Density matrix approach to electron transport in an arbitrary magnetic field and negative magnetoresistance

V. E. Kaminskii*
Institute of RadioEngineering and Electronics, Russian Academy of Sciences, 101999 Moscow, Russia
(Received 30 July 2002; published 6 February 2003)

A solution of a kinetic equation for the one-particle density matrix is obtained for deformation potential scattering. Expressions for conductivity tensor components as sums over states in a magnetic field are derived. It is shown that without a magnetic field they do not differ from classical expressions. In a weak magnetic field, the magnetoresistance of a nondegenerate electron gas is positive for high electron mobility and negative for low mobility. For intermediate mobility, its sign changes with the magnetic-field increase. The magnetoresistance of a degenerate electron gas is nonzero. In a quantizing magnetic field, the conductivity oscillates as a function of the magnetic field. The oscillation amplitude and the mean magnetoresistance value increase with the magnetic field or mobility growth.

DOI: 10.1103/PhysRevB.67.085201 PACS number(s): 72.20.My, 73.43.Qt

I. INTRODUCTION

The classical kinetic theory predicts that magnetoresistance that can be defined as $d=\rho_{xx}(B)/\rho_0-1$ is always positive for a nondegenerate electron gas and zero for a degenerate electron gas. But in semiconductors[1–7] as well as in metals[8] and other conducting objects,[9] negative magnetoresistance (NMR) is frequently observed at sufficiently low temperatures. Extensive experimental studies of this phenomenon have revealed two types of conducting objects with a different NMR behavior. In the first-type objects, NMR is observed at low fields, while at higher fields magnetoresistance become positive.[1,2,7] Altshuler *et al.*[10] have explained this phenomenon by quantum correction to the conductivity, $\Delta\sigma_{xx}^{WL}$ (weak localization WL). The expression $\Delta\sigma_{xx}^{WL}(B)$ is widely used in the analysis of experimental data to extract the phase breaking time and its temperature dependence through experimental curve fitting. It should be pointed out that some deviation of experimental curves from $\Delta\sigma_{xx}^{WL}(B)$ takes place almost without exception. But in some experiments the deviation is significant and NMR correction less than half[11] or even a few percent[12] of the value predicted by weak-localization theory[10] is observed.

In the second-type objects, NMR is observed up to the onset of Shubnikov–de Haas oscillations and may continue as a mean value of oscillations.[3–6] This high-field effect is generally attributed to the electron-electron interaction (EEI) that was considered theoretically in Refs. 13 and 14. The main contribution to NMR in the $B>B_{tr}$ range comes from the weak-localization effect. In the $B>1/\mu$ range, where μ is the mobility, it results from the correction to the conductivity due to the EEI $\Delta\sigma_{xx}^{ee}(B)$. The values $\Delta\sigma_{xx}^{WL}(B)$ and $\Delta\sigma_{xx}^{ee}(B)$ differ significantly for different material systems. This should lead to different temperature dependences $\rho_{xx}(B,T)$, where T is the temperature. But in experiments qualitatively similar $\rho_{xx}(B,T)$ dependences are observed for doped GaAs (Ref. 3) and Ge,[7] AlGaAs/GaAs,[6] GaAs/InGaAs/GaAs,[4] and Ge/SiGe (Ref. 15 and 16) heterostructures. The $\rho_{xx}(B,T)$ decreases sharply at low fields, and continues to decrease slightly at higher fields. The temperature dependence $\rho_{xx}(B,T)$ below a characteristic magnetic field B_e is nonmetallic. At the characteristic field B_e, all curves $\rho_{xx}(B,T)$ for different temperatures intersect each other. According to the theory, the WL and EEI contributions to NMR become more prominent when mobility increases at a fixed electron density. However, in practice, the reverse dependence is observed.[6] All the aforementioned point to the existence of another NMR mechanism.

Recently, a simple classical mechanism for this effect was proposed in Refs. 17–19. It leads to an exponentially small NMR in a weak magnetic field and to its saturation for $\mu B \gg 1$. For this mechanism, NMR also increases when mobility increases at the fixed electron density. The main idea of this mechanism is that for short-range scattering centers, the Boltzmann approach does not work even as the first approximation when $\mu B>1$. In this case, the generalized Boltzmann equation should be used.[19] But at low temperatures, the condition $\mu B>1$ usually corresponds to the $\alpha=\hbar\omega/kT>1$ condition, where $\omega=qB/m$ is the cyclotron frequency. In this case, the magnetic-field quantization should be taken into consideration. Adams and Holstein[20] have considered electron conductivity in the quantizing magnetic field for $\mu B\gg 1$. Isihara and Smrćka[21] have used the Green-function approach to describe the conductivity tensor in low/intermediate magnetic fields ($\mu B\leqslant 1$) for a two-dimensional degenerate electron gas. In both cases, the obtained expressions describe Shubnikov–de Haas oscillations. But in low magnetic fields the expressions obtained in Ref. 20 for $\rho_{xx}(B)$ are not valid, since they diverge for $B\to 0$. In the low-field limit the conductivity tensor[21] coincides with the classical one and the magnetoresistance is zero.

This work eliminates these difficulties by deriving expressions for conductivity tensor components that are correct in an arbitrary magnetic field. In Sec. II, for deformation potential scattering we present a steady-state perturbation solution of the Liouville equation for the one-particle density matrix. General expressions for a conductivity tensor are obtained. In Sec. III, these results are applied to the low-field magnetotransport. The NMR dependence on electron-gas parameters is considered. Theoretical predictions are compared

with available experimental results. In Sec. IV, we consider briefly electron transport in a quantizing magnetic field. In Sec. V, a brief summary and concluding remarks are presented.

II. THEORY

It is well known that in a quantizing magnetic field, matrices of electron momentum components perpendicular to magnetic field have no diagonal elements. In this case, the Boltzmann equation cannot be, in principle, used for the electron-transport description. Therefore the density-matrix approach is needed. We restrict the scope of this paper by the description of the electron transport in an isotropic semiconductor with the parabolic dispersion law. In general, the electric current density can be calculated from the expression

$$\mathbf{j}=\mathrm{Tr}(R\mathbf{J}), \tag{1}$$

where $\mathrm{Tr}(\cdots)$ is the operator trace, R is the statistical operator (density matrix), $\mathbf{J}$ is the operator of current in the magnetic field.

In our case, the one-particle Hamiltonian of a semiconductor is given by the following expression:

$$H=H_0+W+U=H_e+H_e+W+U, \tag{2}$$

where H_e is the electron Hamiltonian in magnetic field, H_p is the phonon Hamiltonian, W is the operator of electron-phonon interaction, $U=-qE_x x$ is the operator of potential energy. For simplicity, we assume that the electric field is directed along the x axis and the magnetic field is directed along the z axis. The statistical operator in the Schrödinger representation is described by the Liouville equation

$$i\hbar\frac{\partial R}{\partial t}=[HR]. \tag{3}$$

To solve this equation, we consider the sum $W+U$ as a perturbation. If we take vector potential in the gauge $\mathbf{A}=(0,Bx,0)$, the wave function (WF) (eigenfunction) of the Hamiltonian H_0 is

$$\psi=\exp[i(k_y y+k_z z)]f(\eta)\chi_s|N\rangle, \tag{4}$$

where

$$\eta=\frac{x+k_y\lambda^2}{\lambda},\quad \lambda^2=\frac{\hbar}{qB},$$

$\mathbf{k}$ is the electron wave vector, χ_s is the spin WF, $|N\rangle$ is the phonon WF in the second quantization representation. The WF (4) along the x axis in the magnetic field is always normalized to the unity and determined by

$$f_n(\eta)=\left(\frac{1}{\sqrt{\pi}\lambda 2^n n!}\right)^{1/2}\exp\left(-\frac{\eta^2}{2}\right)H_n(\eta). \tag{5}$$

The electron energy is

$$E=\frac{\hbar^2 k_z^2}{2m}+\hbar\omega\left(n+\frac{1}{2}\right)+\mu_B g s B, \tag{6}$$

where μ_B is the Bohr magneton, g is the Lande factor, $s=\pm\frac{1}{2}$. It can be seen that in our representation the electron state is described by a set of quantum numbers (n,ky,kz,s).

Since the spin operator does not stand in current and perturbation operators their matrix elements are diagonal with respect to the spin numbers. Assume also that the phonon system is in equilibrium. Then all matrix elements are diagonal with respect to the phonon numbers. To shorten the notation, we omit them in all matrix elements and take into account in sums over the states.

The basis and the procedure of deriving perturbative solution of Eq. (3) are given in detail in Ref. 22. Using the WF (4) for stationary conditions from Eq. (3), we get

$$(E_1-E_2)R_{12}+U_{13}R_{32}-R_{13}U_{32}+i\pi S_{12}=0, \tag{7}$$

where

$$\begin{aligned}S_{12}=\sum w(\mathbf{q})\times\{&\delta(E_2-E_3-\hbar\Omega)[(N+1)M_{14}^{+}\\&\times(\delta_{43}-R_{43})M_{35}R_{52}-N(\delta_{14}-R_{14})M_{43}^{+}R_{35}M_{52}]\\&+\delta(E_2-E_3+\hbar\Omega)[NM_{14}(\delta_{43}-R_{43})M_{35}^{+}R_{52}\\&-(N+1)(\delta_{14}-R_{14})M_{43}R_{35}M_{52}^{+}]\\&+\delta(E_1-E_3-\hbar\Omega)[(N+1)R_{14}M_{45}^{+}(\delta_{53}-R_{53})M_{32}\\&-NM_{13}^{+}R_{34}M_{45}(\delta_{42}-R_{42})]\\&+\delta(E_1-E_3+\hbar\Omega)[NR_{14}M_{45}(\delta_{53}-R_{53})M_{32}^{+}\\&-(N+1)M_{13}R_{34}M_{45}^{+}(\delta_{42}-R_{42})]\},\end{aligned} \tag{8}$$

$\hbar\Omega(\mathbf{q})$ is the phonon energy, N is the phonon filling number, $w(\mathbf{q})$ is the function that describes the electron-phonon interaction, M_{ij} is the matrix element of electron-phonon interaction, the "+" sign denotes the Hermit's conjugation. In the Eq. (8), the sum runs over all electron states, except $|1\rangle$ and $|2\rangle$, and over phonon wave vectors.

It is well known that Eq. (3) does not describe irreversible motion of the electron system. To introduce irreversibility, additional assumptions are needed.[20] This can be done by modification of the total Hamiltonian or by some artificial mathematical method that takes into consideration the interaction between the system and the environment. The latter may be the widely used hypothesis on the initial form of the density-matrix diagonal part, i.e., the distribution function. We represent the density matrix in the form of

$$R_{12}=F_1\delta_{12}+G_{12}\delta(k_{y1}-k_{y2}), \tag{9}$$

where $G_{12}\equiv G_{n1,n2}(k_{z1},k_{z2})$. If the electric field does not break the electron-gas space homogeneity, the Fermi-Dirac distribution function that depends on the electron energy [Eq. (6)] and the chemical potential (Fermi energy) E_f is normally taken as $F_1=F(E_1-E_f)$. At high temperatures, this ap-

proximation is quite sufficient. But at low temperatures, the combined action of the electric field and electron scattering changes the form of the distribution function. In the first nonvanishing scattering order the correction to the distribution function has been obtained in Ref. 20. But it has the physical meaning at high magnetic fields only. The general procedure for deriving corrections is described in Ref. 22. As the matter of fact, it is the Gibbs statistical operator expansion into a series by orders of interaction. In the linear approximation by the electric field, the distribution function can be represented as

$$F_1+\frac{\partial F_1}{\partial E}U_{11}Z,$$

where U_{12} is the potential-energy matrix element, Z is the function that is the result of summation over all scattering orders. This function depends on the magnetic field, parameters of scattering, and distribution function F_1 itself. In this paper, we do not intend to analyze the Z function. It should be noted that it is very small or equal to zero only when $B=0$. From Eq. (9) we get

$$R_{12}=\left(F_1+\frac{\partial F_1}{\partial E}U_{11}Z\right)\delta_{12}+G_{12}\delta(k_{y1}-k_{y2}). \quad (10)$$

Note here that similar to Eqs. (4)–(6), solutions can be obtained if the potential energy U is included into the basic Hamiltonian H_0. Such a solution has been used in Ref. 20. But the WF of this Hamiltonian does not meet the conditions of the Hermit's conjugation and completeness. It becomes obvious in the limit $B\to 0$.

Now we obtain the linear solution by G of Eq. (7). For the acoustic phonon the frequency $\Omega=v_s q$ and for the deformation acoustic-phonon scattering the interaction function is

$$w(\mathbf{q})=w_0 q, \quad (11)$$

where $w_0=\hbar E_A^2/2\rho v_s^2$, v_s is the sound velocity. If the temperature is not very low the phonon number can be written as follows:

$$N\approx\frac{kT}{\hbar v_s q}\gg 1. \quad (12)$$

From the semiconductor theory it is well known that for $T>1$ K the acoustic-phonon energy can be omitted in the energy conservation law. Then, substituting Eqs. (8)–(12) into Eq. (7) after simple integration we get

$$(E_1-E_2)G_{12}+U_{12}(F_2-F_1)-iqE_x\eta_{12}Q_{12}$$
$$+i\frac{kTw_0}{2\pi\lambda^2\hbar}\sum_3\int dk_{z3}\{[\delta(E_2-E_3)+\delta(E_1-E_3)]$$
$$\times[G_{12}-\delta_{12}G_{33}(k_{z1}-k_{z2}+k_{z3},k_{z3})]\}=0,$$

where $\eta_{12}=\int f_1\eta f_2 d\eta$,

$$Q_{12}=\frac{kTw_0}{2\pi\lambda\hbar}\sum_3\int dk_{z3}\left[\frac{\partial F_2}{\partial E_2}(1+F_2-F_1)\delta(E_2-E_3)\right.$$
$$\left.+\frac{\partial F_1}{\partial E_1}(1+F_1-F_2)\delta(E_1-E_3)\right]Z.$$

Note here that in Eq. (8) the integration is simple for detormation acoustic (DA) phonons only, since the product $Nw(\mathbf{q})$ is independent of q. For other scattering mechanisms, it is impossible to obtain such a simple expression. It is easy to check that the solution of this equation is

$$G_{12}=\frac{U_{12}(F_1-F_2)+iq\eta_{12}Q_{12}}{E_1-E_2+i\hbar\nu_{12}}, \quad (13)$$

where

$$\nu_{12}=\frac{kTw_0}{2\pi\lambda^2\hbar^2}\sum_3\int dk_{z3}[\delta(E_2-E_3)+\delta(E_1-E_3)].$$

The current and potential matrix elements are easily calculated and given by the expressions

$$j_{x12}=i\frac{p_{12}}{\lambda^2}D_{12},\quad j_{y12}=\left(\frac{x_{12}}{\lambda^2}-k_{y1}\right)D_{12},$$
$$j_{z12}=-k_{z1}D_{12},\quad U_{12}=-qE_x x_{12},$$

where

$$D_{12}=\frac{q\hbar}{m}\delta(k_{y1}-k_{y2})\delta(k_{z1}-k_{z2}),$$
$$p_{12}=\sqrt{\frac{n_1}{2}}\delta_{n_1,n_2+1}-\sqrt{\frac{n_2}{2}}\delta_{n_1,n_2-1}.$$

Substituting these matrix elements and Eq. (13) into Eq. (1), we obtain a conductivity tensor

$$\sigma_{xx}=\frac{q^2\omega}{(2\pi)^2\hbar}\sum_{s,n}(n+1)$$
$$\times\int dk_z\left[\frac{\nu_{n,n+1}(F_n-F_{n+1})-\omega Q_{n,n+1}}{\omega^2+\nu_{n,n+1}^2}\right], \quad (14)$$

$$\sigma_{xy}=\frac{q^2\omega}{(2\pi)^2\hbar}\sum_{s,n}(n+1)$$
$$\times\int dk_z\left[\frac{\omega(F_n-F_{n+1})+\nu_{n,n+1}Q_{n,n+1}}{\omega^2+\nu_{n,n+1}^2}\right], \quad (15)$$

where

$$F_n\equiv F(E_n),\quad E_n=\frac{\hbar^2k_z^2}{2m}+\hbar\omega\left(n+\frac{1}{2}\right)+\mu_B gsB,$$
$$\nu_{nm}=\frac{\hbar\omega}{4\sqrt{kT}\tau_{DA}}\sum_l\left[\frac{\theta(z_{nl})}{\sqrt{z_{nl}}}+\frac{\theta(z_{ml})}{\sqrt{z_{ml}}}\right],$$

$$Q_{nm}=\left[\frac{\hbar\nu_{nn}}{2}\frac{\partial F_n}{\partial E_n}(1+F_n-F_m)+\frac{\hbar\nu_{mm}}{2}\frac{\partial F_m}{\partial E_m}(1+F_m-F_n)\right]Z,$$

$$z_{nl}=\hbar\omega(n-l)+\frac{\hbar^2k_z^2}{2m},\quad \tau_{DA}=\frac{\pi\hbar^4\rho v_s^2}{\sqrt{2}(mkT)^{3/2}E_A^2}.$$

Using the obtained results it is easy to show that

$$\sigma_{zz}=\frac{q^2\omega}{(2\pi)^2\hbar}\sum_{s,n}\int dk_z\frac{k_z}{\nu_{n,n}}\left(-\frac{\partial F_n}{\partial k_z}\right). \tag{16}$$

If the magnetic field is so strong that $\omega\gg\nu$, Eq. (14) is transformed into the equation that has been obtained in Ref. 20. If in this case for a degenerate electron gas the Poisson's formula is used for summation in Eqs. (14) and (15), we obtain expressions that are identical to ones obtained in Refs. 20 and 21.

The conductivity tensor components (14)–(16) are finite in an arbitrary magnetic field. Their general analysis is very complicated. Therefore, let us consider some simple cases.

III. WEAK-FIELD MAGNETOTRANSPORT

In the weak-field limit at $\alpha\ll1$, the energy of the spin interaction with the magnetic field can be assumed to be zero and integration can be used instead of summation over n in Eqs. (14) and (15). After transformation of variables and integration over angular variables, the conductivity tensor can be written as

$$\sigma_{xx}=\sigma_{Bxx}+\frac{2q^2}{\sqrt{\pi}m}N_c\alpha\int\sqrt{x}dx\frac{\nu}{\nu^2+\omega^2}\left(-\frac{\partial F}{\partial x}\right), \tag{17}$$

$$\sigma_{xy}=\sigma_{Bxy}+\frac{2q^2}{\sqrt{\pi}m}N_c\alpha\int\sqrt{x}dx\frac{\omega}{\nu^2+\omega^2}\left(-\frac{\partial F}{\partial x}\right), \tag{18}$$

where σ_{Bxx}, σ_{Bxy} are the components of Drude conductivity tensor that can be calculated from the Boltzmann equation, N_c is the conductivity band density of states, $x=E/kT$,

$$\nu=\frac{1}{2\tau_{DA}}(\sqrt{x}+\sqrt{x+\alpha}).$$

To derive Eqs. (17) and (18), we assume $Z=0$. In this approximation, the σ_{zz} (16) coincides with a classical expression. It can be seen that there are additional terms in Eqs. (17) and (18), as compared with classical expressions. They reflect the magnetic-field effect on electron motion. As is well known, in the transverse magnetic field the electron motion is diffusion. Using identity $\lambda^2(n+1)\equiv x_{n,n+1}^2=(k_y)_{n,n+1}^2\lambda^2$ it follows from Eqs. (14) and (15) that in the magnetic field the diffusion coefficient is

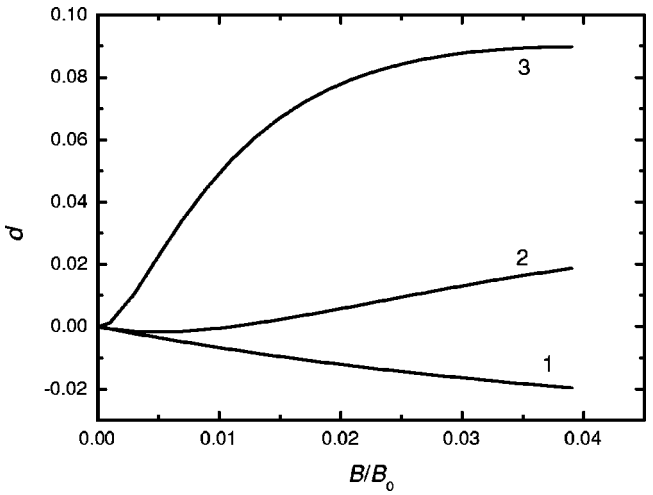

FIG. 1. Classical magnetoresistance field dependence (curve 1, $\gamma=5$; curve 2, $\gamma=20$; curve 3, $\gamma=100$).

$$D_H=\omega^2\left\langle\frac{\nu_{n,n+1}x_{n,n+1}^2}{\nu_{n,n+1}^2+\omega^2}\right\rangle.$$

For $\omega\tau\ll1$, we obtain

$$D_H=\frac{1}{m}\left\langle\frac{E_n+\frac{\hbar\omega}{2}}{\nu_{n,n+1}}\right\rangle=D+\frac{\hbar}{2m}\omega\tau,$$

where D is the diffusion coefficient without the magnetic field, τ is the momentum relaxation time, and the brackets denote statistical averaging. As can be seen in a weak magnetic field, the diffusion coefficient and hence the electron mobility increase. This results in NMR.

Equations (17) and (18) are formally obtained for deformation potential scattering. But at $B=0$ they coincide with classical expressions. Therefore, they should be valid for other scattering mechanisms at least in a weak magnetic field. It is well known that the classical magnetoresistance $\rho_{xx}(B)=\rho(B)>\rho(0)$. In Eq. (17), there is an additional term that is linear by the magnetic field. It results in a different $\rho(B)$ dependence. In the $0<B<B_1$ range it has the minimum $\rho_m<\rho(0)=\rho(B_1)$. The values of B_1 and ρ_m depend on parameters of the electron system and the temperature. It follows from Eq. (17) that for nondegenerate electron gas ($E_f<0$) the value $d=\Delta\rho/\rho(0)$ is independent of the electron density. In this case the parameters $B_0=mkT/q\hbar$ and $\gamma=kT\tau/\hbar$ determine the ρ_m and B_1 values.

Figure 1 shows $d(B)$ dependences for different γ values. They are calculated for a wideband semiconductor with electron conductivity and $E_f/kT=-5$. At constant temperature, the calculated dependences correspond to semiconductors with different electron mobilities. For the room temperature and the effective mass equal to the free-electron mass the mobilities are 200, 800, and 4000 $cm^2/V\,s$, respectively. Under these conditions the magnetic field varies from 0 to 10 T. As can be seen from Fig. 1 for a low electron mobility, MR is negative within the whole range and B_1 is infinite. For a high electron mobility, MR is positive and $B_1=0$. For inter-

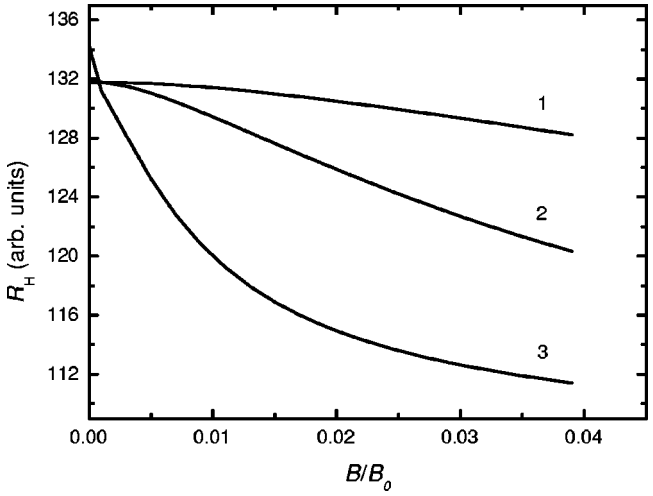

FIG. 2. The Hall constant as a function of the magnetic field (curve 1, $\gamma=5$; curve 2, $\gamma=20$; curve 3, $\gamma=100$).

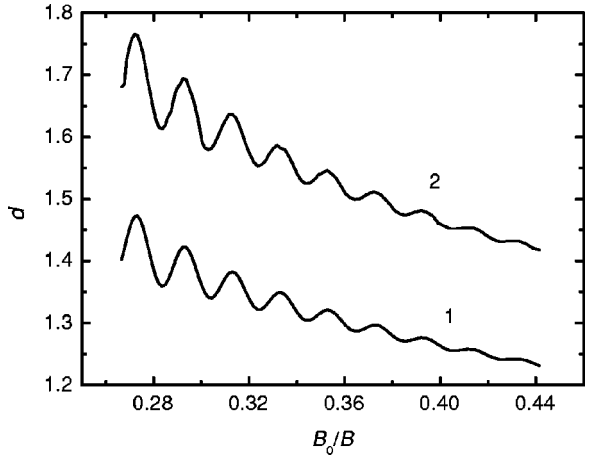

FIG. 3. Low-temperature magnetoresistance field dependence in a degenerate electron gas (curve 1, $\gamma=43$; curve 2, $\gamma=866$).

mediate mobility values B_1 is finite and MR changes the sign. Note here that dependences that are similar to the curve for $\gamma=5$ are observed in three-,[1] two-,[11] and one-dimensional[23] semiconductors. In Fig. 1, the dependences are calculated for the DA-phonon scattering. For other scattering mechanisms, the MR dependences on γ are qualitatively similar. Therefore the NMR temperature dependences should be similar too and the dependences of Fig. 1 can be used for qualitative description of experimental results. The $\nu(E)$ dependence determines quantitative parameters of NMR.

As is known, the scattering rate increases when the doping concentration is increased. As a result the γ value should decrease. It follows from Fig. 1 that NMR decreases when the γ value gets higher. This dependence might be a possible reason for different results obtained in samples[1,7] with different doping concentrations. When the temperature is increased the ionized impurity scattering rate decreases for any doping concentration and the γ value increases. It should qualitatively result in NMR temperature dependences that are observed in Refs. 1, 7, and 23. Figure 2 shows magnetic-field dependence of the Hall constant (R_H) for different γ values. The parameters of Fig. 1 were used for the calculation. It can be seen that the $R_H(B)$ tends to decrease slightly for all γ values. Since the electron concentration is assumed constant for the calculation these dependences reflect the magnetic-field dependence of the Hall factor.

In degenerate electron gas, classical MR is equal to zero and $d=0$. In this case, the derivative of the distribution function in Eq. (17) can be replaced by the δ function and we obtain

$$d=-\frac{\hbar\omega}{E_f+\hbar\omega}<0. \tag{19}$$

Equations (14), (15), (17), and (18) are formally obtained for bulk semiconductor. But in a weak magnetic field there is no strong difference between three- and two-dimensional electron-gas properties. Therefore, these equations can be used for qualitative description of different dependences that are observed in degenerate two-dimensional electron gas.[11,23,24] Particularly, expression (19) describes experimental dependences[11] for $\alpha<1$.

IV. LOW-TEMPERATURE ELECTRON TRANSPORT IN QUANTIZING MAGNETIC FIELD

At low temperatures, MR should be an oscillating function of B^{-1} in a strong magnetic field. Figure 3 shows the $d(B)$ dependence for two different values of γ. At the room temperature they correspond to γ values of 5 and 100. The dependence is calculated by Eq. (14) for DA-phonon scattering at $T=4$ K and $E_f/kT=50$.

It is assumed that $Z=1$. It corresponds to the approximation of Ref. 20. If at $\omega\gg\nu$ we omit the first term in Eq. (14) and the second term in Eq. (15) we obtain the results of Ref. 20. It can be seen from Fig. 3 that for $E_f>\hbar\omega$ the parameter d oscillates with a period that is proportional to B^{-1}. The oscillation amplitude and the mean value are increased when B increases or the scattering rate decreases. The calculation shows that in this case the Hall constant R_H is independent of B. At the same time the oscillation amplitude depends strongly on Z. For $Z\ll1$ (Boltzmann electron gas), there are no oscillations and d smoothly increases with B.

V. CONCLUSION

The modern kinetic theory attributes NMR to the weak localization that needs irregular distribution of scattering centers. In this work, a kinetic equation for the density matrix is solved for deformation potential scattering in quantizing magnetic field. It is shown that transverse NMR can be observed in an electron gas without weak localization. It is the result of consistent consideration of magnetic quantization. It should be noted that the used kinetic approach is valid for a nondegenerate electron gas only at $\gamma\gg1$. The obtained expressions, opposite to Ref. 20, are valid for an arbitrary magnetic field. In a strong magnetic field the prevailing term of σ_{xx} (Eq. 14) is similar to the expression in

Ref. 20. But in this paper it is shown that unlike the implicit assumption of Ref. 20 the value of Z is not equal to unity and depends on the magnetic field and temperature. The calculation shows that the NMR predicted in this work should be observed at low magnetic fields in semiconductors with low electron mobilities. In traditional semiconductors (Si, Ge, GaAs), low electron mobility is realized in heavily doped samples. But in this case the weak localization can be also observed and there is no single interpretation for NMR. But in some undoped semiconductors (such as GaN, graphite), the electron mobility is low. In such semiconductors, NMR should be observed in perfect crystals.

*Electronic address: kamin@mail.cplire.ru

[1] N. Zhdanova, M. Kagan, and E. Landsberg, Zh. Éksp. Teor. Fiz. **117**, 761 (2000) [Sov. Phys. JETP **90**, 662 (2000)].
[2] Y. Proskuryakov, A. Savchenko, S. Safonov, M. Pepper, M. Simmons, and D. Ritchie, Phys. Rev. Lett. **86**, 4895 (2001).
[3] S. Murzin, M. Weiss, A. Jansen, and K. Eberl, Phys. Rev. B **64**, 233309 (2001).
[4] G. Minkov, O. Rut, A. Germanenko, A. Sherstobitov, V. Shashkin, O. Khrykin, and V. Daniltsev, Phys. Rev. B **64**, 235327 (2001).
[5] V. Senz, T. Ihn, T. Heinzel, K. Ensslin, G. Dehlinger, D. Grützmacher, and U. Gennser, Phys. Rev. Lett. **85**, 4357 (2000).
[6] C. Huang, Y. Chang, C. Lee, H. Chou, H. Yeh, C. Liang, Y. Chen, H. Lin, H. Cheng, and G. Hwang, Phys. Rev. B **65**, 045303 (2001).
[7] A. Veinger, A. Zabrodskii, and T. Tisnek, Fiz. Tekh. Poluprovodn. **34**, 774 (2000) [Sov. Phys. Semicond. **34**, 746 (2000)].
[8] G. Bergmann, Phys. Rev. B **25**, 2937 (1982).
[9] L. Langer, V. Bayot, J. Issi, J. Heremans, C. Olk, L. Stockman, C. Van Haesendonck, and Y. Bruynseraede, Phys. Rev. Lett. **76**, 479 (1996).
[10] B. Altshuler, D. Khmelnitzkii, A. Larkin, and P. Lee, Phys. Rev. B **22**, 5142 (1980).
[11] G. Minkov, S. Negashev, O. Rut, A. Germanenko, V. Shashkin, O. Khrykin, and V. Danil'tsev, Phys. Rev. B **61**, 13 172 (2000).
[12] A. Mills, A. Ramirez, X. Gao, I. Pfeiffer, K. West, and S. Simon, cond-mat/0101020 (unpublished).
[13] B. Altshuler, A. Aronov, and P. Lee, Phys. Rev. Lett. **44**, 1288 (1980).
[14] A. Houghton, J. Senna, and S. Ying, Phys. Rev. B **25**, 6468 (1982).
[15] M. Hilke, D. Shahar, S. Song, D. Tsui, and Y. Xie, Phys. Rev. B **62**, 6940 (2000).
[16] S. Song, X. Wang, R. Chang, and J. Ketterson, Phys. Rev. Lett. **78**, 2200 (1997).
[17] A. Mirlin, D. Polyakov, F. Evers, and P. Wölfle, Phys. Rev. Lett. **87**, 126805 (2001).
[18] A. Dmitriev, M. Dyakonov, and R. Jullien, Phys. Rev. B **64**, 233321 (2001).
[19] A. Bobylev, F. Maao, A. Hansen, and E. Hauge, Phys. Rev. Lett. **75**, 197 (1995).
[20] E. Adams and T. Holstein, J. Phys. Chem. Solids **10**, 254 (1959).
[21] A. Isihara and L. Smrćka, J. Phys. C **19**, 6777 (1986).
[22] A. Akhiezer and S. Peletminskii, *Metody Statisticheskoi Fiziki* (Nauka, Moskva, 1977) [*Methods of Statistical Physics* (Pergamon Press, Oxford, New York, 1981)].
[23] A. Visser, V. Kadushkin, V. Kul'bachinskii, V. Kytin, A. Senichkin, and E. Shangina, Pis'ma Zh. Éksp. Teor. Fiz. **59**, 339 (1994) [JETP Lett. **59**, 363 (1994)].
[24] A. Kreshchuk, S. Novikov, and T. Polyanskaya, Fiz. Tekh. Poluprovodn. **31**, 459 (1997) [Sov. Phys. Semicond. **31**, 391 (1997)].