Interaction-Induced Oscillations of the Tunneling Density of States in a Nonquantizing Magnetic Field

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(Received 9 August 1996)

We study tunneling into an interacting disordered two-dimensional electron gas in a nonquantizing magnetic field, which does not cause the standard de Haas–van Alphen oscillations. Interaction induces a new type of oscillation in the tunneling density of states with the characteristic period of cyclotron quantum $\hbar\omega_c$. [S0031-9007(96)02182-5]

PACS numbers: 73.40.Gk, 71.10.Pm

It is well known that a strong magnetic field modifies the single-particle density of states (DOS) of noninteracting electrons due to the Landau quantization. In a two-dimensional electron gas, the quantization leads to a peak structure in the DOS, which is revealed in tunneling experiments as peaks in the dependence of the tunneling conductance on applied bias, see, e.g., Ref. [1]. The form and the width of these peaks are determined [2] by the disorder. Experiments [1,3,4] show also suppression of the conductance at zero bias. This suppression is a manifestation of interaction between electrons and has been studied theoretically both for disordered [5] and for clean [6] two-dimensional conductors in the quantizing magnetic field.

In a weak magnetic field the distance between the Landau levels, $\hbar\omega_c$, is smaller than their disorder-induced width. In such "classical" magnetic field, the Landau quantization peaks in DOS disappear, while the aforementioned interaction correction to tunneling DOS survives and in the limit of zero field evolves into a singular at the Fermi energy negative logarithmical correction predicted originally by Altshuler, Aronov, and Lee [7].

The question arises whether or not in such "classical," $\omega_c\tau_s \ll 1$, limit the magnetic field influences the spatially averaged density of electron states measured in tunneling experiments (here τ_s is the electron quantum lifetime). For the noninteracting system, the effect of the weak magnetic field on DOS is exponentially small [2], $\propto \exp[-2\pi/(\omega_c\tau_s)]$, and can be neglected. The goal of the present paper is to show that, contrary to the noninteracting case, in the *interacting* electron gas the "classical" magnetic field does produce a significant effect on tunneling DOS. This effect becomes pronounced if the disorder potential is weak enough and smooth, with the correlation length much larger than the Fermi wavelength and the amplitude much smaller than the Fermi energy. In such a potential, electrons experience small-angle scattering, and their transport relaxation time $\tau_{\rm tr}$, is much larger than τ_s. Thus there exists a range of magnetic fields in which Landau quantization is suppressed ($\omega_c\tau_s \ll 1$), while classical electron trajectories are strongly affected by the field ($\omega_c\tau_{\rm tr} \gg 1$). In this regime interaction correction to the tunneling DOS, $\delta\nu(\epsilon)$, is strongly enhanced with respect to the zero magnetic field case. Furthermore, we will show that it exhibits peaks as a function of energy with the distance between peaks equal to the cyclotron quantum, $\hbar\omega_c$. The further a peak is away from the Fermi level, the smaller and wider it is. The shape of the nth peak in DOS, $|\epsilon - n\hbar\omega_c| \lesssim \hbar\omega_c/2$, is given by:

$$\frac{\delta\nu(\epsilon)}{\nu} = \frac{\hbar(\omega_c\tau_{\rm tr})^2}{8\pi\epsilon_F\tau_{\rm tr}}\,\frac{1}{n}\,f\left(\frac{\epsilon - n\hbar\omega_c}{\hbar n^2/\tau_{\rm tr}}\right), \qquad (1)$$

where

$$f(x) = \frac{1}{\sqrt{2}}\left[\frac{1+\sqrt{x^2+1}}{x^2+1}\right]^{1/2}, \qquad (2)$$

energy ϵ is measured from the Fermi level, and $\nu = m/\pi\hbar^2$ is the free-electron density of states (m is the electron mass). The peaks overlap strongly for $\epsilon \gtrsim \hbar\omega_c\sqrt{\omega_c\tau_{\rm tr}}$, and the oscillatory structure is washed out.

Sensitivity of tunneling DOS to the classical magnetic field comes from the fact that, as we will show, the interaction correction to tunneling DOS is associated with the self-crossing of classical electron trajectories. We denote the probability for an electron to complete a loop of self-crossing trajectory over time t as $K(t)$. The interaction correction to DOS, $\delta\nu(\epsilon)$, turns out to be related to the Fourier transform of this probability, $\delta\nu(\epsilon) \propto K(\epsilon) = \int_0^\infty dt e^{-i\epsilon t}K(t)$. The strong enough, $\omega_c\tau_{\rm tr} \gg 1$, magnetic field curves the electron trajectories, significantly affects the return probability, and, in turn, affects the tunneling DOS.

For long time scales $t \gg \tau_{\rm tr}$, the function $K(t)$ can be found from the diffusion equation. It gives $K(t) \propto (Dt)^{-1}$ for the two-dimensional case (D is the diffusion coefficient). The Fourier transform, $K(\epsilon)$, is proportional to $\ln(\epsilon)$, which leads to a predicted by Altshuler, Aronov, and Lee [7] logarithmic correction to DOS at small energies, $\epsilon \ll \hbar/\tau_{\rm tr}$, with the renormalized by the magnetic field diffusion coefficient [8].

At short time scales, $t \ll \tau_{\rm tr}$, electrons move ballistically along the cyclotron orbits. Provided that $\omega_c\tau_{\rm tr} \gg 1$, during the time t electron may return to the initial point many times. Thus, at these short time scales the magnetic field significantly increases the return probability $K(t)$.

Multiple periodic returns of the electron produce peaks in the probability Fourier transform $K(\epsilon)$ at energies, which are multiples of the cyclotron quantum. Tunneling density of states oscillates with the same period, which is reflected by Eq. (1).

Now we derive the expression for the interaction correction to DOS valid for arbitrary energies. We will put $\hbar = 1$ in all intermediate formulas. In an ideal pure sample the electron density does not depend on coordinates, $n_e(\mathbf{r}) = n_0$. Due to the sharp Fermi edge, scattering of the electron forming the Fermi sea on impurities results in an interference pattern in electron density. This pattern is commonly referred to as the Friedel oscillation [9]. In general, one can express the density profile of noninteracting electrons in terms of the exact retarded Green function, $\mathcal{G}^R_\epsilon(\mathbf{r},\mathbf{r}') = \sum_l \psi_l^*(\mathbf{r}')\psi_l(\mathbf{r})/(\epsilon - \epsilon_l + i0)$, of an electron in the random potential:

$$n_e(\mathbf{r}) = -\frac{2}{\pi}\int_{-\epsilon_F}^{0} d\epsilon\, \mathrm{Im}\,\mathcal{G}^R_\epsilon(\mathbf{r},\mathbf{r})\,. \tag{3}$$

Single electron wave function $\psi_l(\mathbf{r})$ satisfies the Schrödinger equation for noninteracting electrons, $\hat{H}_0\psi_l = (\epsilon_l + E_F)\psi_l$, where $\hat{H}_0 = -(\hbar^2/2m)\nabla^2 + U_r(\mathbf{r})$, E_F is the Fermi energy, and $U_r(\mathbf{r})$ is the random potential. In the presence of interaction $V(\mathbf{r}-\mathbf{r}')$ between electrons, the Friedel oscillation produces an additional term in the Hamiltonian, $\hat{H}_{HF}$, which can be presented (see, e.g., Ref. [9]) as a sum of Hartree, V_H, and exchange, V_F, terms:

$$H_{HF}(\mathbf{r},\mathbf{r}') = V_H(\mathbf{r})\,\delta(\mathbf{r}-\mathbf{r}') - V_F(\mathbf{r},\mathbf{r}') \tag{4}$$

$$V_H(\mathbf{r}) = \int V(\mathbf{r}-\mathbf{r}'')\delta n_e(\mathbf{r}'')d\mathbf{r}'', \tag{5}$$

$$V_F(\mathbf{r},\mathbf{r}') = \frac{1}{2}\,V(\mathbf{r}-\mathbf{r}')\delta\rho(\mathbf{r},\mathbf{r}'). \tag{6}$$

Here $\delta\rho(\mathbf{r},\mathbf{r}')$ is the perturbation of the density matrix, $\rho(\mathbf{r},\mathbf{r}') = -(2/\pi)\int d\epsilon\,\mathrm{Im}\,\mathcal{G}^R_\epsilon(\mathbf{r},\mathbf{r}')$, by the random potential. Only the electrons with the same spin participate in the exchange interaction, which is reflected by the factor $1/2$ in Eq. (6). The Hartree-Fock energy (5)–(6) oscillates as a function of coordinate in the same manner as $\delta n_e(\mathbf{r})$ does.

The local DOS is related to the retarded Green function of electron, $\nu(\epsilon,\mathbf{r}) = -(2/\pi)\,\mathrm{Im}\,\mathcal{G}^R_\epsilon(\mathbf{r},\mathbf{r})$. Scattering of electron on the Hartree-Fock potential, Eqs. (5)–(6), induces a correction to $\mathcal{G}^R_\epsilon(\mathbf{r},\mathbf{r})$, which, in the Born approximation, can be expressed as:

$$\delta\mathcal{G}^R_\epsilon(\mathbf{r},\mathbf{r}) = \int \mathcal{G}^R_\epsilon(\mathbf{r},\mathbf{r}')V_H(\mathbf{r}')\mathcal{G}^R_\epsilon(\mathbf{r}',\mathbf{r})d\mathbf{r}' - \int \mathcal{G}^R_\epsilon(\mathbf{r},\mathbf{r}')V_F(\mathbf{r}',\mathbf{r}'')\mathcal{G}^R_\epsilon(\mathbf{r}'',\mathbf{r})d\mathbf{r}'d\mathbf{r}''. \tag{7}$$

We will be interested in the spatially averaged density of states, $\delta\nu(\epsilon) = (1/S)\int\delta\nu(\epsilon,\mathbf{r})d\mathbf{r}$, where S is the area of the system. For simplicity we will start with the case of the finite-range interaction potential and will calculate the Hartree contribution to the averaged DOS. Making use of Eqs. (3), (5), and (7), and exploiting the identity, $\int \mathcal{G}^R_\epsilon(\mathbf{r}',\mathbf{r})\,\mathcal{G}^R_\epsilon(\mathbf{r},\mathbf{r}'')d\mathbf{r} = \partial\mathcal{G}^R_\epsilon(\mathbf{r}',\mathbf{r}'')/\partial\epsilon$, we obtain:

$$\delta\nu_H(\epsilon) = \frac{2}{\pi^2 S}\mathrm{Re}\int_{-\epsilon_F}^{0} d\epsilon_1 \int d\mathbf{r}'d\mathbf{r}''V(\mathbf{r}'-\mathbf{r}'') \times \frac{\partial\mathcal{G}^R_\epsilon(\mathbf{r}',\mathbf{r}')}{\partial\epsilon}\left[\mathcal{G}^R_{\epsilon_1}(\mathbf{r}'',\mathbf{r}'') - \mathcal{G}^A_{\epsilon_1}(\mathbf{r}'',\mathbf{r}'')\right], \tag{8}$$

where $\mathcal{G}^A_\epsilon(\mathbf{r},\mathbf{r}') = [\mathcal{G}^R_\epsilon(\mathbf{r}',\mathbf{r})]^*$. We are interested in the correction to the density of states averaged over the realizations of the disorder potential. Average of the product of two retarded Green functions $\mathcal{G}^R_\epsilon(\mathbf{r}',\mathbf{r}')\mathcal{G}^R_{\epsilon_1}(\mathbf{r}'',\mathbf{r}'')$ does not contain contributions associated with the electron trajectories longer than λ_F. Thus, this product does not produce DOS energy dependence at small, as compared to E_F, energies, and can be neglected. On the contrary, averaged product $\mathcal{G}^R\mathcal{G}^A$ is determined by long electron trajectories. Furthermore, it can be expressed [10] in terms of the classical probability density $\mathcal{D}$:

$$\langle\mathcal{G}^R_{\epsilon_1}(\mathbf{r}_1,\mathbf{r}_2)\mathcal{G}^A_{\epsilon_2}(\mathbf{r}_3,\mathbf{r}_4)\rangle = \pi\nu\int\frac{d\phi_1}{2\pi}\int\frac{d\phi_2}{2\pi} \times e^{i\mathbf{p}_1(\mathbf{r}_1-\mathbf{r}_4)}e^{i\mathbf{p}_2(\mathbf{r}_3-\mathbf{r}_2)}\mathcal{D}(\epsilon_1-\epsilon_2;\mathbf{r}_1,\phi_1;\mathbf{r}_2,\phi_2)\,;$$

$$\mathcal{D}(\omega;1;2) = \int_0^\infty dt e^{i\omega t}\mathcal{D}(t;1;2)\,. \tag{9}$$

Here $\mathbf{p}_i = p_F(\cos\phi_i,\sin\phi_i)$, and $\mathcal{D}(t;1;2)$ is the probability density for electron which starts at moment $t=0$ in point $\mathbf{r}_1$ with the direction of momentum ϕ_1 to arrive at moment t to the point $\mathbf{r}_2$ with momentum direction ϕ_2.

Equation (9) is valid as long as the sizes $|\mathbf{r}_1-\mathbf{r}_4|$ and $|\mathbf{r}_2-\mathbf{r}_3|$ of spatial domains defining the ends of a trajectory are small enough, so that electron propagation in these two domains can be described by plane waves. Thus formula (9) is valid if $\epsilon_F\tau_{\mathrm{tr}} \gg 1$ (semiclassical regime), and for the arguments sufficiently close to each other pairwise; in general, the condition $|\mathbf{r}_1-\mathbf{r}_4|, |\mathbf{r}_2-\mathbf{r}_3| \ll v_F\tau_s$ must be satisfied; however, in the special case of Eq. (8), we need the product of Green functions with $\mathbf{r}_1 = \mathbf{r}_3$ and $\mathbf{r}_2 = \mathbf{r}_4$, and this condition can be eased. Indeed, the integral over the angular variables in Eq. (9) is dominated by the close to each other momenta $\mathbf{p}_1$ and $\mathbf{p}_2$, and therefore a weaker requirement, $|\mathbf{r}_1-\mathbf{r}_4|, |\mathbf{r}_2-\mathbf{r}_3| \ll l_{\mathrm{tr}}, R_c$, should be satisfied. Here $R_c = v_F/\omega_c$ is the cyclotron radius, $l_{\mathrm{tr}} = v_F\tau_{\mathrm{tr}}$ is the electron transport relaxation length, and v_F and p_F are the Fermi velocity and momentum, respectively. Further derivation of the Hartree correction requires substitution of Eq. (9), with arguments $\mathbf{r}_1 = \mathbf{r}_2 = \mathbf{r}'$ and $\mathbf{r}_3 = \mathbf{r}_4 = \mathbf{r}''$, into Eq. (8) and integration over the difference between $\mathbf{r}'$ and $\mathbf{r}''$. These coordinates are coupled by the interaction potential $V(\mathbf{r}'-\mathbf{r}'')$. Hence, performing the integration we can exploit Eq. (9) only if the range of the interaction potential d is sufficiently short, $d \ll l_{\mathrm{tr}}, R_c$.

For a macroscopically homogeneous sample, classical probability $\mathcal{D}$ depends only on difference of its coordinates $\mathbf{r} = \mathbf{r}_2 - \mathbf{r}_1$. Using this fact and Eqs. (8) and (9), we find:

$$\frac{\delta\nu_H(\epsilon)}{\nu} = -\frac{2}{\pi}\mathrm{Re}\int_\epsilon^\infty d\omega \int \frac{d\phi_1 d\phi_2}{(2\pi)^2}\frac{\partial\mathcal{D}(\omega;\mathbf{r}=0;\phi_1,\phi_2)}{\partial\omega} \times V(2p_F|\sin[(\phi_1-\phi_2)/2]|), \quad (10)$$

where $V(q)$ is a Fourier transform of the interaction potential. Formula (10) can be easily generalized for the case of the long range Coulomb potential $V(q) = 2\pi e^2/\kappa q$ (here κ is the dielectric constant). For such a potential, $V(q)$ in Eq. (10) should be replaced (see, e.g., Ref. [7]) by the screened potential $V_{\mathrm{scr}}(\mathbf{q}) = 2\pi e^2/[\kappa(q + 2/a_B)]$ with $a_B = \hbar^2\kappa/me^2$ being the effective Bohr radius. Note that the range of screened Coulomb potential, $V_{\mathrm{scr}}(r)$, is of the order of the effective Bohr radius, and, therefore, much smaller than l_{tr} and R_c. Integration over the frequency in Eq. (10) immediately gives the resulting expression for the Hartree contribution:

$$\frac{\delta\nu_H(\epsilon)}{\nu} = \frac{2}{\pi}\mathrm{Re}\int\frac{d\phi_1 d\phi_2}{(2\pi)^2}\mathcal{D}(\epsilon;\mathbf{r}=0;\phi_1,\phi_2) \times V_{\mathrm{scr}}(2p_F|\sin[(\phi_1-\phi_2)/2]|). \quad (11)$$

The exchange correction to DOS can be obtained in a similar way as Eq. (11). The difference is, however, that in the exchange counterpart of Eq. (10), one should put the retarded screened potential $V_{\mathrm{scr}}(\omega,\mathbf{q})$ instead of $V(q)$. The retardation of the interaction potential makes immediate integration over frequencies, that led to Eq. (11), no longer possible. The resulting expression for the exchange correction has the form:

$$\frac{\delta\nu_F(\epsilon)}{\nu} = \frac{1}{\pi}\mathrm{Re}\int_\epsilon^\infty d\omega\int\frac{d\mathbf{q}}{(2\pi)^2}V_{\mathrm{scr}}(\omega,\mathbf{q})\frac{\partial\Delta(\omega;\mathbf{q})}{\partial\omega}, \quad (12)$$

where $V_{\mathrm{scr}}(\omega,\mathbf{q}) = V(q)/[1 + V(q)\Pi(q,\omega)]$, with $\Pi(q,\omega) = \nu[1 + i\omega\Delta(\omega;\mathbf{q})]$ being the polarization operator, and

$$\Delta(\omega;\mathbf{q}) = \int\frac{d\phi_1}{2\pi}\int\frac{d\phi_2}{2\pi}d\mathbf{r}\,e^{-i\mathbf{qr}}\,\mathcal{D}(\omega;\mathbf{r};\phi_1,\phi_2). \quad (13)$$

Equations (11) and (12) express the interaction correction to tunneling DOS in terms of classical probability density $\mathcal{D}(\epsilon;\mathbf{r};\phi_1,\phi_2)$ and are valid at energies $\epsilon \lesssim E_F$. Function $\mathcal{D}(\epsilon;\mathbf{r};\phi_1,\phi_2)$ can be found from the Boltzmann equation describing the scattering of electrons on impurities. In the special case $\tau_s \ll \tau_{\mathrm{tr}}$ we are interested in, scattering on small angles dominates the collision integral. With account for this simplification, the transport equation takes the Fokker–Planck form:

$$(-i\omega + \frac{\mathbf{p}_2}{m_e}\frac{\partial}{\partial\mathbf{r}} + \omega_c\frac{\partial}{\partial\phi_2} - \frac{1}{\tau_{\mathrm{tr}}}\frac{\partial^2}{\partial\phi_2^2})\mathcal{D}(\omega;\mathbf{r};\phi_1,\phi_2) = 2\pi\delta(\phi_1-\phi_2)\,\delta(\mathbf{r}). \quad (14)$$

Equation (14) describes electron motion along the cyclotron orbit accompanied by angular diffusion caused by scattering on a random potential. Analysis of Eq. (14) yields:

$$\mathcal{D}(\omega;\mathbf{q};\phi_1,\phi_2) = \sum_n \mathcal{D}_n(\omega;\mathbf{q};\phi_1,\phi_2), \quad (15)$$

$$\mathcal{D}_n(\omega;\mathbf{q};\phi_1,\phi_2) = \frac{e^{in(\phi_2-\phi_1)}e^{iR_c\mathbf{q}[(\mathbf{p}_1-\mathbf{p}_2)/p_F\times\mathbf{z}]}}{-i(\omega - n\omega_c) + R_c^2q^2/2\tau_{\mathrm{tr}} + n^2/\tau_{\mathrm{tr}}}, \quad (16)$$

where $\mathbf{z}$ is a unit vector parallel to the magnetic field. The solution (15)–(16) is valid for $qR_c \ll \omega_c^2\tau_{\mathrm{tr}}/(|\omega| + \omega_c)$. The obtained solution of transport equation together with Eqs. (11) and (12) enables us to calculate the interaction correction to DOS in the classical magnetic field.

At small frequencies, $\omega \ll \hbar/\tau_{\mathrm{tr}}$, the $n = 0$ term in Eq. (16) contains a part independent on initial, ϕ_2, and final, ϕ_1, directions of the electron momentum. This part dominates in Eq. (15), so that $\mathcal{D}(\omega;\mathbf{q};\phi_1,\phi_2) \approx 1/(-i\omega + R_c^2q^2/2\tau_{\mathrm{tr}})$. This limit corresponds to the diffusion regime studied in Refs. [7,8]. Comparing Eqs. (12) and (11), one sees that in the diffusion regime the exchange correction to DOS contains the interaction potential, $V_{\mathrm{scr}}(\omega,k)$, at very small, $kR_c \lesssim 1$, momentum transfers, while for the Hartree correction the potential at large, $k \lesssim p_F$, momentum transfers is important. Provided that the interaction potential decreases rapidly with the increase of k, one concludes [7] that in the diffusion regime, $\epsilon \ll \hbar/\tau_{\mathrm{tr}}$, the exchange contribution to DOS, Eq. (12), dominates over the Hartree one. The resulting expression for the interaction correction to DOS in this regime yields:

$$\frac{\delta\nu(\epsilon)}{\nu} = -\frac{\hbar(\omega_c\tau_{\mathrm{tr}})^2}{8\pi\epsilon_F\tau_{\mathrm{tr}}}\ln\left[\frac{|\epsilon|\tau_{\mathrm{tr}}}{\hbar}\right]\ln\left[\frac{|\epsilon|\tau_{\mathrm{tr}}a_B^4}{\hbar R_c^4}\right]. \quad (17)$$

Equation (17) corresponds to the Altshuler–Aronov–Lee [7] result with the renormalized by magnetic field diffusion coefficient $D = R_c^2/2\tau_{\mathrm{tr}}$.

At larger frequencies, $\omega \gg \hbar/\tau_{\mathrm{tr}}$, the probability $\mathcal{D}(\omega;\mathbf{q};\phi_1,\phi_2)$ describes quasiballistic motion of the electron along the cyclotron orbit. After the period $2\pi/\omega_c$ the electron approaches the vicinity of the initial point with a momentum only slightly deflected with respect to the initial direction. According to Eq. (11), that means that in the quasiballistic limit the Hartree contribution to DOS contains interaction potential at small momentum transfers. Thus, in contrast to the diffusion regime, there is no special reason why the Hartree contribution should be smaller than the exchange one. Furthermore, the Hartree correction appears to dominate over the exchange one. Indeed, when comparing the magnitudes of the Hartree and the exchange correction to DOS, one should take into account differences in the way screening influences the two corrections. The Hartree correction contains a screened interaction potential at zero frequency. On the other hand, expression for the exchange correction, Eq. (12), contains integration over frequencies. The polarization operator,

$\Pi(\omega, q)$, exhibits singularities at the same frequencies as the return probability, which causes suppression of the interaction potential $V_{\rm scr}(\omega, q)$, and, in turn, suppression of the exchange contribution at energies close to multiples of the cyclotron quantum. As a result, it is the Hartree contribution [11] that determines the peak structure of the interaction correction to tunneling DOS. Calculation of this correction, which consists in substituting Eqs. (15)–(16) and expression for $V_{\rm scr}$ into Eq. (11) and straightforward integration, gives the resulting expressions (1) and (2). Peak structure in tunneling DOS is well pronounced for $n \lesssim \sqrt{\omega_c \tau_{\rm tr}}$. For larger n the width of the peaks becomes comparable with $\hbar\omega_c$, and the oscillating structure of DOS disappears.

The resulting energy dependence of the tunneling DOS of the interacting electron gas in a classical magnetic field, obtained by numerical integration of Eqs. (11) and (12) with account for Eq. (15), is shown in Fig. 1.

The range of magnetic fields B, where the oscillations of the density of states $\nu(\epsilon)$ are caused by the interaction effects, is confined by the condition $1/\tau_{\rm tr} \ll \omega_c \ll 1/\tau_s$, and depends on the sample quality. In a high mobility sample with the typical values [12] of relaxation times, $\tau_s \approx 10$ ps and $\tau_{\rm tr} \approx 300$ ps, this range is $10^{-3} \lesssim B \lesssim 5 \cdot 10^{-2}$ T. Correspondingly, the typical number N of observable peaks, $N \lesssim \sqrt{\tau_{\rm tr}/\tau_s}$, for these samples is $N \approx 6$.

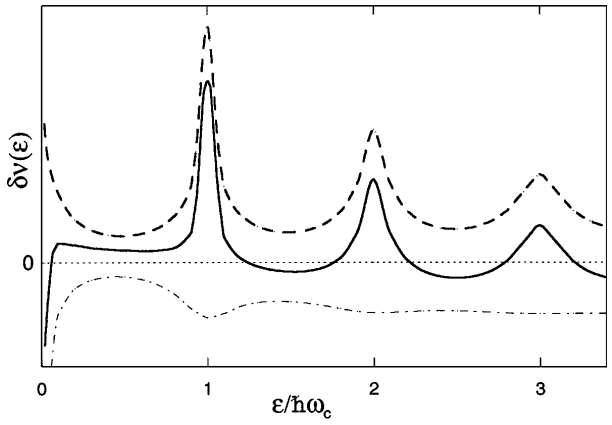

FIG. 1. Energy dependence of the exchange (dashed), the Hartree (dotted), and the total (solid) interaction correction to tunneling DOS in a classical magnetic field. Energy ϵ is measured from the Fermi level. The curves are calculated for $\omega_c \tau_{\rm tr} = 40$.

To conclude, we have shown that classical ($\omega_c \tau_s \ll 1$) magnetic field affects strongly the tunneling density of states of interacting electron gas. The tunneling DOS is found to be an oscillating function of energy with the characteristic period $\hbar\omega_c$.

Discussions with R. C. Ashoori and H. B. Chan are acknowledged with gratitude. This work was supported by NSF Grant DMR-9423244.

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