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2D lattice of coupled Sinai billiards: metal or insulator at $g \ll 1$?

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Abstract

We investigate the transport in a two-dimensional (2D) lattice of coupled Sinai billiards fabricated on the basis of a high-mobility 2D electron gas in a GaAs/AlGaAs heterojunction in going from $g > 1$ to $g \ll 1$ (where $g = \sigma/(e^2/h)$ is dimensionless conductivity). The data obtained suggest that the system studied behaves more like a metal than an insulator at $g \ll 1$ and is not described by the generally accepted picture of metal–insulator transition in 2D electron systems. © 1998 Published by Elsevier Science B.V. All rights reserved.

Keywords: Magnetotransport; Metal–insulator transition; Electron billiards; Heterojunction

Quantum and classical transport in systems with dynamic chaos has been intensively studied during the last few years, since the successes of modern semiconductor technology have made it possible to obtain various experimental realizations of such systems with electron billiards as an example. At the present time two varieties of these systems are studied. The first one is unit billiards (regular or chaotic Bunimovich or Sinai billiards) [1,2], while the second one is macroscopic two-dimensional (antidot lattice) [3–7] or one-dimensional [8] Sinai billiards.

We report for the first time the results of experimental investigation of the new system with dynamic chaos. The system was built on the basis of a two-dimensional (2D) lattice of closely situated antidots fabricated from a high mobility 2D electron gas in a GaAs/AlGaAs heterojunction with a metallic gate evaporated on the top, that permitted us to control the conductivity of the structure in a sufficiently wide range from $g = 0.01$ to $g = 2$.

The square lattice of antidots was fabricated on the basis of a two-dimensional electron gas with electron density $N_S = (2\text{–}3)\cdot 10^{11}$ cm^{-2} and mobility $\mu = (3\text{–}8)\cdot 10^5$ cm^2/V s corresponding to the mean free path $l = 3\text{–}6$ μm by means of electron lithography and subsequent plasma etching. Then the NiAu or TiAu gate was evaporated on the top of the device. We investigated three samples with the lattice period $d = 0.6$ μm and three ones with $d = 0.7$ μm. The lithographic size of antidots $a = 0.2$ μm was the same for all samples. However, due to the depletion layers the actual size was larger, being approximately equal to the lattice period even before the gate evaporation. The experimental

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sample was made up of two Hall bars of the length of 100 μm and the width of 50 μm.

Fig. 1 shows the set of temperature dependences of the conductivity $g(T)$ for the sample AG219 with the period $d=0.6$ μm ($\mu=7\cdot10^5$ cm^2/V s at $N_S=2\cdot10^{11}$ cm^{-2} in the unpatterned part of the sample) for different values of the gate voltage. It is seen that for $g>1$ conductivity is practically temperature independent. More exactly, weak logarithmic decrease of g is observed typical for the weak localization effects. It becomes more noticeable for lower values of conductivity, but it still remains weak even for $g \ll 1$, and is well described by the power law dependence $g(T) \propto T^{\alpha}$ with $\alpha<1$ for all of the tested samples. Specifically, for the dependences of Fig. 1(a) $\alpha=0.1-0.27$ at $g \ll 1$. The behavior of $g(T)$ described above is significantly different from that for the unpatterned 2D electron gas both in silicon MOS-transistors [9] and in AlGaAs/GaAs heterojunctions [10], as well as for antidot lattices with short period [11,12]. In all of these cases at $g \sim 1$ the transition from the weak logarithmic dependence (weak localization regime) to the strong exponential one (strong localization regime) is observed. In our case the weak logarithmic decrease of g (for $g>1$) is followed by a weak power law, that has not been observed in other 2D systems before.

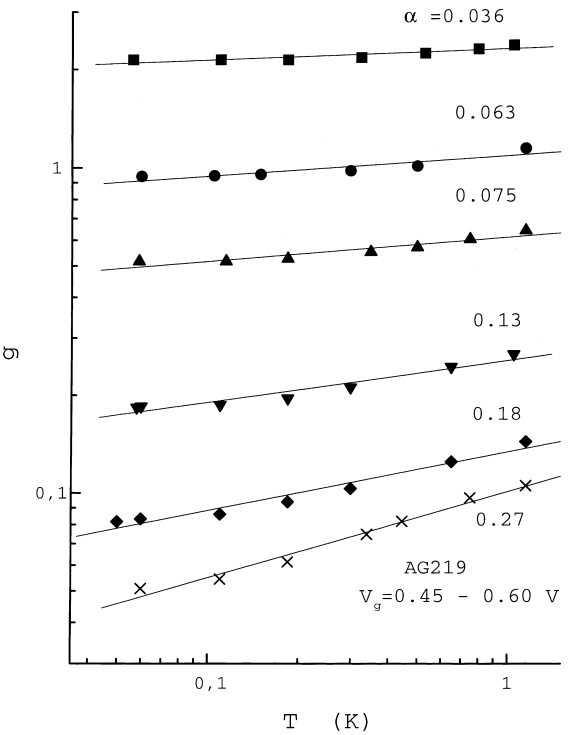

Fig. 1. Temperature dependences of the dimensionless conductivity at the transition from $g>1$ to $g \ll 1$. Solid lines are $g \sim T^{\alpha}$.

Consider the influence of magnetic field. It is well known that in AlGaAs/GaAs heterojunctions at $g<1$ under the influence of magnetic field a transition from an insulator to the quantum-Hall-liquid is observed. This transition is characterized by the critical point B_c and $g_c \approx 0.5$–1 [10]. Our samples exhibit a radically different behaviour. It is seen from Fig. 2 that for all values of $g(B=0)$ in the magnetic fields about $B \approx 1$ T the transition takes place from weak power-law dependence $g(T)$ to no temperature dependence at all. Moreover, this transition is of different kind, for there is no critical point, and the metallic behavior extends for $g \ll 1$. Fig. 2(a) also shows an interesting picture in weak magnetic field. For all values of g, the negative magnetoresistance (NMR) is observed for $B<0.05$ T, followed by a peak at $B \approx 0.2$ T corresponding to the condition $2R_c=d$. This peak is well-known for the antidot lattices at $g>1$ and originates from the so-called pin-ball trajectory that surrounds an antidot not colliding with it. It is seen from Fig. 2(b) that the second commensurate peak is observed at $g \approx 1$. It corresponds to the condition $2R_c=(\sqrt{2}-1)d$. The positions of the commensurability peaks shows that we really deal with the lattice of closely situated antidots with $d \approx a$ and $d \gg d-a$. The Sinai billiards between the antidots have the area $S=d^2\ (1-\pi/4)$ and contain a large number of electrons $N \gg 1$. In our case we have correspondingly $S=0.5$ μm^2, $N \approx 70$ for $d=0.7$ μm, and $S=0.36$ μm^2, $N \approx 50$ for $d=0.6$ μm. It is also important that both the main commensurability peak and NMR are conserved at the transition from $g>1$ to $g \ll 1$.

The behavior of NMR is shown in Fig. 3 in more detail. It is characterized by two distinguishing features: (i) for all states with $0.05<g<2$ NMR is cut off at the same magnitude of magnetic

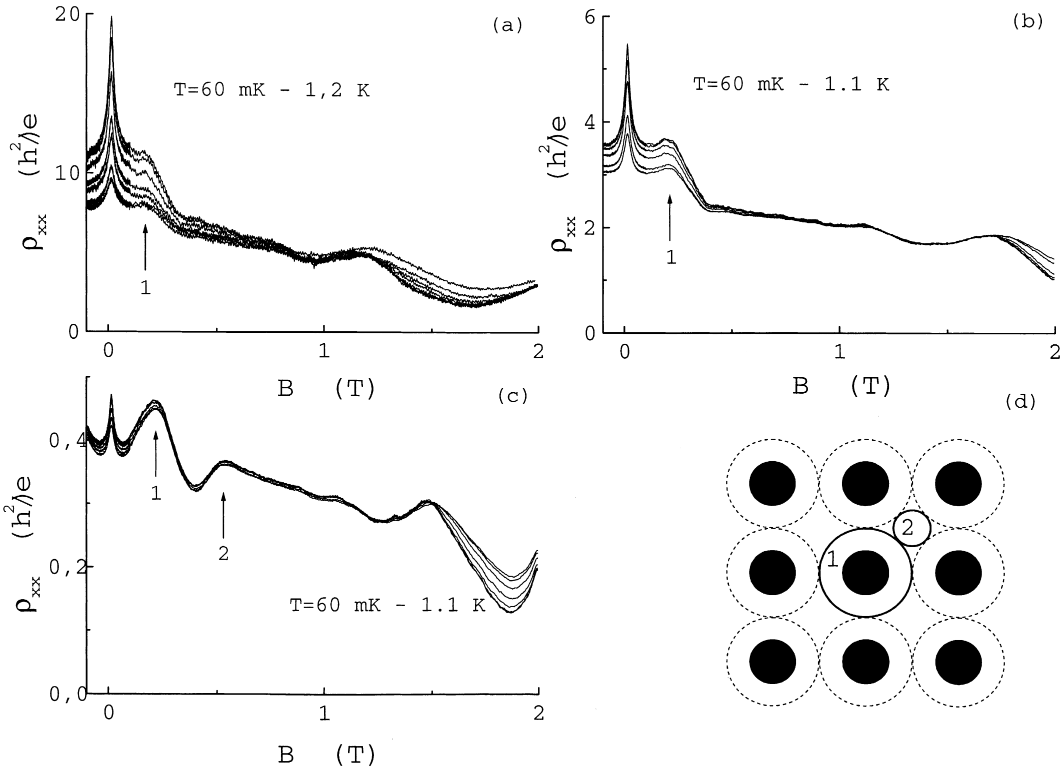

Fig. 2. Magnetoresistance traces $\rho_{xx}(B)$ ($\rho_{xx} = 1/g$) for different values of the conductivity and at different temperatures: (a) $g = 0.05$, (b) $g = 0.18$, (c) $g = 2.2$. (d) Schematic view of antidot lattice (black points show the etched regions, broken lines show the boundary of depletion region, 1 – an electron trajectory around antidot, 2 – an electron trajectory between antidots).

field $B \approx 0.05$ T; (ii) NMR noticeably increases with decreasing g (Fig. 3(a) shows that at $g = 0.05$ it reaches a considerable value about 40%), and its temperature dependence becomes stronger. For the states with the highest resistivity it was more stronger than $g(T)$). For $g > 1$ NMR can be attributed to the effects of weak localization in open chaotic billiards [2], because it has relatively small amplitude and Lorentzian line-shape. The behavior of NMR for $g \ll 1$ is surprising. It increases by an order of magnitude reaching a considerable value comparable to the total resistance of the sample, while the line-shape of NMR is described by a Lorentz curve of the same width. The width is equal to $\Delta B_1 = 27 \pm 2$ mT for the sample AG219. It corresponds to a half magnetic flux quantum through the area of the billiard that is equal to d^2 $(1 - \pi/4)$. The fact that the width is determined by the magnetic flux quantum is well seen from the comparison of NMR for the samples with two different periods 0.6 and 0.7 μm. As it is clearly seen from the Fig. 3(c) the width of NMR curve for the period 0.7 μm equals $\Delta B_2 = 20 \pm 2$ mT, that is $\Delta B_1/\Delta B_2 = (0.7)^2/(0.6)^2$. Thus at $g \ll 1$ we observe NMR which is very similar to weak localization NMR in chaotic open billiards [2].

Let us discuss the results obtained. First turn to the temperature dependence at the transition from $g > 1$ to $g \ll 1$. It differs from the accepted picture of the metal–insulator transition in 2D electron systems. This picture is based on the concept of Anderson localization of electrons, be it the model of minimal metallic conductivity (MMC) or the scaling theory (ST) [13]. According to it at $k_F l \sim 1$ or $g \sim 1$ in a macroscopic 2D system the transition should occur from the metallic behavior (as in MMC model) or from the weak localization (as in ST) to the strong localization characterized by

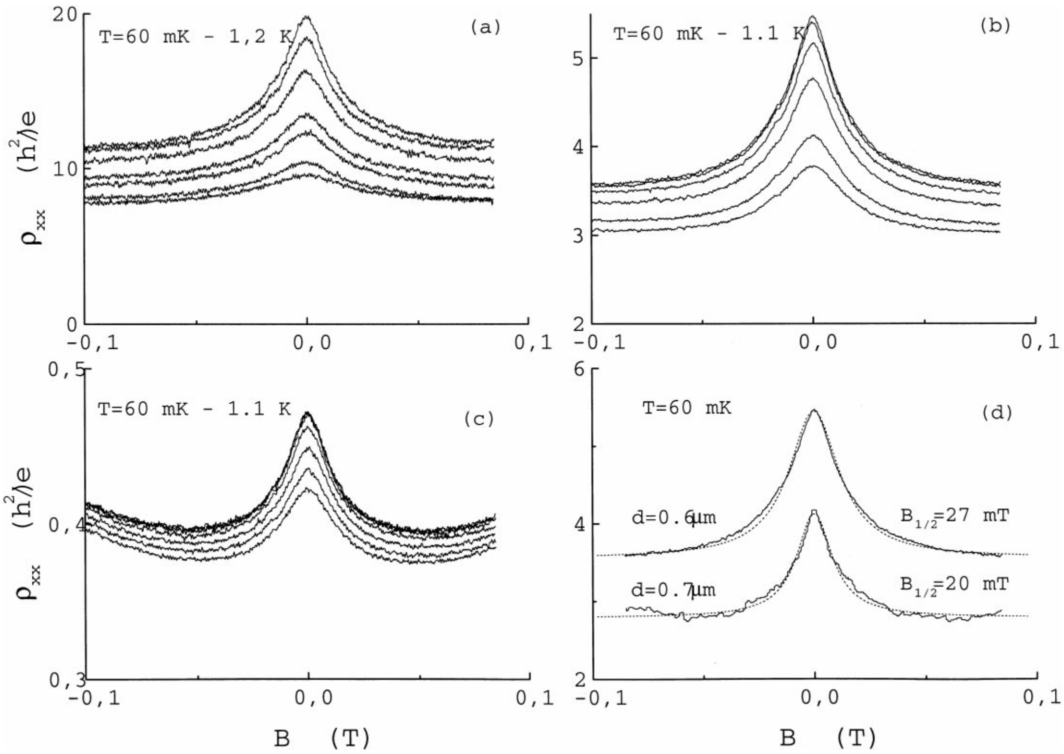

Fig. 3. (a,b,c) NMR curves for the sample. (d) Experimental and calculated NMR curves for the antidot lattices with two different periods $d = 0.6$ μm and $d = 0.7$ μm. Solid lines are the experimental curves, broken lines are the calculated Lorenz ones, $B_{1/2}$ is the width of the Lorenz curves.

the exponential temperature dependence of the activation type (hopping conductivity) or of the Mott type (variable range hopping conductivity). In real systems the transition can be quite complicated, but for $g \ll 1$ the state with hopping conductivity is always realized [9,10]. Recently in [14] the model of metal grains with $g \gg 1$ coupled by tunneling in a way that the conductivity of the macroscopic sample was low $g \ll 1$ was considered theoretically. At $g \ll 1$, due to inelastic electron-electron scattering the linear dependence $g(T)$ was obtained in a wide temperature range. The dependence changes to the exponential one at $T \ll e^2/C$, where C is the capacitance of a metal grain. Our results are not in agreement with the predictions of this theory. Firstly, in our case $g(T)$ is weaker than linear. Secondly, it is observed at $T \ll e^2/C$, because the Coulomb energy for our samples $e^2/2C \approx 15$–25 K. This means that the low conductivity of the lattice of the Sinai billiards cannot be explained by the model of metal lakes coupled by weak tunnelling junctions. The effects of Coulomb blockade are not manifested in the experiment, because one observes no features in $g(V_g)$ and $g(V_d)$ (where V_d is voltage applied to the current probes) dependences. So we have to assume that at $g \ll 1$ the coupling between the billiards is stronger than that provided by tunneling. The behavior of the lattice in the magnetic field supports this assumption. In Fig. 2(a) one can see the Shubnikov–de Haas oscillations, which give the electron density in the lattice saddle points connecting the billiards. It coincides with the density determined from the Hall effect that should give the concentration in these points. The electron density in these points weakly changes with the strong change of g. It decreases only by about 30% while g drops by a factor of 30, and its magnitude, equaled to $4.1 \cdot 10^{10}$ cm^{-2}, is only three times less than N_S inside the billiard even for the state with the lowest value of g. This means that the Fermi level in the saddle point lies several meV above the barrier at $g \ll 1$, and an electron should ballistically move through the "bottle neck" to go from one billiard to another. The behavior of commensurability peak and of NMR support this picture. Hence, we should come to a paradoxical conclusion that 2D lattice of Sinai billiards coupled via conducting bottle-neck can have very low conductivity $g \ll 1$ but simultaneously exhibits the properties typical for metallic ballistic systems rather than for the insulators. This conclusion is in a drastic contradiction with the standard picture of metal–insulator transition in 2D systems and presents a challenge to the modern theory of quantum transport in the condensed matter.

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