Leading Temperature Corrections to Fermi-Liquid Theory in Two Dimensions

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We calculate the basic parameters of the Fermi liquid: the scattering vertex, the Landau interaction function, the effective mass, and physical susceptibilities for a model of two-dimensional (2D) fermions with a short-ranged interaction at nonzero temperature. The leading temperature dependences of the spin components of the scattering vertex, the Landau function, and the spin susceptibility are found to be linear. T-linear terms in the effective mass and in the "charge-sector" quantities are found to cancel to second order in the interaction, but the cancellation is argued not to be generic. The connection with previous studies of the 2D Fermi-liquid parameters is discussed.

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The question of the low-energy behavior of two-dimensional (2D) Fermi liquid (FL) is of long-standing and fundamental importance. One crucial motivation has been the non-Fermi-liquid behavior observed in high-T_c superconductors above T_c [1]. In this context the existence and stability of the FL in $d > 1$ has been extensively investigated [2]. At least for short-ranged interactions and in the absence of coupling to gauge fields, the FL is stable, provided that standard conditions [3] are satisfied.

Rather surprisingly, the issue of the leading temperature corrections to the parameters of a *stable* FL, which is of intrinsic theoretical interest and has important implications for the theory of quantum critical phenomena (as explained in Ref. [4]), remains a subject of controversy. For example, it was found [5] that the leading T correction to the specific heat coefficient $\gamma = C/T$ was $T^2 \ln T$ in 3D and T in 2D [6,7]. Whether the spin and charge susceptibilities display a similar anomalous (i.e., non-T^2) temperature dependence is a subject of a contradictory literature: see, e.g., Ref. [8], and references therein. The prevailing conclusion was that of Carneiro and Pethick [8] who found no leading $T^2 \ln T$ correction to the spin susceptibility of the 3D FL. Their arguments imply that terms $\propto T$ are absent in 2D.

The heuristic *a posteriori* argument for the absence of anomalous terms in T or in q in response functions is that, although these terms are known to occur in individual diagrams, they cancel in "consistently calculated" physical quantities due to Ward identities. We note that in the existing literature concerning this point it is assumed that the crucial coupling is between quasiparticles and long-wavelength collective modes. However, the existence of "$2k_F$ singularities," i.e., anomalous temperature terms coming from processes involving large ($\sim 2k_F$) momentum transfers, was already pointed out by Misawa for 3D FL in the early 1970's [9]. Apparently due to the lack of experimental evidence of a $T^2 \ln T$ term in the susceptibility of a generic 3D FL and also because Misawa's results rely on the analysis of selected diagrams [cf. the beginning of this paragraph], they were widely disregarded in favor of those of Carneiro and Pethick. In the context of semiconductor physics, Stern was the first to note [10] that in a 2D electron gas the electron scattering rate was proportional to T due to $2k_F$ effects. The consequences of the $2k_F$ effects for the leading T dependence of 2D FL quantities were not considered in the literature until recently.

The issue of the leading correction to 2D FL parameters was recently revived by several papers. Belitz *et al.* [11] concluded from perturbative analysis combined with power counting that the leading q dependence of the spin susceptibility was $|q|$ in 2D ($q^2 \ln q$ in 3D). They did not find the analogous T correction explicitly, but concluded that one should generally expect a linear-T term in the 2D FL susceptibility ($T^2 \ln T$ in 3D).

Sénéchal and one of us [12] predicted the occurrence of the linear-T corrections to the FL vertices from one-loop renormalization group (RG) calculations. This implies the appearance of the linear-T corrections in other FL quantities, but explicit calculations were not done.

Hirashima and Takahashi [13] performed numerical analyses of perturbative expressions which appeared to confirm the prediction of Belitz *et al.* [11] for 2D susceptibility. However, due to numerical difficulties in handling divergences in some terms they were unable even to determine the sign of the coefficient in the leading q term. Also contrary to Belitz *et al.* who focused on the long-wavelength contributions, the authors of Ref. [11] emphasized the crucial role of $2k_F$ contributions in their findings. Following this, Misawa conjectured the phenomenological form for the free energy [7] which results in the linear-T term in the 2D spin susceptibility and in the coefficient γ, and agrees well with the numerical calculations in lowest order [14].

Despite this work a systematic analytic study of anomalous terms in the FL parameters and their relationship to Ward identities is warranted. To elucidate these issues in the most transparent way, we apply a perturbation theory for 2D contact-interacting spin-$\frac{1}{2}$ fermions, starting from a microscopic action. We present what is apparently the first analytic calculation of the leading T dependence of the effective mass, Landau parameters, and response functions of a 2D electron gas, to second order in the interaction strength, *including all channels and all momentum* *processes (scales)*. We take into account the Ward identities explicitly.

The model.—We treat interacting fermions at finite temperature in the standard path integral Grassmannian formalism. The partition function is given by $\mathcal{Z} = \int \mathcal{D}\bar{\psi}\, \mathcal{D}\psi \exp(S_0 + S_{\text{int}})$, where

$$S_0 = \int_{(\mathbf{1})} \bar{\psi}_\alpha(\mathbf{1})[i\omega_1 + \mu - \epsilon(\mathbf{k}_1)]\psi_\alpha(\mathbf{1}). \qquad (1)$$

We have adopted the condensed notations, $(\mathbf{i}) \equiv (\mathbf{k}_i, \omega_i)$ and $\int_{(\mathbf{i})} \equiv \frac{1}{\beta}\int \frac{d\mathbf{k}_i}{(2\pi)^2}\sum_{\omega_i}$, where β is the inverse temperature, μ is the chemical potential, ω_i is the fermion Matsubara frequencies, and $\psi_\alpha(\mathbf{i})$ is a two-component Grassmann field with a spin index α. Summation over repeated indices is implicit throughout this paper. We set $k_B = 1$ and $\hbar = 1$. We consider mainly electrons with the bare spectrum of a free gas $\epsilon(\mathbf{k}) = k^2/2m$ and the circular Fermi surface, but discuss the consequences of generic spectra and a noncircular Fermi surface below. We take

$$S_{\text{int}} = -\frac{u}{4\nu_0}\int_{(\mathbf{1},\dots,\mathbf{4})} \bar{\psi}_\alpha(\mathbf{1})\bar{\psi}_\beta(\mathbf{2})\psi_\gamma(\mathbf{3})\psi_\varepsilon(\mathbf{4})T^{\alpha\beta}_{\gamma\varepsilon} \times \delta^{(2+1)}(\mathbf{1} + \mathbf{2} - \mathbf{3} - \mathbf{4}). \qquad (2)$$

Here $\nu_0 = m/2\pi$ is the free 2D density of states per spin, the spin antisymmetric operator $T^{\alpha\beta}_{\gamma\varepsilon} \equiv \delta_{\alpha\varepsilon}\delta_{\beta\gamma} - \delta_{\alpha\gamma}\delta_{\beta\varepsilon}$. For this model we calculate the FL parameters by assuming $0 < u < 1$, which corresponds to a weak repulsion, and by assuming that we are above the Kohn-Luttinger temperature.

Four-point 1PI vertex and the FL vertices.—The 1PI vertex $\hat{\Gamma}(\mathbf{1},\mathbf{2};\mathcal{Q})$ is defined in the standard way, where the transfer $\mathcal{Q} = (\mathbf{q}, \Omega)$, and Ω is a bosonic Matsubara frequency. To shorten notations we denote operators in spin space with a circumflex. To find the FL vertices we need to calculate $\hat{\Gamma}(\mathbf{1},\mathbf{2};\mathcal{Q})$ in the limit of zero transfer $\mathcal{Q}$ with incoming momenta lying on the Fermi surface. Since the Fermi surface is circular, the vertex can be parametrized by the relative angle between incoming momenta. It is known [3] that the limit $\mathcal{Q} \to 0$ is not unique. Two vertices, $\hat{\Gamma}^q(\theta_{12}) \equiv \hat{\Gamma}(\theta_{12}; \mathbf{q} \to 0, \Omega = 0)$ and $\hat{\Gamma}^\Omega(\theta_{12}) \equiv \hat{\Gamma}(\theta_{12}; \mathbf{q} = 0, \Omega \to 0)$, can be defined unambiguously at $\theta_{12} \neq 0$ and then continued to $\theta_{12} \to 0$ (see, e.g., Refs. [3,12] on this subtlety). They can be related to the components of the physical scattering vertex (A, B) and the Landau interaction function (F, G), respectively. Namely [3],

$$A\delta_{\alpha\gamma}\delta_{\beta\varepsilon} + B\sigma^a_{\alpha\gamma}\sigma^a_{\beta\varepsilon} = -2\nu_R Z^2 \Gamma^{\alpha\beta(q)}_{\gamma\varepsilon}, \qquad (3)$$

where $\nu_R = m^*/2\pi$, Z is the field renormalization constant, and $\hat{\sigma}^a$ are Pauli's matrices. Two components (F, G) of the Landau function are defined by an analogous equation, with the substitution $q \mapsto \Omega$, $A \mapsto F$, $B \mapsto G$.

Scattering vertex and Landau function.—The one-loop approximation for $\hat{\Gamma}(\mathbf{1},\mathbf{2};\mathcal{Q})$ in diagrammatic form is given in Fig. 1. In this approximation we calculate the FL vertices (scattering vertex, Landau function) using definition (3). At the one-loop level we can put $\nu_R = \nu_0$ $(m^* = m)$ and $Z = 1$. By doing the direct analytic evaluation of each diagram's contribution to the vertex we find the Fourier components of the angular-dependent FL vertices in terms of the temperature series. This series comes from the contributions of the ZS′ and Bardeen-Cooper-Schrieffer (BCS) loops. The details will be given in a companion paper [15]. We find the leading temperature corrections to the first two Fourier components of the vertices:

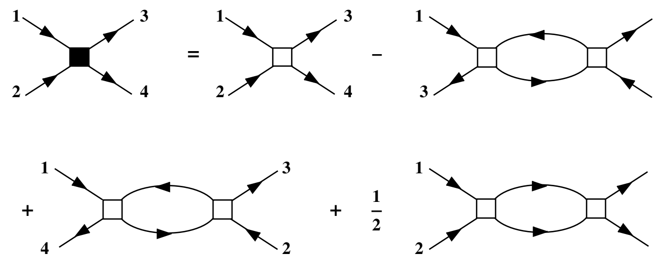

FIG. 1. Diagrammatic equation for the four-point vertex at the one-loop level. The one-loop graphs are called ZS, ZS′, and BCS in the order they appear on the rhs of this equation.

$$\delta A_0 = \delta F_0 = \delta A_1 = \delta F_1 = -u^2\frac{\pi^2}{24}\frac{T^2}{E_F^2}, \qquad (4)$$

$$\delta B_0 = \delta G_0 = -\delta B_1 = -\delta G_1 = -u^2\frac{T}{E_F}\ln 2. \qquad (5)$$

We should mention that, taken separately, each contribution of the ZS′ or BCS bubble gives a leading T term to the Fourier components of the vertices. However, a cancellation of such terms coming from two graphs occurs in the "charge sector" (i.e., in A, F components), while the linear-T terms survive in the "spin sector" (B, G components). The temperature dependence of the ZS′ contribution to the Fourier components of the vertices comes from integration around the "effective transfer" through the loop $|\mathbf{k}_1 - \mathbf{k}_2| \sim 2k_F$, i.e., when incoming momenta $\mathbf{k}_1 \sim -\mathbf{k}_2$. Similarly, the T dependence of the BCS contribution comes from regions of small $\mathbf{k}_1 + \mathbf{k}_2$, i.e., again when $\mathbf{k}_1 \sim -\mathbf{k}_2$. In other words, the temperature dependence comes from "$2k_F$ effects."

We expect that the cancellation of the T-linear terms in the charge sector of the vertices (4) is an artifact of our simple model calculation in which all three one-loop terms have the same factor u^2. If we had a bare coupling *function* of, say, two incoming momenta and transfer, then the coupling factors would have been different in each of the three graphs, and the anomalous T corrections would not have canceled.

The linearity of the leading T corrections to the vertices seems to be generic. The same T dependence (apart from presumably model-sensitive prefactors) was obtained in the previous RG analysis [12] of the *effective action* for 2D spinless fermions with a linearized one-particle spectrum and a momentum-dependent coupling function. According to Ref. [11] such temperature behavior can be understood from dimensional arguments.

Effective mass.— It is defined by the following equation:

$$\frac{m^*}{m} = \frac{1 - \frac{\partial \Sigma(\mathbf{1})}{\partial i\Omega}}{1 + \frac{m}{k_F}\frac{\partial \Sigma(\mathbf{1})}{\partial \mathbf{k}}\frac{\mathbf{k}_1}{k_F}}\Bigg|_{\mathbf{k}_1\in S_F}^{\omega_1=0}. \tag{6}$$

To second order we have

$$\frac{m^*}{m} = 1 - \left[\frac{\partial \Sigma(\mathbf{1})}{\partial \omega} + \frac{m}{k_F}\frac{\partial \Sigma(\mathbf{1})}{\partial \mathbf{k}}\frac{\mathbf{k}_1}{k_F}\right]\Bigg|_{\mathbf{k}_1\in S_F}^{\omega_1=0} + \mathcal{O}(u^3). \tag{7}$$

By using the two Ward identities following from the charge conservation and Galilean invariance [3], the above equation can be written as

$$\frac{m^*}{m} = 1 - \frac{1}{2}\int_{(\mathbf{2})} \frac{\mathbf{k}_1\mathbf{k}_2}{k_F^2}\,\Gamma^{\alpha\beta}_{\alpha\beta}(\mathbf{1},\mathbf{2};\Omega\to 0)\,\Delta(\mathbf{2})|^{\omega_1=0}_{\mathbf{k}_1\in S_F}, \tag{8}$$

where

$$\Delta(\mathbf{n}) = \frac{\beta}{4}\frac{\delta(\omega_n - \xi_{\mathbf{k}_n})}{\cosh^2(\frac{1}{2}\beta\xi_{\mathbf{k}_n})}. \tag{9}$$

Within our accuracy we can use the one-loop approximation for the vertex in Eq. (8). One can easily verify that in the the zero-temperature limit Eq. (8) recovers the standard result of the FL theory (FLT) [3], i.e., $m^*(T=0)/m = 1 + F_1(T=0)$. A straightforward extension of this relationship to finite temperatures like $m^*(T)/m = 1 + F_1(T)$ is not valid since, according to Eq. (8), $m^*(T)$ contains an extra contribution from the "off-shell" integration over k_2 ($\xi_{\mathbf{k}_2}$) normal to the Fermi surface, albeit the factor $\beta/\cosh^2(\beta\xi_{\mathbf{k}_2}/2)$ makes this contribution well localized near the Fermi surface. In other words the vertex entering the right-hand side (rhs) of Eq. (8) is not exactly the FL vertex $F(T)$ (up to the normalization factor) as it is defined in the FLT, since one of its momenta (namely, $\mathbf{k}_2$) is not confined to the Fermi surface. After calculations we find that the linear-temperature terms, coming essentially from two one-loop contributions (ZS$'$, BCS) to the vertex, cancel, resulting in

$$\frac{m^*}{m} = 1 + \frac{1}{2}u^2 + \mathcal{O}(u^2T^2). \tag{10}$$

In close analogy with the cancellation of the linear-temperature terms in the Fourier components of the FL vertices A (F), here the cancellation occurs between additive linear-T corrections coming from both "on-shell" (i.e., linear-T term coming from the $2k_F$ contribution to the vertex) and off-shell (i.e., the small-momentum contribution) integrations in two diagrams. Moreover, the on-shell (off-shell) T term of the ZS$'$ graph cancels the on-shell (off-shell) T term of the BCS graph, correspondingly. We expect that the cancellation does not occur at higher orders in the interaction. We also evaluated (7) for a generic 2D Fermi surface without explicitly using Ward identities, finding a T-linear term [15]. The result may be expressed as the sum of two terms, one arising from $2k_F$ processes and one from long-wavelength processes. The two contributions cancel for a circular Fermi surface and contact interaction, but not generically.

Calculations of the order u^2 term in the free energy for the model [(1) and (2)] show similarly that there is no T-linear term in the specific heat coefficient γ, contrary to the results of Ref. [6], where $2k_F$ contributions were missed. However, the model's modifications and/or higher orders in u destroy the cancellation of the linear term. This agrees with the recent experimental data on liquid ^{3}He films [7].

Response functions.— By using the same Ward identities as in the effective mass calculation, we found, for the dynamic zero-transfer limit ($\Omega = 0$, $\mathbf{q}\to 0$) of the density response function,

$$\begin{aligned}\varkappa &= \frac{m}{\pi}\{1 + u^2 + f_1 - f_0\},\\ f_1 &\equiv -\frac{\pi}{m}\int_{(\mathbf{1},\mathbf{2})}\frac{\mathbf{k}_1\mathbf{k}_2}{k_1^2}\,\Delta(\mathbf{1})\,\Gamma^{\alpha\beta}_{\alpha\beta}(\mathbf{1},\mathbf{2};\Omega\to 0)\,\Delta(\mathbf{2}),\\ f_0 &\equiv -\frac{\pi}{m}\int_{(\mathbf{1},\mathbf{2})}\Delta(\mathbf{1})\,\Gamma^{\alpha\beta}_{\alpha\beta}(\mathbf{1},\mathbf{2};\Omega\to 0)\,\Delta(\mathbf{2}).\end{aligned} \tag{11}$$

At $T=0$ we can read from Eq. (11) $\varkappa(T=0) = \frac{m}{\pi}(1 + F_0^2 + F_1 - F_0)$ which is nothing but the FLT result [3] $\varkappa^{\mathrm{FLT}} = \frac{m}{\pi}\frac{1+F_1}{1+F_0}$, expanded up to the third order over the interaction. Adding into consideration the Ward identity following from the total spin conservation, we derived, for the uniform spin susceptibility (for details see Ref. [15]),

$$\begin{aligned}\mathcal{X} &= \frac{m\mathfrak{g}^2}{4\pi}\{1 + u^2 + f_1 - g_0\},\\ g_0 &\equiv -\frac{\pi}{3m}\int_{(\mathbf{1},\mathbf{2})}\Delta(\mathbf{1})\,\sigma^a_{\gamma\varepsilon}\sigma^a_{\beta\alpha}\,\Gamma^{\alpha\varepsilon}_{\beta\gamma}(\mathbf{1},\mathbf{2};\Omega\to 0)\,\Delta(\mathbf{2}),\end{aligned} \tag{12}$$

where $\mathfrak{g}$ stands for the gyromagnetic ratio. Once again, one can see that in the zero-temperature limit the above equation gives $\mathcal{X}(T=0) = \frac{m\mathfrak{g}^2}{4\pi}(1 + G_0^2 + F_1 - G_0)$, reproducing thus the second-order expansion of the the FLT result [3] $\mathcal{X}^{\mathrm{FLT}} = \frac{m\mathfrak{g}^2}{4\pi}\frac{1+F_1}{1+G_0}$.

We were able to analytically calculate the integrals on the rhs of Eqs. (11) and (12) in the leading order of their temperature dependence. We found that the leading linear-T corrections, which can be traced back to the ZS$'$- and BCS-loop contributions to the vertex, cancel in each of the integral terms f_0 and f_1 in Eq. (11) separately. The result for the density response is

$$\varkappa = \frac{m}{\pi}\left(1 - u - \frac{1}{2}u^2 + u^2\ln\frac{2\Lambda}{k_F}\right) + \mathcal{O}(u^2T^2), \tag{13}$$

where $\Lambda \gg k_F$ is the ultraviolet cutoff we introduced to regularize the BCS loop. Note that the compressibility $K = \varkappa/n^2$, where n is the electron density [3].

We calculate the susceptibility in the same way. In this case the second integral term g_0 in Eq. (12) does not contain the contribution of the ZS$'$ loop, so the linear-T term coming from the BCS loop survives. Thus the spin susceptibility has a linear leading temperature correction:

$$\frac{\mathcal{X}(T)}{\mathcal{X}(0)} \approx 1 + u^2 \frac{T}{E_F}, \tag{14}$$

where $\mathcal{X}(0) = \frac{m\mathfrak{g}^2}{4\pi}[1 + u - u^2(\ln\frac{2\Lambda}{k_F} - \frac{3}{2})]$, and the first omitted term is $\mathcal{O}(u^2T^2)$.

It is useful to keep in mind that, albeit the response functions in Eqs. (11) and (12) are explicitly expressed in terms of the vertex only, those contributions indeed entangle both "purely vertex" and self-energy corrections. The latter are expressed just in terms of the vertex via the Ward identities.

Let us return to the argument for the cancellation of anomalous terms in the response functions due to Ward identities. We calculated the vertices at the one-loop level $\mathcal{O}(u^2)$. Through the Ward identities the self-energy corrections were taken into account with the same accuracy. There are no more terms of the order $\mathcal{O}(u^2)$ to cancel the temperature dependence (14). Thus, the linear-T dependence of susceptibility (or weaker T dependence of the compressibility) does not contradict the *exact* Ward identities known to us; moreover, in our results for the response functions, both vertex and self-energy corrections are included on the same footing *by using the Ward identities*.

Conclusions.—We have systematically examined the leading temperature corrections to FLT in 2D. Our results reveal the crucial importance of $2k_F$ processes. We find for the model of a 2D electron gas with a contact interaction that to order u^2 the leading T dependence of the FL parameters in the spin sector is T; for the others it is T^2.

We find that the relationships known from the classical FLT derivations at $T = 0$ for the parameters of Galilean-invariant FL (e.g, the effective mass, response functions vs components of the Landau function) are violated by finite-temperature terms. The coefficients in the temperature corrections to these relationships subtly involve contributions from small and large ($\sim 2k_F$) momentum processes.

The particularly interesting result we found is the leading linear-temperature dependence of the spin susceptibility (14). According to the perturbative calculations of Belitz *et al.* [11], the 2D FL susceptibility has a leading linear correction in $|\mathbf{q}|$ at $T = 0$ with a positive coefficient which is of second order in interaction, i.e., their result has the structure of Eq. (14). This also agrees with the phenomenology of Misawa [7] and the numerical results [13].

For more realistic models of electrons in (quasi-)2D crystals, i.e., for various tight-binding spectra and fillings, the free-gas-like square-root $2k_F$ singularities (with k_F depending on a chosen direction in $\mathbf{q}$ space) are known to exist in the Lindhard functions [16]. We think this is enough to result in linear-T terms in physical quantities analogous to what we found in this paper. We argue that the cancellation of the T terms in some FL parameters is special to second-order perturbation theory and the model considered, while the leading linear-T corrections are a generic feature of the 2D FL.

We hope our results may be experimentally tested in real 2D FL systems. For example, a very naive fit of the temperature dependence of the spin susceptibility in Sr_2RuO_4 system [17] when it is in the 2D metallic regime (above a 3D crossover temperature) shows that the data are compatible with the form (14). We expect that our results will stimulate a more detailed examination of the leading temperature dependences of response functions in 2D systems.

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