diagrams.

2. In the present letter we generalize the result of Refs. 1 and 2 on the case of a strong magnetic field. To be specific we consider magnetic fields larger than paramagnetic limit for the B phase: $H > H_p = T_C/\mu_B$. In this case $S_z=0$—the projection of the spin $S=1$ of the triplet Cooper pair is totally suppressed. In other words the $\Delta_{\uparrow\downarrow}$ component of the 2×2 order parameter matrix $\Delta_{\alpha\beta}$ is zero and there are only two critical temperatures $T_C^{\uparrow\uparrow}$ and $T_C^{\downarrow\downarrow}$, corresponding to $S_z=1$ ($\Delta_{\uparrow\uparrow}$) and $S_z=-1(\Delta_{\downarrow\downarrow})$. Exact analytical expressions for $T_C^{\uparrow\uparrow}$ and $T_C^{\downarrow\downarrow}$ as a functional of the gas parameter and magnetic field (or spin polarization $\eta = N_\uparrow - N_\downarrow / N_\uparrow + N_\downarrow$) were obtained in Ref. 6. The results show a strongly nonmonotonic behavior of $T_C^{\uparrow\uparrow}$, with a large and broad maximum at $\eta=0.48$. $T_C^{\downarrow\downarrow}$ is a monotonically decreasing function of η. Note that $T_C \sim 1$ mK for real ^{3}He, and hence $H_p \sim 1$ T corresponds to $\eta_p \approx (2/3)(\mu_B H_p/\varepsilon_F) \approx 0.004$.

The Ginzburg–Landau free-energy functional in strong magnetic fields can be rewritten in the following convenient form:

$$\Delta F \equiv F_S - F_N = a_\uparrow \Delta_{0\uparrow\uparrow}^2 M_+ + b_\uparrow \Delta_{0\uparrow\uparrow}^4 (2M_+^2 + |N_+|^2) + a_\downarrow \Delta_{0\downarrow\downarrow}^2 M_- + b_\downarrow \Delta_{0\downarrow\downarrow}^4 (2M_-^2 + |N_-|^2) + \Delta_{0\uparrow\uparrow}^2 \Delta_{0\downarrow\downarrow}^2 (\alpha M_+ M_- + \beta |R|^2 + \gamma |P|^2), \qquad (2)$$

where $M_+ = A_{+k}^* A_{+k}$, $M_- = A_{-k}^* A_{-k}$, $N_+ = A_{+k} A_{+k}$, $N_- = A_{-k} A_{-k}$, $R = A_{+k} A_{-k}$, and $P = A_{+k} A_{-k}^*$. In formula (2):

$$A_{+k} = iA_{2k} - A_{1k}; \quad A_{-k} = iA_{2k} + A_{1k}; \quad a_\uparrow = \frac{1}{3} N(0) \ln \frac{T}{T_C^{\uparrow\uparrow}};$$

$$a_\downarrow = \frac{1}{3} N(0) \ln \frac{T}{T_C^{\downarrow\downarrow}}; \quad b_\uparrow = b_\downarrow = \frac{N(0)}{15 T_C^2} \left[1 + O\left(\lambda^2 \frac{T_C}{\varepsilon_F} \right) \right];$$

$\alpha = 2\beta_2 + 2\beta_5$, $\beta = 4\beta_1 + 2\beta_3$ and $\gamma = 2\beta_4 + 2\beta_5$ are the coefficients in front of the strong-coupling invariants with the structure $\Delta_{0\uparrow\uparrow}^2 \Delta_{0\downarrow\downarrow}^2$. Note that the 3×3 matrix A_{ik} in (2) is connected with the 2×2 matrix $\Delta_{\alpha\beta}$ via the standard formula[3,7]

$$\Delta_{\alpha\beta} = \Delta_{0\alpha\beta}(T) i (\sigma_2 \sigma_i)_{\alpha\beta} A_{ik} n_k \quad \text{and} \quad A_{3k} = 0 \quad \text{for} \quad H > H_p.$$

Let us consider first the case of weak coupling. In this case $\alpha = \beta = \gamma = 0$, and the Ginzburg–Landau free-energy functional reduces to that of two independent superfluids with $S_z=1$ ($\Delta_{\uparrow\uparrow}$) and $S_z=-1$ ($\Delta_{\downarrow\downarrow}$). Direct minimization of the free energy, $\delta \Delta F / \delta A_{\pm k} = \delta \Delta F / \delta A_{\pm k}^* = 0$, yields two minima of the free energy.

For $T < T_C^{\downarrow\downarrow} < T_C^{\uparrow\uparrow}$ the first one corresponds to

$$\Delta_{0\uparrow\uparrow}^2 M_+ = -\frac{a_\uparrow}{4b_\uparrow}; \quad \Delta_{0\downarrow\downarrow}^2 M_- = -\frac{a_\downarrow}{4b_\downarrow}; \quad N_+ = N_- = 0,$$

while the second one corresponds to

$$\Delta_{0\uparrow\uparrow}^2 M_+ = -\frac{a_\uparrow}{6b_\uparrow}; \quad \Delta_{0\downarrow\downarrow}^2 M_- = -\frac{a_\downarrow}{6b_\downarrow}; \quad |N_+| = M_+; \quad |N_-| = M_-.$$

The Ginzburg–Landau free energy at the first extremum is given by

$$\Delta F = -\frac{a_\uparrow^2}{8b_\uparrow} - \frac{a_\downarrow^2}{8b_\downarrow} \tag{3}$$

while at the second we have

$$\Delta F = -\frac{a_\uparrow^2}{12b_\uparrow} - \frac{a_\downarrow^2}{12b_\downarrow}. \tag{4}$$

Expressions (3), (4) show that the global minima correspond to the first case with $N_+ = N_- = 0$. For $T_C^{\downarrow\downarrow} < T < T_C^{\uparrow\uparrow}$ the global extremum conditions are satisfied for three degenerate phases:

$$A1 = \frac{1}{2}\begin{pmatrix} 1 & i & 0 \\ i & -1 & 0 \\ 0 & 0 & 0 \end{pmatrix}; \quad E = \frac{1}{\sqrt{2}}\begin{pmatrix} 1 & 0 & 0 \\ 0 & 1 & 0 \\ 0 & 0 & 0 \end{pmatrix}; \quad \sigma_3 = \frac{1}{\sqrt{2}}\begin{pmatrix} 1 & 0 & 0 \\ 0 & -1 & 0 \\ 0 & 0 & 0 \end{pmatrix}.$$

(Note that E phase coincides with the well-known planar phase.)

For $T < T_C^{\downarrow\downarrow}$ the global extremum corresponds again to the E and σ_3 phases together with the $A2$ phase. The last one has the form

$$A2 = \frac{1}{\sqrt{2(1-\delta^2)}}\begin{pmatrix} 1 & i & 0 \\ i\delta & -\delta & 0 \\ 0 & 0 & 0 \end{pmatrix}, \quad \delta = \frac{\Delta_{0\uparrow\uparrow} - \Delta_{0\downarrow\downarrow}}{\Delta_{0\uparrow\uparrow} + \Delta_{0\downarrow\downarrow}}.$$

The appearance of the strong-coupling term proportional to $\Delta_{0\uparrow\uparrow}^2 \Delta_{0\downarrow\downarrow}^2$ mixes the $S_z = 1$ and $S_z = -1$ superfluids below the lower critical temperature $T_C^{\downarrow\downarrow}$. There are two consequences of this fact: the first is an increase of $T_C^{\downarrow\downarrow}$ due to the interaction between the up–up and down–down Bose condensates. Of course $\Delta_{0\uparrow\uparrow}$ is also renormalized below $T_C^{\downarrow\downarrow}$. The second consequence is a lifting of the degeneracy between the phases which correspond to the global minima. Indeed, for the $A2$ phase we have

$$N_+ = N_- = 0; \quad R = 0; \quad |P|^2 = M_+ M_-, \tag{5}$$

and the Ginzburg–Landau free energy at this extremum is given by

$$\Delta F_{A2} = -\frac{a_\uparrow^2}{8b_\uparrow} - \frac{a_\downarrow^2}{8b_\downarrow} + (\alpha + \gamma)\frac{a_\uparrow a_\downarrow}{16 b_\uparrow b_\downarrow}. \tag{6}$$

At the same time, for the planar and σ_3 phases

$$N_+ = N_- = 0; \quad P = 0; \quad |R|^2 = M_+ M_-, \tag{7}$$

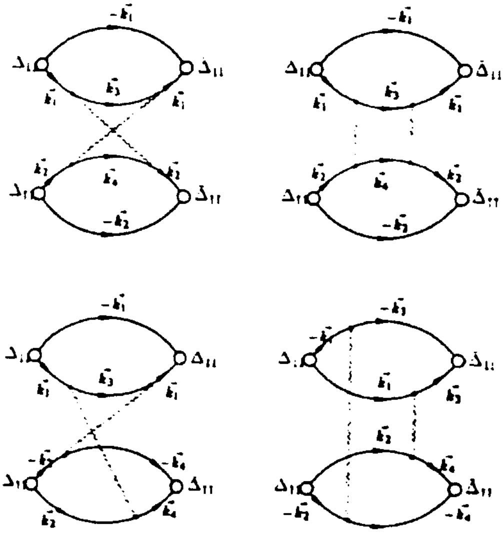

FIG. 1. Diagrams which determine strong-coupling coefficients α, β, and γ.

and, accordingly, the Ginzburg–Landau free energy reads

$$\Delta F_{E,\sigma_3} = -\frac{a_\uparrow^2}{8b_\uparrow} - \frac{a_\downarrow^2}{8b_\downarrow} + (\alpha+\beta)\frac{a_\uparrow a_\downarrow}{16 b_\uparrow b_\downarrow}. \quad (8)$$

Formulas (6) and (8) show that to answer the question of which phase has a lower energy we need to calculate the combinations $(\alpha+\beta)$ and $(\alpha+\gamma)$ of the strong-coupling coefficients. These calculations can be performed in the same way and from the same strong-coupling diagrams as in Refs. 1 and 9. These diagrams are shown in Fig. 1. In analytical form they can be written as

$$-\frac{T^3}{2}\sum_{\omega_1\omega_2\omega_3}\int \frac{d^3\mathbf{k}_1 d^3\mathbf{k}_2 d^3\mathbf{k}_3}{(2\pi)^9}|\Gamma_{+-,+-}(\mathbf{k}_1,\mathbf{k}_2;\mathbf{k}_3,\mathbf{k}_4)|^2$$

$$\times K_\downarrow(\mathbf{k}_1,\omega_1)K_\uparrow(\mathbf{k}_4,\omega_4)G_\downarrow(\mathbf{k}_3,\omega_3)G_\uparrow(\mathbf{k}_2,\omega_2)|\Delta_{\downarrow\downarrow}(\mathbf{k}_1)|^2|\Delta_{\uparrow\uparrow}(\mathbf{k}_4)|^2$$

$$-\frac{T^3}{2}\sum_{\omega_1\omega_2\omega_3}\int \frac{d^3\mathbf{k}_1 d^3\mathbf{k}_2 d^3\mathbf{k}_3}{(2\pi)^9}|\Gamma_{+-,+-}(\mathbf{k}_1,\mathbf{k}_2;\mathbf{k}_3,\mathbf{k}_4)|^2$$

$$\times K_\downarrow(\mathbf{k}_1,\omega_1)K_\uparrow(\mathbf{k}_2,\omega_2)G_\downarrow(\mathbf{k}_3,\omega_3)G_\uparrow(\mathbf{k}_4,\omega_4)|\Delta_{\downarrow\downarrow}(\mathbf{k}_1)|^2|\Delta_{\uparrow\uparrow}(\mathbf{k}_2)|^2$$

$$+2T^3\sum_{\omega_1\omega_2\omega_3}\int \frac{d^3\mathbf{k}_1 d^3\mathbf{k}_2 d^3\mathbf{k}_3}{(2\pi)^9}\Gamma_{+-,+-}(\mathbf{k}_1,\mathbf{k}_2;\mathbf{k}_3,\mathbf{k}_4)\Gamma_{+-,+-}(\mathbf{k}_3,-\mathbf{k}_2;\mathbf{k}_1,-\mathbf{k}_4)$$

$$\times K_\downarrow(\mathbf{k}_1,\omega_1)G_\downarrow(\mathbf{k}_3,\omega_3)F_\uparrow(\mathbf{k}_2,\omega_2)F_\uparrow(\mathbf{k}_4,\omega_4)|\Delta_{\downarrow\downarrow}(\mathbf{k}_1)|^2\Delta_{\uparrow\uparrow}(\mathbf{k}_2)\Delta^+_{\uparrow\uparrow}(\mathbf{k}_4)$$

$$-T^3\sum_{\omega_1\omega_2\omega_3}\int \frac{d^3\mathbf{k}_1 d^3\mathbf{k}_2 d^3\mathbf{k}_3}{(2\pi)^9}|\Gamma_{+-,+-}(\mathbf{k}_1,\mathbf{k}_2;\mathbf{k}_3,\mathbf{k}_4)|^2F_\downarrow(\mathbf{k}_1,\omega_1)F_\downarrow(\mathbf{k}_3,\omega_3)$$

$$\times F_\uparrow(\mathbf{k}_2,\omega_2)F_\uparrow(\mathbf{k}_4,\omega_4)\Delta_{\downarrow\downarrow}(\mathbf{k}_1)\Delta^+_{\downarrow\downarrow}(\mathbf{k}_3)\Delta_{\uparrow\uparrow}(\mathbf{k}_2)\Delta^+_{\uparrow\uparrow}(\mathbf{k}_4), \tag{9}$$

where $\mathbf{k}_4=-\mathbf{k}_3+\mathbf{k}_1+\mathbf{k}_2$, $G_\uparrow=(i\omega-\xi_\uparrow)^{-1}$ and $G_\downarrow=(i\omega-\xi_\downarrow)^{-1}$ are the Matsubara Green functions for up and down spins, $\omega=\pi T(2n+1)$ is a Matsubara frequency, $\xi_{\uparrow(\downarrow)}=(p^2-p^2_{F\uparrow(\downarrow)})/(2m)$ is the spectrum for up (down) spins, $K_\uparrow(\mathbf{p},\omega)=G^2_\uparrow(\mathbf{p},\omega)G_\uparrow(-\mathbf{p},-\omega)$, $F_\downarrow(\mathbf{p},\omega)=G_\downarrow(-\mathbf{p},-\omega)G_\downarrow(\mathbf{p},\omega)$, $\Delta_{\uparrow\uparrow}(\mathbf{p})=\Delta_{0\uparrow\uparrow}(T)A_{+k}n_k$, $\Delta_{\downarrow\downarrow}(\mathbf{p})=\Delta_{0\downarrow\downarrow}(T)A_{-k}n_k$, $\Gamma_{+-,+-}$ is a total vertex for up and down incoming and outgoing spins. In the first two orders of perturbation theory the total vertex is given by:

$$\Gamma_{+-,+-}\equiv\Gamma_{+-}=g+g^2C_{+-}(\mathbf{p}_1+\mathbf{p}_2)+g^2\Pi_{+-}(\mathbf{p}_1-\mathbf{p}_4), \tag{10}$$

where $g=4\pi a/m$ is a coupling constant, C_{+-} is a Cooper loop formed by up and down spins:

$$C_{+-}(k)=\frac{m}{8\pi^2 k}\left[Q(p_{F\uparrow})+Q(p_{F\downarrow})\right],$$

$$Q(p_{F\uparrow})=p^2_{F\uparrow}\ln\frac{\left(p_{F\uparrow}+\frac{k}{2}\right)^2-a^2}{\left(p_{F\uparrow}-\frac{k}{2}\right)^2-a^2}+\left(\frac{k}{2}+a\right)^2\ln\left|\frac{p_{F\uparrow}-\frac{k}{2}-a}{p_{F\uparrow}+\frac{k}{2}+a}\right|+2kp_{F\uparrow}$$

$$+\left(\frac{k}{2}-a\right)^2\ln\left|\frac{p_{F\uparrow}-\frac{k}{2}+a}{p_{F\uparrow}+\frac{k}{2}-a}\right|,$$

$$a^2=\Delta^2-\frac{k^2}{4};\quad \Delta^2=\frac{p^2_{F\uparrow}+p^2_{F\downarrow}}{2};\quad k=|\mathbf{p}_1+\mathbf{p}_2|,$$

Π_{+-} is an exchange contribution which is the same as that of a polarization loop for a short-range interaction:

$$\Pi_{+-}(q)=\frac{m}{8\pi^2 q}\left\{\left[p^2_{F\uparrow}-\left(\frac{q^2+\tilde{\Delta}^2}{2q}\right)^2\right]\ln\left|\frac{q^2+\tilde{\Delta}^2+2p_{F\uparrow}q}{q^2+\tilde{\Delta}^2-2p_{F\uparrow}q}\right|+\frac{q^2+\tilde{\Delta}^2}{q}p_{F\uparrow}\right.$$
$$\left.+\left[p^2_{F\downarrow}-\left(\frac{q^2-\tilde{\Delta}^2}{2q}\right)^2\right]\ln\left|\frac{q^2-\tilde{\Delta}^2+2p_{F\downarrow}q}{q^2-\tilde{\Delta}^2-2p_{F\downarrow}q}\right|+\frac{q^2-\tilde{\Delta}^2}{q}p_{F\downarrow}\right\},$$

$$q=|\mathbf{p}_1-\mathbf{p}_4|;\quad \tilde{\Delta}^2=p^2_{F\uparrow}-p^2_{F\downarrow}.$$

Unfortunately, the complete analytical dependence on $\nu=p_{F\downarrow}/p_{F\uparrow}=(1-\eta/1+\eta)^{1/3}$ for the strong-coupling coefficients α, β, and γ can be recovered only for the leading terms ($\sim\lambda^2(T_C/\varepsilon_F)|\beta_1^{WC}|$). The results are the following:

$$\alpha^{(2)}=126.2\,\frac{T_C}{2\varepsilon_F}\,\lambda^2|\beta_1^{WC}|\left(\frac{2\nu^3}{1+\nu^3}\right)^{1/3}\{-0.18-1.68\nu^2\},$$

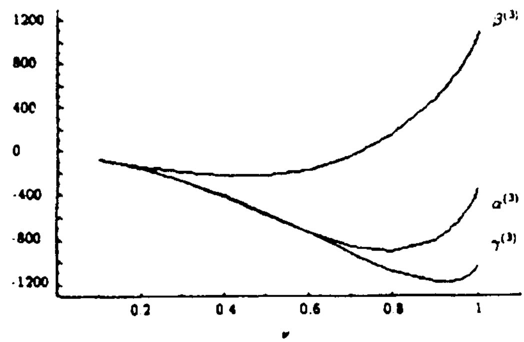

FIG. 2. Magnetic field dependence of strong-coupling coefficients α, β, and γ in third order in gas parameter $\lambda=2ap_F/\pi$ (in units $(T_C/2\varepsilon_F)\lambda^3|\beta_1^{WC}|$).

$$\beta^{(2)}=126.2\,\frac{T_C}{2\varepsilon_F}\,\lambda^2|\beta_1^{WC}|\left(\frac{2\nu^3}{1+\nu^3}\right)^{1/3}\{-1.0-1.51\nu^2\}, \tag{11}$$

$$\gamma^{(2)}=126.2\,\frac{T_C}{2\varepsilon_F}\,\lambda^2|\beta_1^{WC}|\left(\frac{2\nu^3}{1+\nu^3}\right)^{1/3}\{-1.0-1.51\nu^2\}.$$

For the maximum magnetic field $H=20$ T ($\nu=0.95$, $\eta=0.08$) which we can create by brute force, α, β, and γ differ only by 10% from their zero-field values. Coefficients β and γ coincide identically in this order even in nonzero magnetic field. That is why the $A2$ phase, the planar phase, and the σ_3 phase are still degenerate. To lift the degeneracy we need to calculate α, β, and γ in the next order. In zero magnetic field we have from (1):

$$\alpha^{(3)}=-302.9\,\frac{T_C}{2\varepsilon_F}\,\lambda^3|\beta_1^{WC}|;\quad \beta^{(3)}=1114\,\frac{T_C}{2\varepsilon_F}\,\lambda^3|\beta_1^{WC}|;$$

$$\gamma^{(3)}=-1021\,\frac{T_C}{2\varepsilon_F}\,\lambda^3|\beta_1^{WC}|<\beta.$$

These values show that $A2$ is energetically favorable. The complete magnetic field dependence of α, β, and γ in third order in $\lambda=2ap_F/\pi$ is presented in Fig. 2. At small spin-polarizations $\eta\approx 3/2(1-\nu)$ all the coefficients α, β and γ behave linearly in η. For $H=20$ T they again differ only by 10% from their zero-field values. However for spin polarization $\eta=0.3$ ($\nu=0.8$), which can be created by a rapid melting of the ^{3}He crystal,[10] this difference becomes rather significant. Moreover α and β have a nonmonotonic dependence on ν with a pronounced minimum, while γ is a monotonically decreasing function. For $\nu\to 0$ ($\eta\to 1$) there are no down spins in the system, and that is why α, β, and γ saturate to the same zero value. The most important thing is that γ is smaller than β for all magnetic fields. Hence the $A2$ phase is a global extremum of the free energy.

3. In conclusion, we have calculated the strong-coupling corrections to the free energy of a triplet superfluid Fermi gas in high magnetic fields and have found that the phase diagram of the system contains only the $A1$ and $A2$ phases. Note that our theory is exact for low-density systems (e.g. ^{3}He–^{4}He mixtures) where $\lambda<1$. In pure ^{3}He one has $\lambda\sim 1$, and our results can be used only as a qualitative estimate. Nevertheless due to the *intrinsic* character of a superfluid transition in our model the philosophy of the authors is the following: if there are no exotic phases at low densities, there is little chance of obtaining them at higher densities. Hence we consider the present calculations as a strong argument against the possibility of a new phase of ^{3}He (see Refs. 11–13) in a high magnetic field. In view of our previous results in zero magnetic field (no new phase either),[1,2] we consider any novelty in a phase diagram of superfluid ^{3}He to be very unlikely (at least in three dimensions).

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