Magnetotransport in a 2D system with strong scatterers: renormalization of Hall coefficient caused by non-Markovian effects

© A.P. Dmitriev, V.Yu. Kachorovskii¶

Ioffe Physicotechnical Institute, Russian Academy of Sciences,
194021 St. Petersburg, Russia

(Получена 6 февраля 2008 г. Принята к печати 11 февраля 2008 г.)

We show that a sharp dependence of the Hall coefficient R on the magnetic field B arises in two-dimensional electron systems with strong scatterers. The phenomenon is due to classical memory effects. We calculate analytically the dependence of R on B for the case of scattering by antidots (modeled by hard disks of radius a), rendomly distributed with concentration $n_0 \ll 1/a^2$. We demonstrate that in very weak magnetic fields $(\omega_c\tau_{\mathrm{tr}} \lesssim n_0a^2)$ memory effects lead to a considerable renormalization of the Boltzmann value of the Hall coefficient: $\delta R/R \sim 1$. With increasing magnetic field, the relative correction to R decreases, then changes sign, and saturates at the value $\delta R/R \sim -n_0a^2$. We also discuss the effect of the smooth disorder on the dependence of R on B.

PACS: 05.60.+w, 73.43.Qt, 73.50.Bk, 73.50.Jt, 73.63.Hs

The problem of magnetoresistance (MR) in metal and semiconductor structures has been intensively discussed in literature during the past three decades. A large number of both theoretical and experimental papers on this subject was published. Most of these works were devoted to the case of the degenerate two dimensional electron gas where electrons move in the plane perpendicular to the magnetic field and scatter on a random impurity potential, so that only electrons with energy close to the Fermi energy participate in conductance. The simplest theoretical description of such a situation is based on the Boltzmann equation which yields the well-known expressions for the components of the conductivity tensor:

$$\sigma_{xx} = \frac{\sigma_0}{(1+\omega_c^2\tau_{\mathrm{tr}}^2)}, \quad \sigma_{xy} = \frac{\sigma_0(\omega_c\tau_{\mathrm{tr}})}{(1+\omega_c^2\tau_{\mathrm{tr}}^2)}. \tag{1}$$

Here $\sigma_0 = e^2n\tau_{\mathrm{tr}}/m$ is the Drude conductivity at $B=0$, $\omega_c = |e|B/mc$ is the cyclotron frequency, τ_{tr} is the transport scattering time and n is the electron concentration. The resistivity tensor, which can be obtained by inverting the conductivity tensor, has even simpler form:

$$\rho_{xx} = \frac{m}{e^2n\tau_{\mathrm{tr}}}, \quad \rho_{xy} = \frac{m\omega_c}{e^2n} = -RB, \tag{2}$$

where $R = 1/enc < 0$ is the Hall coefficient. Thus, in the frame of the Boltzmann approach, ρ_{xx} and R do not depend on magnetic field B. Experimental measurements of ρ_{xx} and R are widely used to find τ_{tr} and n.

It is known, that Eqs. (2) may become invalid due to a number of effects of both quantum and classical nature. The most remarkable of them is the quantum Hall effect. Another quantum effect, weak localization, leads to the decrease of ρ_{xx} with B, concentrated in the region of weak magnetic fields [1]. Besides, the dependence of ρ_{xx} on B appears due to quantum effects related to electron–electron interaction [2] (see also [3] for review). At the same time, both weak localization and electron–electron interaction (in frame of standard Altshuler–Aronov theory) do not result in any dependence of R on B (see [4] and [2,4–6], respectively), though such a dependence arises in the regime of strong localization [7].

The dependence of ρ_{xx} on B may also be caused by classical memory effects (ME) which are neglected in the Boltzmann approach. Such effects arise as a manifestation of non-Markovian nature of electron dynamics in a static random potential. Physically, a diffusive electron returning to a certain region of space „remembers" the random potential landscape in this region, so its motion is not purely chaotic as it is assumed in the Boltzmann picture. For $B=0$, non-Markovian corrections to kinetic coefficients are usually small. In particular, in the case of hard-core scatterers of radius a randomly distributed with concentration n_0, ME-induced relative correction to the resistivity is proportional to the gas parameter $\beta_0 = a/l = 2n_0a^2 \ll 1$ (here $l = v_{\mathrm{F}}\tau$, v_{F} is the Fermi velocity and τ is the mean free time). However, for $B \neq 0$ the role of ME is dramatically increased due to a strong dependence of the return probability on B. In particular, two years prior to Ref. [1] there appeared a publication [8] where a classical mechanism of strong negative magnetoresistance was discussed. The mechanism was investigated by the example of a gas of non-interacting electrons scattering on hard disks (antidots). It was shown that with increasing magnetic field there is an increasing number of closed electron orbits which avoid scatterers and therefore are not diffusive (see also recent discussions [9–11] of this mechanism). Electrons occupying these orbits do not participate in diffusion (so-called „circling electrons"). As a result, the longitudinal resistance turns out to be proportional to the factor $1-P$, where $P = \exp(-2\pi/\omega_c\tau)$ is the probability of the existence of the circular closed orbit Another classical mechanism was presented in Ref. [12], where the MR due to non-Markovian dynamics of electrons trapped in some region of space was discussed.

¶ E-mail: Kachor@vip1.ioffe.ru

Notwithstanding these developments, the role of classical effects in magnetotransport was underappreciated for a long time. A new boost to the research in this direction was given by Ref. [13], where it was shown that if electrons move in a smooth disorder potential and in a sufficiently strong magnetic fields a phenomenon called „classical localization“ occurs. This phenomenon leads to the exponential suppression of the longitudinal resistance: most electrons are trapped in localized equipotential trajectories and do not participate in diffusion. This work was followed by a series of works [14–21] discussing different aspects of classical magnetotransport in 2D systems. It was shown [15] that for lower magnetic fields near the onset of the classical localization the magnetoresistance is positive, i. e. the longitudinal resistance grows with increasing magnetic field. In Refs. [16,17] the combination of smooth disorder and strong scatterers (antidots) was considered. It was shown that in this system under certain conditions there are several regimes of the behavior of magnetoresistance depending on the strength of the magnetic field: first the longitudinal resistance decreases with growing field, then it saturates and then begins to grow. The role of non-Markovian effects in the cyclotron resonance was also discussed [18].

In Refs. [8–18] magnetoresistance was studied in a situation where the magnetic field is classically strong, that is where the parameter $\beta = \omega_c \tau$ is large. Recently, the region of classically small magnetic fields $\beta \ll 1$ was investigated numerically [19,20] for the case of electrons scattering on strong scatterers. It was shown [19] that memory effects due to double scattering of an electron on the same disk lead to a negative parabolic magnetoresistance (in the Ref. [8], where these processes were not taken into account, exponentially small MR was predicted). The numerical simulations [20] discovered a low-field classical anomaly of the MR. The anomaly was attributed to the memory effects specific for backscattering events. The simulations were performed for the 2D Lorenz gas which is a system of 2D electrons scattering on hard disks randomly distributed in plane with average concentration n. Magnetotransport in this system is characterized by two dimensionless parameters: $\beta = \omega_c \tau$, and the gas parameter β_0. The anomaly was observed in the case $\beta \ll 1$, $\beta_0 \ll 1$. Both the numerical simulations and the qualitative considerations [20] indicated that at zero temperature the MR can be expressed in terms of a dimensionless function $f(z)$ via

$$\frac{\delta\rho_{xx}}{\rho} = -\beta_0 f\left(\frac{\beta}{\beta_0}\right), \tag{3}$$

where ρ is the resistivity for $B = 0$. The analytical theory of the effect was developed in Ref. [21] where it was shown that the function $f(z)$ has the following asymptotics

$$f(z) = \begin{cases} 0.32z^2 & \text{for} \quad z \to 0 \\ 0.39 - 1.3/\sqrt{z} & \text{for} \quad z \to \infty, \end{cases} \tag{4}$$

and can be well approximated by linear function

$$f(z) \approx 0.032(z - 0.04), \tag{5}$$

in the interval $0.05 \lesssim z \lesssim 2$.

In spite of large number of publications, devoted to the study of the influence of the non-Markovian effects on the ρ_{xx}, the dependence of R on B induced by such effects was investigated (to the best of our knowledge) only in the context of „circling electrons“ [8]. It was found that though the existence of circling orbits leads to a strong dependence of ρ_{xx} on B in the region of classically strong B ($\omega_c\tau_{\mathrm{tr}} \gg 1$), the corresponding dependence of R on B is very weak in the whole range of B [8].

In this paper, we propose another mechanism of dependence of R on B. It does not rely upon the existence of non-colliding electrons but, in contrast, assumes that transport properties of colliding electrons are modified by classical ME. The mechanism turns out to be especially effective in the region of very weak fields, $\omega_c\tau_{\mathrm{tr}} \lesssim \beta_0$.

We will study dependence of R on B in 2D degenerated electron gas in a system of randomly located classical scatterers modeled by impenetrable disks of radius $a \gg \lambda_{\mathrm{F}}$, where $\lambda_{\mathrm{F}} = \hbar/mv_{\mathrm{F}}$ is the Fermi wavelength. The simplest realization of such a system is a quantum well with random array of antidots. We restrict ourselves to the case $\omega_c\tau_{\mathrm{tr}} \ll 1$. The electron dynamics is studied classically. The role of quantum effects is briefly discussed at the end of the paper.

We start with recalling that in the frame of the Boltzmann approach, the collision with a single scatterer is described by a differential scattering cross-section $\sigma(\theta)$ (see Fig. 1, a) and the collisions with different scatterers are independent. Inverting in time the process shown in Fig. 1, a we get a process shown in Fig. 1, a', corresponding to scattering by the angle $-\theta$. This implies an important property of a single scattering — the symmetry with respect to replacement of θ by $-\theta$ (reciprocity theorem): $\sigma(\theta) = \sigma(-\theta)$ [22]. This is the property which provides that R does not depend on B. If, for any reason, scattering cross-section acquires an asymmetric correction $\delta\sigma(\theta) \neq \delta\sigma(-\theta)$, the expression for ρ_{xy} becomes

$$\rho_{xy} = \frac{m(\omega_c + \Omega)}{e^2 n} = -B(R + \delta R), \tag{6}$$

where

$$\Omega = -n_0 v_{\mathrm{F}} \int d\theta\, \delta\sigma(\theta) \sin\theta, \qquad \frac{\delta R}{R} = \frac{\Omega}{\omega_c}. \tag{7}$$

In particular, such an asymmetric correction arises due to ME specific for processes of double scattering on a scatterer after return to it (see Fig. 1, b, b', c, c'). Though such processes are beyond the Boltzmann picture, they can be formally included into the kinetic equation by a slight modification of the Boltzmann collision integral. Specifically, one can introduce a small change of the scattering cross-section $\sigma(\theta) \to \sigma(\theta) + \delta\sigma(\theta)$ on

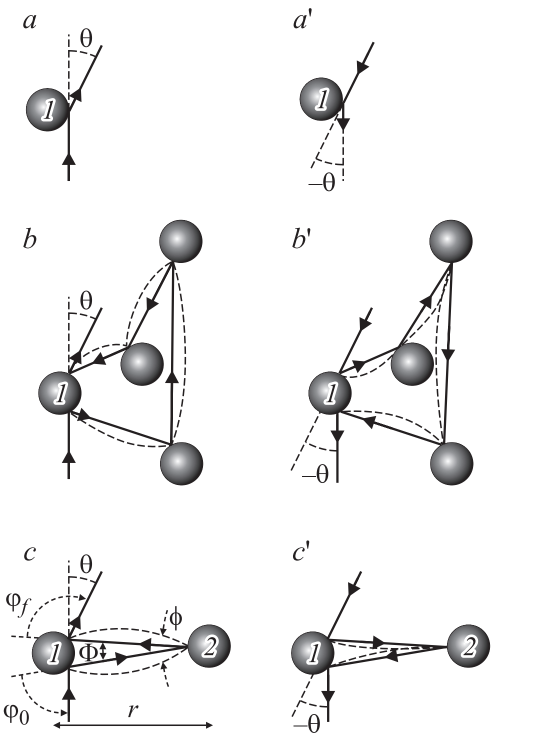

Figure 1. Processes of single scattering by angle θ (a) and $-\theta$ (a') characterized by a scattering cross-section $\sigma(\theta)$ ($\sigma(\theta)=\sigma(-\theta)$ both for $B=0$ and for $B\neq 0$), and processes of scattering on complexes of scatterers (b, b', c, c') including double scattering on the scatterer *1*. Correction to the cross-section due to multi-scattering processes remains symmetric for $B=0$. Magnetic field bends trajectories as shown in b, b', c, c' by dashed lines. As a result, the symmetry with respect to inversion of θ is broken, so that $\delta\sigma(\theta)\neq\delta\sigma(-\theta)$ for $B\neq 0$.

the disk where double scattering takes place (disk *1* in Fig. 1, b, b', c, c') [21,23]. For $B=0$, cross-section remains symmetric: $\delta\sigma(\theta)=\delta\sigma(-\theta)$. However, for $B\neq 0$ the time inversion symmetry is broken, so that the cross-section becomes asymmetric: $\delta\sigma(\theta)\neq\delta\sigma(-\theta)$. The point is that the influence of the magnetic field is different for the processes where closed return path is passed counterclockwise (Fig. 1, b, c) and clockwise (Fig. 1, b', c').

The return after one scattering (see Fig. 1, c, c') needs special attention because the probability of such a process very sharply depends on B due to „empty corridor effect" (ECE) [20,21]. The mechanism of this phenomenon proposed in Ref. [20] is linked to the memory effects arising in backscattering events. It has a close relation to the well known non-analyticity of the virial expansion of transport coefficients [24–28] which we briefly recall. For $B=0$ the leading nonanalytic correction to resistivity, $\delta\rho$, is due to the processes of return to a scatterer after a single collision on another scatterer (see Fig. 2, a). The relative correction, $\delta\rho/\rho$, is proportional to the corresponding backscattering probability, given by the product of $e^{-r/l}d\Phi dr/l$ (which is the probability to reach scatterer *2* without collision and scatter in the angle $d\Phi$) and the probability p to return without collisions from *2* to *1* (here l is the mean free pass). Assuming $p=\exp(-r/l)$ and integrating over intervals $0<\Phi<a/r$, $a<r<\infty$, one obtains [24–28]

$$\delta\rho/\rho\sim\int\limits_a^\infty\frac{dr}{l}\int\limits_0^{a/r}d\Phi\,e^{-2r/l}\sim\beta_0\ln(1/2\beta_0). \tag{8}$$

In Ref. [20] it was shown that the probability p is actually larger than $\exp(-r/l)$. Indeed, the exponent $\exp(-r/l)$ can be written as $\exp(-nS)$, where $S=2ar$. It represents the probability of the existence of an empty corridor (free of the centers of the disk) of width $2a$ around the electron trajectory from *2* to *1*. However, the passage of a particle from *1* to *2* ensures the absence of the disks centers in the region of width $2a$ around this part of trajectory (from *1* to *2*). This reduces the scattering probability on the way back. The correct value of p can be estimated as

$$\begin{aligned}p(R,\Phi)&=\exp\big[-n(S-S_0)\big]\\&=\exp\big[-r/l+nS_0(r,\Phi)\big],\end{aligned} \tag{9}$$

where

$$S_0(r,\Phi)=2ar-r^2|\Phi|/2 \tag{10}$$

is the area of the overlap of the two corridors (see Fig. 2, a). For example, for $\Phi=0$, we have $S_0=2ar$ and $p=1$, which reflects the obvious fact that the particle cannot scatter if it travels back along the same path. Taking into account the effect of „empty corridor", we get

$$\frac{\delta\rho}{\rho}\sim\int\limits_a^\infty\frac{dr}{l}\int\limits_0^{a/r}d\Phi\,e^{-(2r/l)+nS_0}\approx\beta_0\ln\left(\frac{C}{2\beta_0}\right), \tag{11}$$

where C is a constant of the order of unity. Thus, for $B=0$ the „empty corridor" effect simply changes the constant in the argument of the logarithm.

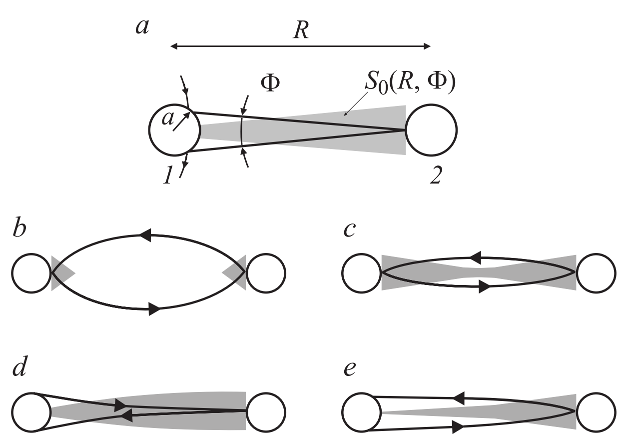

Figure 2. a — backscattering process responsible for leading nonanalytic contribution to the resistivity at $B=0$. b — for $B\neq 0$, the overlap area, S_B, between two corridors is small at large B. c — for $\Phi=0$, S_B decreases with B. d — for $\Phi\neq 0$ and small B, the values of S_B-S_0 for time reversed trajectories have opposite signs.

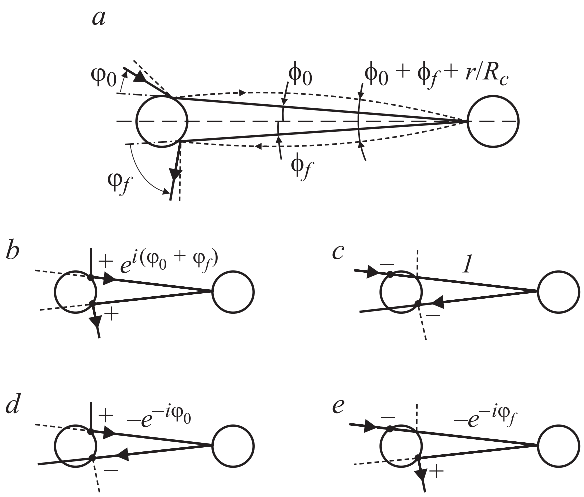

Figure 3. Backscattering process is parameterized by the angles φ_0, φ_f. The magnetic field changes the backscattering angle $\phi = \phi_0 + \phi_f + r/R_c$. *a*: the solid (dashed) line represents electron trajectory for $B = 0$ ($B \neq 0$). *b–e* — different processes contributing to cross-section renormalization.

The key idea suggested in Ref. [20] was that for $B \neq 0$ the area of the overlap of the two corridors, S_B, sharply depends on B, resulting in the observed MR. Indeed, it is seen from Fig. 2,*b* that for $\beta \gtrsim \beta_0$ $S_B \to 0$ resulting in sharp negative MR

$$\frac{\delta\rho_{xx}}{\rho} \sim \int_0^\infty \frac{dr}{l} \int_0^{a/r} d\phi\, e^{-2r/l}(e^{nS_B} - e^{nS_0}). \tag{12}$$

The following qualitative explanation of the observed linear MR was presented in Ref. [20]. The value $n(S_B - S_0)$ was estimated for $\phi = 0$ (see Fig. 2,*c*) to the first order in B as $-nr^2/R_c = -r^3/2alR_c$, where R_c is the cyclotron radius. Assuming that this estimate also works at $\phi \neq 0$ and expanding $e^{nS_B} - e^{nS_0}$ to the first order in B, one gets $\delta\rho_{xx}/\rho \sim -l/R_c = -\omega_c\tau$.

In fact, the physical picture of the phenomenon is more subtle. The contribution of any trajectory with $\phi \neq 0$ is cancelled to the first order in B by the contribution of the time-reversed trajectory, since the values of $S_B - S_0$ are opposite for these paths (see Fig. 2,*d, e*). The cancellation does not occur only at very small $\phi \sim \beta$. The integration in Eq. (12) over $\phi < \beta$ yields $\delta\rho_{xx}/\rho \sim -\beta^2/\beta_0$. Larger values of ϕ also give a quadratic in β contribution to the MR. This contribution is positive and comes from the second order term in the expansion of $e^{nS_b} - e^{nS_0}$ in B. A more rigorous approach [21] demonstrated that the contribution of small angles is dominant resulting in a negative parabolic MR and that the parabolic MR crosses over to linear at very small $\beta \approx 0.05\beta_0$, which explains why the parabolic MR was not seen in numerical simulations [20] (see Eqs. (4), (5)).

The calculation of R is quite analogous to the calculation of ρ_{xx} presented in [21]. As was shown in this paper, the correction to cross-section arises due to four scattering processes (see Fig. 3). In the process $(+,+)$ (Fig. 3,*b*) an electron has two real scatterings on a disk placed at point $\mathbf{r}$. The process $(-,\ -)$ (Fig. 3,*c*) does not correspond to any real scattering at point $\mathbf{r}$. It just allows us to calculate correctly the probability for an electron to pass twice the region of the size a around point $\mathbf{r}$ without scattering. To interpret the process $(+,-)$, note that in the Boltzmann picture, which neglects correlations, electron can scatter on a disk and later passes through the region occupied by this disk without scattering (Fig. 3,*d*). The $(+,-)$ correction to the cross-section modifies the Boltzmann result by substracting the contribution of such unphysical process. Analogous consideration is valid for the process $(-,+)$ shown in Fig. 3,*e*. The rigorous method of calculation $\delta\sigma(\theta)$ accounting for both four processes was developed in Ref. [21]. The calculations yield

$$\delta\sigma(\theta) = \frac{1}{4l}\int_a^\infty \frac{dr}{r}\, e^{-2r/l} \int_0^{2\pi} d\varphi_0 \int_0^{2\pi} d\varphi_f\, \sigma(\varphi_0)\sigma(\varphi_f) e^{n_0 S_B}$$
$$\times \left[\delta(\theta - \varphi_{\varphi_0,\varphi_f}) + \delta(\theta-\pi) - \delta(\theta - \varphi_{\varphi_0,0}) - \delta(\theta - \varphi_{0,\varphi_f})\right]. \tag{13}$$

Here $\varphi_{\varphi_0,\varphi_f} = (\pi + \varphi_0 + \varphi_f)(\mathrm{mod}2\pi)$, $\sigma(\varphi) = (a/2) \times |\sin(\varphi/2)|$ is the single scattering cross-section,

$$S_B = \int_0^r dr'(2a - |\phi r' - r'^2/R_c|)\Theta\left[2a - |\phi r' - r'^2/R_c|\right],$$

$\Theta[\cdots]$ is the Heaviside step function, $\phi = \Phi + r/R_c$, R_c is the cyclotron radius and $\Phi \approx (a/r)[\cos(\varphi_0/2) + \cos(\varphi_f/2)]$ (see Fig. 1,*c*). Four terms $[\delta(\theta - \varphi_{\varphi_0,\varphi_f}) + \delta(\theta-\pi) - \delta(\theta - \varphi_{\varphi_0,0}) - \delta(\theta - \varphi_{0,\varphi_f})]$ in Eq. (13) correspond to four types of non-Markovian processes shown in Fig. 3. Introducing dimensionless variables $T = r/l$, $z = \omega_c\tau/\beta_0$ and using Eq. (7), we get

$$\frac{\delta R}{R} = g(z)$$
$$= -\int_0^\infty \frac{dT}{T}\, e^{-2T} \int_0^\pi d\alpha \int_0^\pi d\gamma\, \sin(\alpha+\gamma)\sin^2\alpha\sin^2\gamma\, \frac{e^{s_z} - e^{s_0}}{2z}. \tag{14}$$

Here

$$s_z = \int_0^T dt\left(1 - \left|\xi t - \frac{zt^2}{2}\right|\right)\Theta\left(1 - \left|\xi t - \frac{zt^2}{2}\right|\right),$$

$\xi = (\cos\alpha + \cos\gamma)/2T + zT/2$, $s_0 = s_{z\to 0}$. Function $g(z)$ calculated numerically with the use of Eq. (14) is plotted in Fig. 4. For $z \ll 1$, $g(z) \approx 0.064 - 4z^2$. For $z \gg 1$, $g(z)$ decreases as $0.35/z^{3/2}$. It worth emphasizing that $\delta R/R \sim 1$ for $z \lesssim 1$. This means that the correction is not parametrically small in a gas parameter β_0 which is usually considered as expansion parameter for ME-induced corrections.

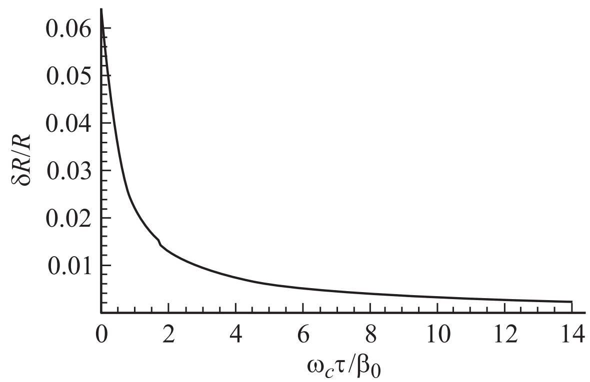

Figure 4. Magnetic field dependence of the relative correction to the Hall coefficient caused by ECE.

Next we calculate δR for stronger fields, $\beta_0 \ll \omega_c\tau_{\rm tr} \ll 1$. At such fields empty corridor effect is suppressed and returns after one scattering (Fig. 1, c, c') and after a number of scatterings (Fig. 1, b, b') equally contribute to δR. In this case, one can also introduce the effective scattering cross-section [23] which turns out to be frequence-dependent and for $\omega = 0$ reads[1]

$$\delta\sigma(\theta-\theta') = v_{\rm F}\int \big[\sigma(\theta-\varphi) - \sigma_0\delta(\theta-\varphi)\big]\tilde{G}(0,\varphi-\varphi') \times \big[\sigma(\varphi'-\theta') - \varphi_0\delta(\varphi'-\theta')\big]d\varphi d\varphi'. \quad (15)$$

Here $\sigma_0 = \int d\varphi\sigma(\varphi)$ is the total cross-section for single scattering, $\tilde{G}(0,\varphi-\varphi') = \tilde{G}(\mathbf{r},\varphi,\varphi')|_{\mathbf{r}\to 0}$, $\tilde{G}(\mathbf{r},\varphi,\varphi') = G(\mathbf{r},\varphi,\varphi') - G^{\rm ball}(\mathbf{r},\varphi,\varphi')$, $G(\mathbf{r},\varphi,\varphi')$ is the Green function of the stationary Boltzmann equation,

$$G^{\rm ball}(\mathbf{r},\varphi,\varphi') = \frac{\exp(-\theta_r/\beta)}{v_{\rm F} r\cos(\theta_r/2)}\,\delta(\varphi-\varphi_{\mathbf{r}}+\theta_r/2) \times \delta(\varphi'-\varphi_{\mathbf{r}}-\theta_{\mathbf{r}}/2)$$

is the Green function of the Boltzmann equation without in-scattering term, $\varphi_{\mathbf{r}}$ is the angle of vector $\mathbf{r}$, and $\theta_r = 2\arcsin(\beta r/2l)$. Substituting Eq. (15) into Eq. (7) and using the property $\int d\varphi d\varphi_0 G(0,\varphi,\varphi')\sin(\varphi-\varphi') = 0$,[2] we get after some algebra

$$\frac{\delta R}{R} = -\frac{n_0\sigma_{\rm tr}^2}{2\pi} \ll 1. \quad (16)$$

where $\sigma_{\rm tr} = \int d\theta\sigma(\theta)(1-\cos\theta) = 8a/3$. Hence, with increasing B relative correction decreases according to Eq. (14), then changes sign and saturates at small negative values. It is noteworthy that, as follows from the above derivation, Eq. (16) is valid not only for the case of impenetrable disks but also for any type of well-separated scatterers.

Above we discussed an idealized system where only strong scatterers are present. Let us now assume (see [29–31]) that in addition to strong scatterers there is a weak smooth random potential $U(\mathbf{r})$ with the rms amplitude U and the correlation length d ($a \ll d \ll l$). The presence of such a potential does not influence the ECE provided that $\Lambda \gg l$, where $\Lambda \sim d(E_{\rm F}/U)^{2/3}$ is the Lyapunov length, characterizing the divergence of the electron trajectories in the potential $U(\mathbf{r})$. In the opposite limit, $\Lambda \ll l$, one should restrict integration over r in Eq. (13) by Λ. In this case, relative correction to R decreases: $\delta R/R \sim \Lambda/l$. On the other hand, the field needed for suppression of the ECE increases and can be found from the following estimate $\omega_c\tau_{\rm tr} \sim \beta_0(l/\Lambda)^2$ (at such a field two corridors corresponding to passage *1* $\to$ *2* and *2* $\to$ *1* (see Fig. 2, a) between disks *1* and *2* separated by a distance $r \sim \Lambda$ cease to overlap). One can also show, that at stronger fields the effect of smooth disorder leads to appearing of a very weak parabolic dependence of R on B: $\delta R/R \sim -(\omega_c\tau)^2(d/l)^2(U/E_{\rm F})^2$.

Finally, we briefly discuss the role of quantum effects. As was mentioned in Ref. [21], for $\lambda_{\rm F} > a^2/l$ the corridor effect is suppressed by diffraction on the disk's edges. In this case, the integration over r in Eq. (13) is limited by $\Lambda' = a^2/\lambda_{\rm F}$. Hence, $\delta R/R \sim \Lambda'/l < 1$, and the field needed for suppression of the ECE can be found from the following estimate $\omega_c\tau_{\rm tr} \sim \beta_0(l/\Lambda')^2$.

To conclude, we have shown that the classical memory effects might strongly renormalize Hall coefficient. The most interesting phenomena arise due to empty corridor effect which leads to a sharp field dependence of return probability and, in turn, results in a sharp dependence of R on B. The analytical calculation of dependence $R(B)$ was presented for 2D system with random array of antidots modeled by hard-core spherical scatterers. It was demonstrated that empty corridor effect leads to a very sharp dependence of R on B concentrated in the region of very weak fields ($\omega_c\tau \lesssim a/l$). The total variation of R in this region of fields is on the order of the Boltzmann value of R. At larger fields, where $a/l \ll \omega_c\tau_{\rm tr} \ll 1$, the ME lead to a small field-independent correction to R and (in a presence of smooth disorder) to a very weak parabolic dependence.

We acknowledge I. Gornyi and G. Minkov for fruitful discussion. The work was supported by RFBR, by grant of Russian Scientific School, and by programmes of the RAS. V.Yu.K. was supported by Dynasty foundation.

[1] Eq. (15) is obtained by integration over time Eq. (13) of [23] and extracting contribution of ballistic term. The latter one represents propagation without collision and, evidently, can not give any contribution to return processes.

[2] To obtain this property we first integrate the Boltzmann equation (in q-space) over φ which yields $\int G_{\mathbf{q}}(\varphi,\varphi')i\mathbf{q}\mathbf{v}d\varphi = 1$. One can show that integration over φ' yields the same result $\int G_{\mathbf{q}}(\varphi,\varphi')i\mathbf{q}\mathbf{v}'d\varphi' = 1$. Using these identities we get $\int d\varphi d\varphi' d\mathbf{q} G_{\mathbf{q}}(\varphi,\varphi')\sin(\varphi-\varphi') = \int d\varphi\, d\varphi' d\,\mathbf{q}\, G_{\mathbf{q}}(\varphi,\varphi')\,[\sin(\varphi-\varphi_{\mathbf{q}})\cos(\varphi_{\mathbf{q}}-\varphi') + \cos(\varphi-\varphi_{\mathbf{q}}) \times \sin(\varphi_{\mathbf{q}}-\varphi')] = 0$.

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Редактор Л.В. Шаронова