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PHASE DIAGRAM FOR THE 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 the global minima in strong magnetic field correspond in the case of Fermi-gas with repulsion only to A_1 and A_2 phases discovered by Anderson, Brinkman and Morel [5].

1. In our previous papers [1, 2] we calculated from the 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 $\lambda = \frac{2ap_F}{\pi}$, where a is the scattering length and p_F is a Fermi momentum. They yield the following expressions for the standard $\beta_1 \ldots \beta_5$ coefficients [3]:

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

where $T_C \simeq \varepsilon_F \exp\left\{-\frac{128}{(\pi\lambda)^2}\right\}$ is a superfluid transition temperature of a Fermi–gas with repulsion in a triplet p-wave state [4], ε_F is a Fermi energy, $\beta_1^{WC} = -\frac{N(0)}{T_C^2}\frac{7\zeta(3)}{240\pi^2}$ is the Ginzburg–Landau coefficient β_1 in a weak-coupling approximation [3], $\zeta(z)$ is the Riemann zeta-function ($\zeta(3) = 1.202$) and $N(0) = \frac{mp_F}{\pi^2}$ is the density of states at the Fermi surface. The analysis of the phase diagram performed in paper [2] with the coefficients β_i from [5] shows the tendency of shifting of the global minima from the B to A phase when λ – increases approaching unity. (Formal application of [5] gives the transition from B to A phase for $\lambda = 1.26$.) 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 the local minima of the energy.

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

2. In the present letter we generalize the result of papers [1,2] on the case of a strong magnetic field. To be specific we consider magnetic fields larger then paramagnetic limit for B phase: $H > H_p = \frac{T_C}{\mu_B}$. In this case $S_z = 0$ – projection of the spin $S = 1$ of the triplet Cooper pair is totally suppressed. In other words $\Delta_{\uparrow\downarrow}$ component of 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 polarisation $\eta = \frac{N_\uparrow - N_\downarrow}{N_\uparrow + N_\downarrow}$) were obtained in [6]. The results show 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 for real ^{3}He $T_C \sim 1\,\mathrm{mK}$ and hence $H_p \sim 1\,\mathrm{T}$ corresponds to $\eta_p \approx \frac{2}{3}\frac{\mu_B H_p}{\varepsilon_F} \approx 0.004$.

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

$$\begin{aligned} \Delta F \equiv F_S - F_N &= a_\uparrow \Delta^2_{0\uparrow\uparrow} M_+ + b_\uparrow \Delta^4_{0\uparrow\uparrow}(2M_+^2 + |N_+|^2) + \\ &+ a_\downarrow \Delta^2_{0\downarrow\downarrow} M_- + b_\downarrow \Delta^4_{0\downarrow\downarrow}(2M_-^2 + |N_-|^2) + \\ &+ \Delta^2_{0\uparrow\uparrow}\Delta^2_{0\downarrow\downarrow}(\alpha M_+ M_- + \beta |R|^2 + \gamma |P|^2), \end{aligned} \tag{2}$$

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

$$A_{+k} = iA_{2k} - A_{1k}; \quad A_{-k} = iA_{2k} + A_{1k};$$

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

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

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

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

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

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

while the second one corresponds to

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

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

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

while on the second one 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 corresponds 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 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 $S_z = 1$ and $S_z = -1$ superfluids below the lower critical temperature $T_C^{\downarrow\downarrow}$. There are two consequences of this fact: the first one is the increase of $T_C^{\downarrow\downarrow}$ due to the interaction between up-up and down-down bose-condensates. Of course $\Delta_{0\uparrow\uparrow}$ is also renormalized below $T_C^{\downarrow\downarrow}$. The second one is lifting the degeneracy between the phases which correspond to the global minima. Indeed, for $A2$ phase we have

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

and the Ginzburg–Landau free energy on 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

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

and, accordingly, the Ginzburg–Landau free energy reads

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

Formulae (6) and (8) show that to answer the question 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

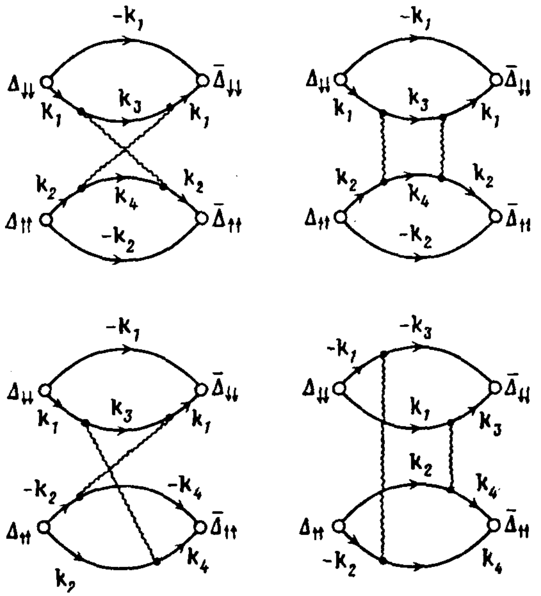

Fig.1. Diagrams which determine strong-coupling coefficients α, β and γ

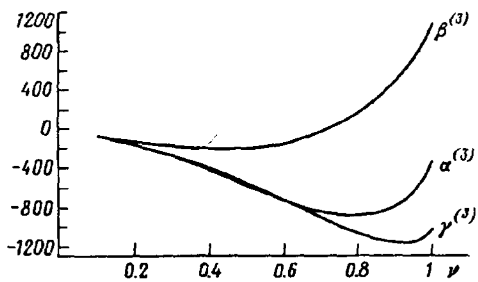

Fig.2. Magnetic field dependence of strong - coupling coefficients α, β and γ in third order in gas parameter $\lambda = \frac{2ap_F}{\pi}$ (in units $\frac{T_C}{2\varepsilon_F}\lambda^3 \left|\beta_1^{W.C.}\right|$

and from the same strong-coupling diagrams as in papers [1, 9]. These diagrams are shown on Fig.1. In analytical form they can be written as

$$-\frac{T^3}{2}\sum_{\omega_1\omega_2\omega_3}\int\frac{d^3\mathbf{k}_1 d^3\mathbf{k}_2 d^3\mathbf{k}_3}{(2\pi)^9}\left|\Gamma_{+-,+-}(\mathbf{k}_1,\mathbf{k}_2;\mathbf{k}_3,\mathbf{k}_4)\right|^2\times$$
$$\times K_\downarrow(\mathbf{k}_1,\omega_1)K_\uparrow(\mathbf{k}_4,\omega_4)G_\downarrow(\mathbf{k}_3,\omega_3)G_\uparrow(\mathbf{k}_2,\omega_2)|\Delta_{\downarrow\downarrow}(\mathbf{k}_1)|^2|\Delta_{\uparrow\uparrow}(\mathbf{k}_4)|^2-$$
$$-\frac{T^3}{2}\sum_{\omega_1\omega_2\omega_3}\int\frac{d^3\mathbf{k}_1 d^3\mathbf{k}_2 d^3\mathbf{k}_3}{(2\pi)^9}\left|\Gamma_{+-,+-}(\mathbf{k}_1,\mathbf{k}_2;\mathbf{k}_3,\mathbf{k}_4)\right|^2\times$$
$$\times K_\downarrow(\mathbf{k}_1,\omega_1)K_\uparrow(\mathbf{k}_2,\omega_2)G_\downarrow(\mathbf{k}_3,\omega_3)G_\uparrow(\mathbf{k}_4,\omega_4)|\Delta_{\downarrow\downarrow}(\mathbf{k}_1)|^2|\Delta_{\uparrow\uparrow}(\mathbf{k}_2)|^2+ \qquad (9)$$
$$+2T^3\sum_{\omega_1\omega_2\omega_3}\int\frac{d^3\mathbf{k}_1 d^3\mathbf{k}_2 d^3\mathbf{k}_3}{(2\pi)^9}\Gamma_{+-,+-}(\mathbf{k}_1,\mathbf{k}_2;\mathbf{k}_3,\mathbf{k}_4)\Gamma_{+-,+-}(\mathbf{k}_3,-\mathbf{k}_2;\mathbf{k}_1,-\mathbf{k}_4)\times$$
$$\times K_\downarrow(\mathbf{k}_1,\omega_1)G_\downarrow(\mathbf{k}_3,\omega_3)F_\uparrow(\mathbf{k}_2,\omega_2)F_\uparrow(\mathbf{k}_4,\omega_4)|\Delta_{\downarrow\downarrow}(\mathbf{k}_1)|^2\Delta_{\uparrow\uparrow}(\mathbf{k}_2)\Delta^+_{\uparrow\uparrow}(\mathbf{k}_4)-$$
$$-T^3\sum_{\omega_1\omega_2\omega_3}\int\frac{d^3\mathbf{k}_1 d^3\mathbf{k}_2 d^3\mathbf{k}_3}{(2\pi)^9}\left|\Gamma_{+-,+-}(\mathbf{k}_1,\mathbf{k}_2;\mathbf{k}_3,\mathbf{k}_4)\right|^2F_\downarrow(\mathbf{k}_1,\omega_1)F_\downarrow(\mathbf{k}_3,\omega_3)\times$$
$$\times F_\uparrow(\mathbf{k}_2,\omega_2)F_\uparrow(\mathbf{k}_4,\omega_4)\Delta_{\downarrow\downarrow}(\mathbf{k}_1)\Delta^+_{\downarrow\downarrow}(\mathbf{k}_3)\Delta_{\uparrow\uparrow}(\mathbf{k}_2)\Delta^+_{\uparrow\uparrow}(\mathbf{k}_4),$$

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 a spectrum for up (down) spins, $K_\uparrow(\mathrm{p},\omega) = G^2_\uparrow(\mathrm{p},\omega)G_\uparrow(-\mathrm{p},-\omega)$, $F_\downarrow(\mathrm{p},\omega) = G_\downarrow(-\mathrm{p},-\omega)G_\downarrow(\mathrm{p},\omega)$, $\Delta_{\uparrow\uparrow}(\mathrm{p}) = \Delta_{0\uparrow\uparrow}(T)A_{+k}n_k$, $\Delta_{\downarrow\downarrow}(\mathrm{p}) = \Delta_{0\downarrow\downarrow}(T)A_{-k}n_k$, $\Gamma_{+-,+-}$ - 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:

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

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

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

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

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

Π_{+-} is an exchange contribution which coincides with a polarisation loop for a short-range interaction:

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

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

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

$$\alpha^{(2)} = 126.2\frac{T_C}{2\varepsilon_F}\lambda^2|\beta_1^{WC}|\left(\frac{2\nu^3}{1+\nu^3}\right)^{1/3}\{-0.18 - 1.68\nu^2\},$$
$$\beta^{(2)} = 126.2\frac{T_C}{2\varepsilon_F}\lambda^2|\beta_1^{WC}|\left(\frac{2\nu^3}{1+\nu^3}\right)^{1/3}\{-1.0 - 1.51\nu^2\}, \tag{11}$$
$$\gamma^{(2)} = 126.2\frac{T_C}{2\varepsilon_F}\lambda^2|\beta_1^{WC}|\left(\frac{2\nu^3}{1+\nu^3}\right)^{1/3}\{-1.0 - 1.51\nu^2\}.$$

For the maximal magnetic field $H = 20\,\mathrm{T}$ ($\nu = 0.95$, $\eta = 0.08$) which we can create by brute force α, β and γ differ only on 10% from their zero-field values. Coefficients β and γ identically coincide in this order even in nonzero magnetic field. That is why the $A2$-phase, planar phase and 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):

$$\alpha^{(3)} = -302.9\frac{T_C}{2\varepsilon_F}\lambda^3\left|\beta_1^{WC}\right|; \ \beta^{(3)} = 1114\frac{T_C}{2\varepsilon_F}\lambda^3\left|\beta_1^{WC}\right|; \ \gamma^{(3)} = -1021\frac{T_C}{2\varepsilon_F}\lambda^3\left|\beta_1^{WC}\right| < \beta.$$

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-polarisations $\eta \approx \frac{3}{2}(1-\nu)$ all the coefficients α, β and γ behave linearly in η. For $H = 20\,\mathrm{T}$ they again differ only on 10% from their zero-field values. However for spin polarisation $\eta = 0.3$ ($\nu = 0.8$) which can be created by a rapid melting of ^{3}He crystal [10] this difference becomes rather significant. Moreover α and β have a nonmonotonic dependence from 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 $A2$ 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 a high magnetic fields and found that the phase diagram of the system contains only $A1$ 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 $\lambda \sim 1$ 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 a small chance to obtain them at higher densities. Hence we consider the present calculations as a strong argument against the possibility of a new phase of ^{3}He [11–13] in a high magnetic field. Having in mind 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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