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Quantum magneto-oscillations in a two-dimensional Fermi liquid

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Quantum magneto-oscillations provide a powerful tool for quantifying Fermi-liquid parameters of metals. In particular, the quasiparticle effective mass and spin susceptibility are extracted from the experiment using the Lifshitz-Kosevich formula, derived under the assumption that the properties of the system in a nonzero magnetic field are determined uniquely by the zero-field Fermi-liquid state. This assumption is valid in three dimensions (3D) but, generally speaking, erroneous in 2D where the Lifshitz-Kosevich formula may be applied only if the oscillations are strongly damped by thermal smearing and disorder. In this work, the effects of interactions and disorder on the amplitude of magneto-oscillations in 2D are studied. It is found that the effective mass diverges logarithmically with decreasing temperature signaling a deviation from the Fermi-liquid behavior. It is also shown that the quasiparticle lifetime due to inelastic interactions does not enter the oscillation amplitude, although these interactions do renormalize the effective mass. This result provides a generalization of the Fowler-Prange theorem formulated originally for the electron-phonon interaction.

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Patterns of quantum magneto-oscillations in thermodynamic (de Haas–van Alphen effect) and transport (Shubnikov–de Haas effect) quantities encode three important parameters of a Fermi-liquid (FL) metal. The period of the oscillations gives the area of the extremal cross section of the Fermi surface, the slope of the temperature dependence of the oscillations amplitude provides the quasiparticle effective mass, and the phase shift between oscillations of spin-up and spin-down electrons yields the (renormalized) spin susceptibility. Magneto-oscillations studies of the FL state in two dimensions (2D) date back to early 1970's, when semiconductor heterostructures first became available.[1] Another surge of the activity in this field, which occurred in mid 1990's, was stimulated by the discovery of the metallic state at $\nu = 1/2$.[2] Recently, Shubnikov–de Haas oscillations have been used to determine the parameters of the ''anomalous'' metallic state in Si metal-oxide-semiconductor field-effect transistor (MOSFET's) and other semiconductor heterostructures exhibiting an apparent metal-insulator transition in zero magnetic field.[3–6] Despite the long and successful history of quantifying FL's in 3D via magneto-oscillations, this method remains controversial in 2D. The primary goal of our paper is to resolve some of the open issues.

The first controversy is related to the applicability of the current theory of magneto-oscillations to the two-dimensional case. The analysis of the experimental data in 2D is often based on the premise that the classic result for magneto-oscillations in a Fermi liquid for the three-dimensional case, known as the ''Lifshitz-Kosevich (LK) formula,''[7–9] is transferrable to 2D upon a trivial change in the electron spectrum. The crucial features of the LK formula, i.e., its validity for arbitrarily strong interactions (without destroying the Fermi liquid) and the fact that the FL parameters entering the formula are taken at zero magnetic field, survive on this premise. The deviations of the observed oscillation pattern in stronger fields from that predicted by the LK formula are ascribed to oscillations in the effective g factor[1] and the effective mass.[10] On the other hand, there have been warnings that the three-dimensional LK formula is nontransferrable to 2D (Refs. 11,12) for any field strength. Hence the situation needs to be clarified. The second controversy—not specific to 2D—is related to the effect of quasiparticle damping. It is often mentioned in the literature that any scattering of quasiparticles, elastic and inelastic, contributes to the smearing of magneto-oscillations via the effective Dingle temperature (scattering rate) for a given process. Alternatively, Fowler and Prange[13] showed that the electron-phonon scattering rate does not appear in the oscillations amplitude due to the cancellation of two T-dependent parts of the Matsubara self-energy (cf. also Ref. 14). To the best of our knowledge, this cancellation has never been discussed for other interactions, including the electron-electron one, which is of a primary importance for two-dimensional electron systems. The last issue to be addressed in this paper (and not discussed previously in the literature) is the effect of interference between electron-impurity and electron-electron scattering on magneto-oscillations, neglected in the LK formula. The theory of interference effects in the ballistic regime,[15] when $T\tau \gg 1$, where τ is the electron-impurity scattering time (we set $\hbar = k_B = 1$ throughout the paper), offers a plausible explanation of the metallic temperature dependence in the metallic phase of the two-dimensional metal-insulator transition. Unusual (within the LK framework) temperature dependences of the oscillation amplitude are also commonly observed in Si MOSFETs,[6,16] but the proper theory is currently lacking.

Our answers to these open questions are as follows. (i) Although it is true that the LK formula does not work in 2D at $T = 0$ and in the absence of disorder, it is still applicable to the situation when finite temperature and/or disorder cause the oscillations to be exponentially small. (ii) The cancellation of the scattering rate term in the Matsubara self-energy is pertinent to *any* inelastic interaction, including the electron-electron one. Due to this cancellation, the scattering rate of inelastic processes does not enter the oscillation amplitude. (iii) Interference between electron-impurity and electron-electron interactions gives a new $T \ln T$ dependence of the amplitude's argument, that can be interpreted equiva-

lently either as a ''T-dependent'' effective mass or Dingle temperature. The functional form of this dependence is the *same* in the diffusive ($T\tau \ll 1$) and ballistic ($T\tau \gg 1$) regimes.

We limit our analysis to the de Haas–van Alphen effect and assume the chemical potential is fixed. The main features of the results for the de Haas–van Alphen effect, in particular, the T dependence of the oscillation amplitude, are commonly expected to apply to the Shubnikov–de Haas effect as well, although a rigorous proof of that is currently lacking. Assuming a fixed chemical potential is not essential for the case of small oscillations, which is the focus of this paper (see below). We begin with a brief reminder of how the LK formula is derived in the Luttinger formalism.[8,9] The key issue here is whether the zero-field FL parameters determine uniquely the oscillation pattern in a finite (and not small) field. The (Matsubara) self-energy (that encodes all FL parameters) consists of two parts $\Sigma=\Sigma_0+\Sigma_{\rm osc}$, where Σ_0 may contain a monotonic (nonoscillatory) dependence on magnetic field and $\Sigma_{\rm osc}$ oscillates with the field. For electron-electron interactions in 3D $|\Sigma_{\rm osc}|/|\Sigma_0| \sim N^{-3/2}$, where N is the number of occupied Landau levels, whereas the leading term in the oscillatory part of the thermodynamic potential Ω falls off as $N^{-5/2}$ for $\Sigma_{\rm osc}=0$. Expanding Ω in $\Sigma_{\rm osc}$ up to the second order (the first order term vanishes due to the property $\delta\Omega/\delta\Sigma=0$), one finds that the oscillatory part of Σ can always be neglected in the semiclassical regime ($N\gg 1$). With this simplification and for a momentum-independent self-energy, the amplitude of the kth harmonic in Ω is given by

$$A_k=\frac{4\pi^2 kT}{\omega_c}\sum_{\varepsilon_n>0}\exp\left(-\frac{2\pi k[\varepsilon_n+i\Sigma_0(i\varepsilon_n,T)]}{\omega_c}\right), \quad (1)$$

where $\varepsilon_n=\pi(2n+1)T$ and $\omega_c=eB/mc$. Notice that $i\Sigma_0$ is real and does not contain a constant term. The second argument of Σ_0 emphasizes the fact that the temperature enters Σ_0 in two ways: via the Matsubara frequency and via the thermal distribution of electrons and other degrees of freedom. For a generic Fermi liquid and in the presence of short-range impurities $i\Sigma_0(i\varepsilon_n,T)=\alpha\varepsilon_n+{\rm sgn}\varepsilon_n/2\tau$, so that the effective mass is defined as $m^*=m(1+\alpha)$. The amplitude then assumes a familiar form

$$A_k=\frac{2\pi^2 kT/\omega_c}{\sinh(2\pi^2 kT/\omega_c^*)}\exp\left(-\frac{2\pi^2 T_D}{\omega_c}\right), \quad (2)$$

where $\omega_c^*=eB/m^*c$, and $T_D=1/2\pi\tau$ is the Dingle temperature. Momentum dependence of Σ_0 of the form $\beta v_F(p-p_F)$ results in a change of the effective mass in Eq. (2) to $m^*=m(1+\alpha)(1+\beta)^{-1}$ and in multiplying Eq. (2) by $Z_s m^*/m$, where Z_s is the renormalization factor.

In arbitrary dimensionality D, the estimates for the ratio of oscillatory to monotonic-in-field parts of the self-energy and for the leading oscillatory term in Ω change to $N^{-D/2}$ and $N^{-(D+2)/2}$, respectively. For $D=2$, the oscillations in the self-energy are as important as in the thermodynamic potential itself.[12] The Luttinger expansion at $T=T_D=0$ breaks down and the LK formula is not, generally speaking, valid.[12] The physical reason is that the ground states of an interacting system at $B=0$ and in a finite field are not adiabatically connected in 2D. This fact has been emphasized by recent findings that the ground state of a two-dimensional electron liquid is not a Fermi liquid even for $N\gg 1$, but rather a charge-ordered state.[17] Nevertheless, an absence of the full LK formula in 2D does not preclude a canonical analysis of magneto-oscillations, if under more restrictive conditions, as the FL behavior is restored at higher energies.

The power-counting argument for (against) neglecting the oscillatory part of the self-energy in 3D (2D) is valid at $T=T_D=0$. If the real and/or Dingle temperatures are sufficiently high, i.e.,

$$2\pi^2(T/\omega_c^*, T_D/\omega_c)\gtrsim 1, \quad (3)$$

the amplitudes of *all* oscillatory quantities, including the self-energy, are exponentially small. Neglecting the oscillatory part of the self-energy, the amplitude of the first harmonic takes the form

$$A_1=(4\pi^2 T/\omega_c)\exp\{-2\pi[\pi T+i\Sigma_0(i\pi T,T)]/\omega_c\}, \quad (4)$$

where we assumed that T is high [in the sense of condition (3)] and limited the Matsubara sum by the first term $\varepsilon_0=\pi T$. The oscillatory part of Σ results in a correction to A_1 which is itself of order A_1 (with exponential accuracy). The net contribution to Ω is of order A_1^2, which is of the same order as A_2 for $\Sigma_{\rm osc}=0$. Thus, harmonics with $k\geqslant 2$ are affected by the oscillations in Σ and a two-dimensional analog of the LK formula, which includes the sum over all k, can only be derived in a perturbation theory for a specific interaction but not for a generic Fermi liquid. However, the $k=1$ harmonic does not include $\Sigma_{\rm osc}$ and, as long as Eq. (3) is satisfied, the analysis can proceed as in the three-dimensional case. In what follows, we assume that Eq. (3) is satisfied and the amplitude of the first (and only important) harmonic is given by Eq. (4).

Next we discuss whether the quasiparticle relaxation rate affects the amplitude of magneto-oscillations. We set $T_D=0$ temporarily. Suppose that a quasiparticle relaxation rate is measured in a clean Fermi liquid, e.g., via electron heat conductivity, with a result that $1/\tau_{\rm e\text{-}e}\propto T^2$. It seems natural to assume that the same rate contributes also to the Dingle temperature of magneto-oscillations. That this is *not* the case was shown for the electron-phonon interaction by Fowler and Prange.[13] Here we generalize their arguments for the electron-electron interaction in 3D and 2D, and then give a general result for an arbitrary interaction. For a generic Fermi liquid in 3D, the Matsubara self-energy, up to the quadratic in ε_n and T terms, can be written as

$$i\Sigma_0(i\varepsilon_n,T)=\alpha\varepsilon_n+i\beta v_F(p-p_F)+\gamma[(\pi T)^2-\varepsilon_n^2]. \quad (5)$$

In addition to a direct calculation, the validity of the quadratic term in Eq. (5) is readily established by noticing that upon analytic continuation $i\varepsilon_n\to\varepsilon+i0$ this term gives the correct form for the imaginary part of the on-shell self-energy: $-{\rm Im}\Sigma_0^R\propto(\pi T)^2+\varepsilon^2$.[18] The amplitude of the first harmonic (4) contains $i\Sigma_0(i\pi T,T)$, in which the quadratic term vanishes identically. The T^2 terms from higher Matsub-

ara frequencies (legitimately considered within this scheme in 3D) increase the amplitude and cannot be interpreted as "damping."

In 2D, the integral over momentum transfers diverges logarithmically at the lower limit, changing the behavior of $\mathrm{Im}\Sigma_0^R$ to $E^2\ln E$, where $E=\max\{\varepsilon,T\}$. This change does not alter the principal result. Consider the simplest case of a contact interaction. To the second order in this interaction, the quadratic term in Eq. (5) is replaced by

$$i\tilde{\Sigma}_0(i\varepsilon_n,T)=-\frac{U^2m}{\pi^2v_F^2}T\sum_{\omega_m=0}^{\varepsilon_n-\pi T}\omega_m\ln(\varepsilon_F/\omega_m). \qquad (6)$$

Although the sum in Eq. (6) does not have an analytic solution, it obviously vanishes for $\varepsilon_n=\pi T$.

To analyze a general case of a finite range and dynamic interaction, including the screened Coulomb one, it is convenient to find the imaginary part of the retarded self-energy first and then continue back to Matsubara frequencies. On the mass shell, $\mathrm{Im}\Sigma_0^R$ is given by

$$\mathrm{Im}\Sigma_0^R(\varepsilon)=-\frac{\pi}{2(2\pi)^D}\int d\omega F(\omega)\left[\coth\frac{\omega}{2T}-\tanh\frac{\omega-\varepsilon}{2T}\right],$$

where $F(\omega)=\int d^Dq\,\delta(\omega-\mathbf{v}_F\cdot\mathbf{q})\mathrm{Im}V^R(\omega,q)$ and $V^R(\omega,q)$ is the retarded interaction potential. As a function of a complex variable z, $f(z)\equiv\mathrm{Im}\Sigma^R(z)$ has the following properties in the upper half-plane: (i) all lines $\mathrm{Im}z=\pi(2n+1)T$ are branch cuts on which $\mathrm{Re}\,f$ is continuous but $\mathrm{Im}\,f$ changes jumpwise; (ii) due to the fact that $\tanh[x-i\pi(n+1/2)]=\coth x$, all points $z=i\pi(2n+1)T$ are zeroes of $f(z)$. Thus, function $f(z)$ is analytic in the band $0\leqslant\mathrm{Im}z<\pi T$ including the point $z=i\pi T$. Analytic continuation from the real axis into this band is legitimate and at $z=i\pi T$ it yields the Matsubara self-energy $\tilde{\Sigma}_0(i\varepsilon_0=i\pi T,T)$, which is equal to zero. Zeroes of $f(z)$ at $z=i\pi(2n+1)T$ with $n\geqslant1$ do not lead to vanishing of $\tilde{\Sigma}_0(i\varepsilon_n,T)$ for $n\geqslant1$ because those zeroes are separated from the real axis by branch cuts and thus are not accessible by analytic continuation. The two-dimensional case is special only in that the q integration results in the $\ln\omega$ factor in $F(\omega)$ which does not change the reasoning given above. In particular, for a dynamically screened Coulomb interaction in 2D, $F(\omega)\propto\omega\ln|\omega|$ and still $\tilde{\Sigma}_0(i\pi T,T)=0$. As this result does not depend on the particular form of the interaction, it can be viewed as a generalization of the Fowler-Prange theorem. Note that in 3D the Fowler-Prange theorem is of limited applicability because nothing prevents one from considering $k>1$ and lower values of $T+T_D$, when the effect of $\Sigma_0(i\varepsilon_{n>0},T)$ needs to be taken into account. In 2D, one is bound to consider only $\Sigma_0(i\varepsilon_{n=0},T)$ within the Luttinger approximation.

Finally, we discuss the effect of interference between electron-electron and electron-impurity scattering on magneto-oscillations, extending the analysis of the interference corrections to the self-energy in 2D from the diffusive ($T\tau\ll1$) (Ref. 19) to the ballistic ($T\tau\gg1$) limit. The general form of the interference correction to the Matsubara self-energy is (see Fig. 1)

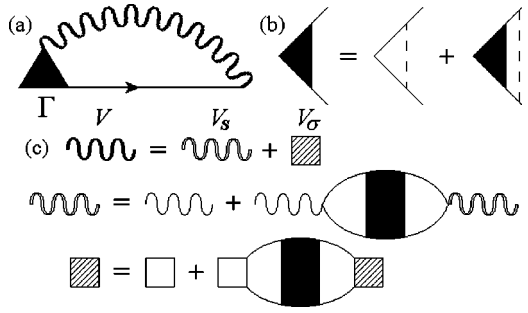

FIG. 1. (a) The interference correction to the self-energy. (b) The vertex correction is assigned to either one of the vertices in (a) because the self-energy arises as insertions into the thermodynamic potential (closed loops). (c) Singlet/triplet-channel contributions to the effective potential.

$$\Sigma_0^{\mathrm{int}}(i\varepsilon_n,\mathbf{p})=-2T\sum_{\varepsilon_n(\omega_m-\varepsilon_n)>0}\int\frac{d^2q}{(2\pi)^2}V(i\omega_m,q)\times\Gamma(i\omega_m,q)G(i\varepsilon_n-i\omega_m,\mathbf{p}-\mathbf{q}), \qquad (7)$$

where $G(i\varepsilon_n,\mathbf{p})=(i\varepsilon_n-\xi_p+i\mathrm{sgn}\varepsilon_n/2\tau)^{-1}$ and the effective interaction $V=V_s+3V_t$ contains the contributions from both singlet and triplet channels,[15]

$$V_s^{-1}(i\omega_m,q)=(2\pi e^2/q+F_\rho^0/\nu)^{-1}-\Pi(i\omega_m,q),$$

$$V_\sigma^{-1}(i\omega_m,q)=\nu/F_\sigma^0-\Pi(i\omega_m,q),$$

where $\nu=m/\pi$. The factor of two in Eq. (7) accounts for two possibilities of including the vertex correction in Fig. 1(a). The Fermi-liquid constants F_ρ^0 and F_σ^0 (Ref. 20) determine the renormalized charge and spin susceptibilities, respectively. The vertex $\Gamma(i\omega_m,q)=[\sqrt{(|\omega_m|\tau+1)^2+(qv_F\tau)^2}-1]^{-1}$ reduces to $\Gamma=(D\tau q^2+|\omega_m|\tau)^{-1}$ and $\Gamma=\tau^{-1}[|\omega_m|^2+(qv_F)^2]^{-1/2}$ in the diffusive and ballistic limits, respectively. The general form of the (small q) polarizarion operator

$$\Pi(i\omega_m,q)=-\nu[1-|\omega_m|\tau\Gamma(i\omega_m,q)] \qquad (8)$$

reduces to $\Pi(i\omega_m,q)=-\nu Dq^2/(Dq^2+|\omega_m|)$ in the diffusive limit, where $D=v_F^2\tau/2$, and to $\Pi^0(i\omega_m,q)=-\nu[1-|\omega_m|/\sqrt{(v_Fq)^2+\omega_m^2}]$ in the ballistic one. Omitting the details of lengthy but straightforward calculations, we give just the result for the self-energy valid to logarithmic accuracy

$$i\Sigma_0^{\mathrm{int}}(i\pi T,T)=-T\ln(\varepsilon_F/T)Q(T\tau,F_\sigma^0)/2\varepsilon_F\tau,$$

where $Q(T\tau,F_\sigma^0)=g_\rho(T\tau)+[3F_\sigma^0/(1+F_\sigma^0)]g_\sigma(T\tau)$ and $g_{\rho/\sigma}(x)$ are slowly varying functions which interpolate between the diffusive and ballistic regimes. The limiting values of $g_{\rho/\sigma}(x)$ are as follows: $g_\rho(0)=1, g_\rho(x\gg1)=3/2, g_\sigma(0)=1, g_\sigma(x\gg1)=1/2$. Apart from the numerical coefficients, the T dependence of $i\Sigma_0^{\mathrm{int}}(\pi T,T)$ is the same in the diffusive and ballistic regimes. In that sense, the behavior of the self-energy is similar to that of the tunneling density of states.[21] Notice that the interference correction to the *scattering rate* in the ballistic regime $|\mathrm{Im}\Sigma_0^R|_{\mathrm{int}}\sim(T/\varepsilon_F\tau)\ln(\varepsilon_F/T)$ is smaller than the scattering rate in a clean FL, $|\mathrm{Im}\Sigma_0^R|_{\mathrm{int}}$

$\sim(T^2/\varepsilon_F)\ln(\varepsilon_F/T)$, in the parameter $(T\tau)^{-1}\ll 1$. However, due to the cancellation of the $T^2\ln T$ term in $i\Sigma_0(i\pi T,T)$ discussed earlier in this paper, the interference correction is the main nonlinear T-dependent term in the Matsubara self-energy, leading to a modification of the LK formula. The $T\ln T$ dependence of the self-energy can be interpreted as a logarithmic T dependence of the effective mass. Following this interpretation, the interference effect leads to a replacement of the effective mass in the argument of the exponential in A_1 by

$$m^*(T)=m^*\left(1-\frac{m}{m^*}\frac{\ln(\varepsilon_F/T)}{2\pi\varepsilon_F\tau}Q(T\tau,F_\sigma^0)\right). \qquad (9)$$

Note that the effective mass is reduced by the singlet-channel interaction but enhanced by the ferromagnetic $(F_\sigma^0<0)$ interaction in the triplet channel. The $\ln T$ dependence of the effective mass is a characteristic feature of the ''marginal Fermi-liquid'' model.[22,23]

Equivalently, the nonlinear T dependence of $i\Sigma_0(i\pi T,T)$ may be interpreted as a T dependent Dingle temperature:

$$T_D(T)=T_D[1-(T/\varepsilon_F)\ln(\varepsilon_F/T)Q(T\tau,F_\sigma^0)]. \qquad (10)$$

One of the empirical procedures used in Ref. 6 to account for the observed deviations from the LK formula was to assume that the effective Dingle temperature has the same T dependence as the zero-field resistivity. Comparing Eq. (10) with the result for the interference correction to the resistivity,[15] we see that, although this recipe is not precise, it has some theoretical justification: the general structure of the $T_D(T)$ and $\rho(T)$ depedencies is similar in the ballistic regime, except for a factor of $\ln(\varepsilon_F/T)$ present in $T_D(T)$ but not in $\rho(T)$.

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