Observation of large $h/2e$ oscillations in semiconductor antidot lattices

F. Nihey, S. W. Hwang, and K. Nakamura
Fundamental Research Laboratories, NEC Corporation, 34 Miyukigaoka, Tsukuba, Ibaraki 305, Japan
(Received 25 October 1994)

We report the observation of large-amplitude oscillations in the magnetoresistance of a two-dimensional electron gas in a GaAs-$Al_xGa_{1-x}As$ heterostructure with an artificially imposed antidot lattice potential. The period of the oscillations is about $h/2e$ as a function of the magnetic flux threading through the unit cell of the antidot lattice. The oscillations are definitely visible in the antidot array with a hexagonal lattice configuration, the amplitude reaching 20% of the sample resistance, but almost invisible with a square lattice configuration. We believe that the hexagonal lattice potential is suitable for electrons to return to their starting points and to contribute to coherent backscattering. We also show the oscillations cannot be described by the theory for the diffusive regime established by Al'tshuler, Aronov, and Spivak.

The resistance of a conductor in the shape of a hollow cylinder oscillates as a function of the magnetic flux threading through the hollow with a period of $h/2e$. This was predicted by Al'tshuler, Aronov, and Spivak[1] (AAS) for the diffusive regime where the mean free path of the electrons is much smaller than the sample size. The conductance amplitude of the oscillations is of the order of e^2/h and depends on the phase coherence length over which an electron maintains its phase coherence. The oscillations are caused by coherent backscattering of an electron, where a pair of backscattered partial waves with time-reversal symmetry interfere with each other. This prediction has been proved experimentally by using cylindrical metal films[2–5] and other geometric structures such as networks.[6–8] Those experiments were done using polycrystalline metal where electrons move diffusively due to impurity scattering. The amplitudes of the oscillations reported so far have been very small, i.e., less than 0.1% of the sample resistance.

Negative magnetoresistance, which has been observed in the diffusive regime, is also caused by the coherent backscattering of electrons. A recent experiment demonstrated[9] that the negative magnetoresistance emerges even in ballistic microstructures. This can be explained by the semiclassical theory of quantum billiards,[10] where chaotic electron trajectories contribute to the interference. From this theory one may expect that $h/2e$ oscillations should also occur in the magnetoresistance of ballistic microstructures in the shape of a cylinder or network. However, there have been very few experiments showing the existence of the $h/2e$ oscillations in the ballistic regime.

Recently, a two-dimensional electron gas (2DEG) in a GaAs-$Al_xGa_{1-x}As$ heterostructure with an artificially imposed two-dimensional array of potential peaks has attracted much attention. In the antidot lattice,[11] the periodic potential is so strong that it forms a two-dimensional array of small depletion regions in the 2DEG plane.[12] When the period of the antidots is smaller than the mean free path of electrons, a sequence of peaks appears in the magnetoresistance.[13] The origin of the peaks has been explained with a classical pinball model where the formation of localized orbits enclosing a set of antidots increases the magnetoresistance. When the antidot system is coherent, i.e., the antidot period is smaller than the phase coherence length of electrons, quantum mechanical phenomena occur in the antidots. On top of the magnetoresistance peak, fine oscillations appear with a period of about h/e as a function of the magnetic flux threading through the unit cell of the lattice.[14,15] These oscillations are caused by the quantization of the localized orbits encircling a single antidot.

In this paper, we report the observation of large-amplitude $h/2e$ oscillations in the low-field magnetoresistance of GaAs-$Al_xGa_{1-x}As$ antidot lattices. Because the antidots have a large diameter compared to the antidot period, conducting channels in between the antidots form a honeycomb network. Although this network has a geometry similar to those made of metal where $h/2e$ oscillations have been observed,[6–8] it is apparent that our network is in the ballistic regime rather than the diffusive one.

An experiment has been performed using antidots with a small diameter compared to their period and with a square lattice configuration.[16] The magnetoresistance showed faint anomalies, suggesting the existence of the $h/2e$ oscillations in the ballistic regime. However, this needs verification because those anomalies are very weak and easily distorted by background changes of the magnetoresistance. On the contrary, we have succeeded in inducing definite $h/2e$ oscillations whose amplitude reaches 20% of the sample resistance. One of the key points to this success is the effective squeezing of electron channels with an array of antidots that are large compared to their period. Furthermore, we have found that the amplitude of the oscillations largely depends on the antidot lattice configuration. Large oscillations appear in a hexagonal lattice whereas a square lattice shows almost no oscillations. We attribute the $h/2e$ oscillations to the manifestation of the coherent backscattering of *ballistic* electrons that are steered by the hexagonal antidot lattice potential. We show that the antidot potential with the hexagonal lattice configuration is suitable for electrons to return to their starting points.

The samples were fabricated from a modulation-doped GaAs-$Al_{0.3}Ga_{0.7}As$ heterostructure grown by molecular-beam epitaxy with a 1-μm GaAs buffer layer on top of a semi-insulating GaAs substrate, a 300-Å undoped $Al_{0.3}Ga_{0.7}As$ spacer, a 600-Å Si-doped $Al_{0.3}Ga_{0.7}As$ layer, and a 50-Å Si-doped GaAs cap layer. The density of the two-dimensional electron gas was 2.8×10^{11} cm^{-2} and the mobility was 1.2×10^6 $cm^2/V\,s$ at $T = 1.5$ K in the dark. The mean free path determined from these values was 10 μm. Following the fabrication of Hall bars with a channel width of 16 μm and a distance between the potentiometric probes of 16 μm, a two-dimensional array of shallow holes was patterned and etched (see the insets in Fig. 1) in the whole area of the Hall bars using electron-beam lithography and wet chemical etching. Due to the reduction of the Si-doped layer thickness, small regions beneath these holes are depleted. Antidot arrays with a square lattice configuration as well as a hexagonal lattice configuration were prepared for comparison. The period (a) of both lattices was 200 nm and the diameter of the etched holes of both lattices was about 90 nm after etching. These length scales are much shorter than the mean free path in the starting material and possibly shorter than the phase coherence length [estimated to be ~ 1 μm at 1 K (Ref. 17)]. The magnetoresistance ρ_{xx} and the Hall resistance ρ_{xy} of these samples were obtained using a low-frequency lock-in technique (17–75 Hz) with an excitation current of 1–10 nA.

Magnetoresistance traces ρ_{xx} as a function of the perpendicular magnetic field B are shown in Fig. 1 for square and hexagonal antidot lattices. These data were obtained at $T = 1.5$ K in the dark. At high fields ($B > 2.4$ T for the square lattice and $B > 3$ T for the hexagonal lattice) these traces show the quantum Hall profile where

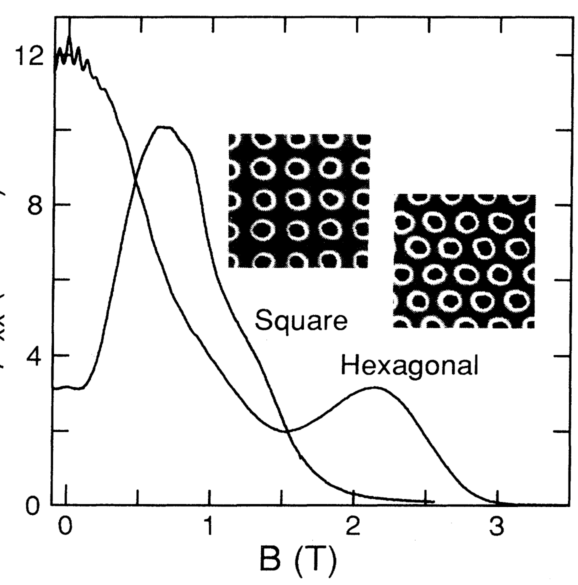

FIG. 1. Magnetoresistance traces for a hexagonal and a square lattice of antidots measured at $T = 1.5$ K in the dark. Each sample has an antidot period of $a = 200$ nm and almost the same antidot diameter. The insets are scanning electron microscope photographs of the antidot lattices.

each ρ_{xx} drops to zero and ρ_{xy} (not shown) is quantized to $h/2e^2$, indicating that there are few impurities in the conducting region. At low fields each trace shows an enhanced profile compared to the starting material. The magnetoresistance of the square lattice shows a peak around $B = 0.7$ T, which is caused by the formation of localized orbits encircling a single antidot.[13] There are no further low-field peaks corresponding to the formation of the orbits including more than one antidot (e.g., four antidots) because our samples have antidots with a large diameter compared with their period. On the other hand, the low-field magnetoresistance of the hexagonal lattice is quite different from that of the square lattice. At $B = 0$, magnetoresistance ρ_{xx} is about four times as large as that of the square lattice, and drops monotonically as the magnetic field increases for $B < 1.4$ T. The increase at $B = 2.2$ T is a Shubnikov–de Haas (SdH) oscillation peak, which disappears as the temperature increases. The monotonic change in the low-field magnetoresistance for the hexagonal lattice also contrasts with those of the previously reported hexagonal antidots[18] in which several magnetoresistance peaks appeared but no increase at $B = 0$ was observed. This is because, as will be discussed later, the electron motion at $B = 0$ is effectively stagnated by the large-diameter antidots with a hexagonal lattice configuration.

We evaluate the electron density (n_s) and the antidot diameter (d) of the hexagonal antidot lattice. We estimate n_s by assuming that at the SdH peak of $B = 2.2$ T the Fermi level coincides with the center of the second-lowest Landau levels. We obtain $n_s = 1.6 \times 10^{11}$ cm^{-2}, which is about 60% of that of the starting material. To estimate the antidot diameter d, which is usually different from the structural diameter of the etched holes, we assume that the magnetoresistance falls to zero at a field where $2l_{\text{cycl}} = a - d$ is satisfied.[13] Here, $l_{\text{cycl}} = \hbar k_F/eB$ is the classical cyclotron radius and $k_F = \sqrt{2\pi n_s}$ is the Fermi wave number. From this assumption, we obtain $d \sim 150$ nm from the field of $B = 2.9$ T. The obtained diameter is larger than the etched-hole diameter possibly due to the lateral depletion around the etched holes.

The most prominent feature in the trace of the hexagonal antidot is the appearance of large-amplitude oscillations at $B \lesssim 0.3$ T. To show the oscillations more clearly, two sets of low-field magnetoresistance traces for the hexagonal lattice are shown in Fig. 2(a) and (b) for the temperature ranges 1.5 K $< T <$ 14 K and 0.2 K $< T <$ 1 K, respectively. It is clear from these traces that they are periodic in B up to 0.3 T. Even though two sets of traces are obtained from two different cool-down procedures, the oscillation periods have the same value of 60 mT in both cases, almost equal to $h/2eA$, where A is the area of the unit cell of the lattice. Furthermore, the maximum of the oscillations always shows up at $B = 0$, irrespective of different temperatures and different cool-down procedures. The amplitude of these oscillations largely depends on the temperature. The inset in Fig. 2(a) shows the temperature dependence of the zero-field resistance $\rho_{xx}(B = 0)$ and the conductivity amplitude $\Delta\rho_{xx}/\rho_{xx}^2$. Although there is a gap especially for $\rho_{xx}(B = 0)$ at $T = 1$ K due to the different cool-

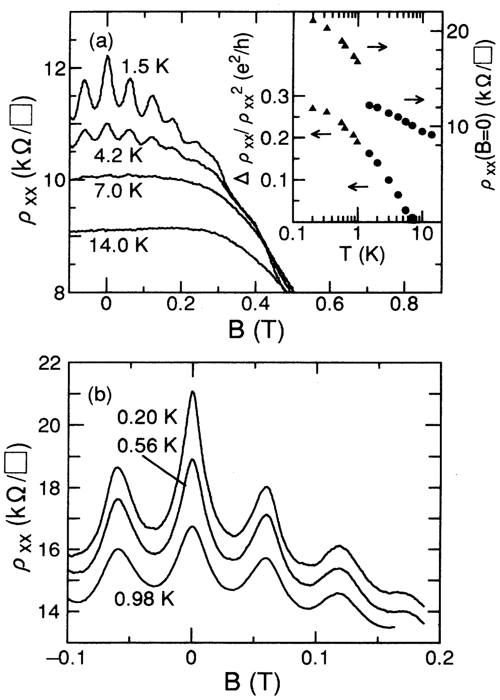

FIG. 2. (a) Magnetoresistance traces in a hexagonal lattice of antidots in a temperature range from 1.5 K to 14 K in the dark. The inset shows the temperature dependence of zero-field magnetoresistance $\rho_{xx}(B = 0)$ and the magnitude of the first oscillation $\Delta\rho_{xx}/\rho_{xx}^2(B = 0)$. Here, $\Delta\rho_{xx}$ is $\rho_{xx}(B = 0)$ minus ρ_{xx} at the first dip. The circles and triangles correspond to the measurements of (a) and (b), respectively. (b) Magnetoresistance traces in the range of 0.2 K $\leq T < 1$ K. The electron density is a little different from that of (a). Note that no traces in (a) or in (b) are offset.

down procedure, a logarithmic increase with decreasing temperature can be seen.

The above experimental results suggest the oscillations have almost the same origin as the $h/2e$ oscillations in the diffusive regime. The period of the oscillations is determined from one-half of the magnetic flux quantum for normal metal conductors threading through the mean area of the circular paths, which may be almost the same as the area of the unit cell of the lattices. The temperature dependence is similar to the previously reported results that show the logarithmic dependence on temperature. The most peculiar feature in our data is that the oscillation amplitude is very large, reaching 20% of the zero-field resistance, more than two orders of magnitude larger than that observed in metallic systems. This can also be understood in the framework of the AAS theory because even in our sample the conductance amplitude of the oscillations is of the order of e^2/h as shown in the inset in Fig. 2(a). This is consistent with the theoretically predicted value which is independent of the sample material in the diffusive regime. This originally comes from the fact that the conductivity of the antidot sample is much lower than that of metallic samples. For this reason, the relative amplitude of the oscillations becomes large.

It is worth noting that the amplitude of the oscillations depends on the lattice configuration. The oscillations are visible in the hexagonal lattice, but are almost invisible in the square lattice even though the square lattice antidots have the same period (a) and almost the same diameter (d) as the hexagonal antidots. This great difference can be explained by the lattice configuration dependence of the macroscopic-scale diffusion length of ballistic electrons. When the antidot diameter normalized by the antidot period d/a is much smaller than unity, electrons can move straight over a long distance at $B = 0$. However, when d/a becomes large, electrons can no longer go straight and thus the electron motion is stagnated in a scale much larger than the antidot period.[19] This macroscopic-scale diffusion depends largely on the lattice configuration. We have calculated the electron motions in both antidot lattices based on the model[20] using smooth antidot potentials. Figure 3(a) shows an electron trajectory in a square lattice potential. The antidot diameter is set to one-half of the antidot period. The trajectory shows electrons tend to move straight in the square lattice. From calculations for the hexagonal lattice, we have obtained a result that contrasts with the square lattice case. Affected by the hexagonal potential, electrons cannot go straight and have a high possibility of returning to their starting point with a certain circular orbit, as can be seen in Fig. 3(b). Since the interference of this kind of circular orbit is considered as the main origin of the $h/2e$ oscillations, the calculation is consistent with our experimental result that the oscillations can definitely be observed in the hexagonal lattice. We mentioned before that the zero-field resistance of the hexagonal lattice is much larger than that of the square lattice (Fig. 1). One of the reasons for this is that the macroscopic-scale diffusion length of electrons in the hexagonal potential is smaller than that in the square potential. Another reason, especially for the increase of the zero-field resistance and the emergence of background negative magnetoresistance at low temperatures for the hexagonal lattice, is the coherent backscattering of all chaotic trajectories as well as circular trajectories in Fig. 3(b).[10]

Finally, we point out that the oscillations cannot be described by the AAS theory for the diffusive regime. We have evaluated the phase coherence length by applying the calculation for $h/2e$ oscillations in metal networks to those of our hexagonal antidot lattice. This calcula-

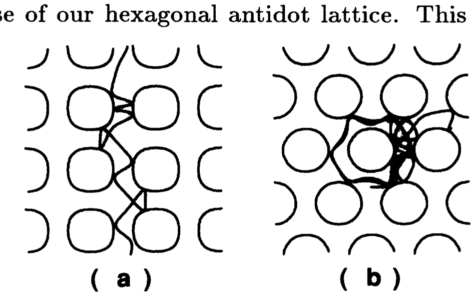

FIG. 3. Calculated electron trajectories in antidot potential at $B = 0$. The antidots are indicated by the contour lines where the potential energy is equal to the Fermi energy. (a) A trajectory passing through a square lattice. (b) A trajectory circulating around an antidot in a hexagonal lattice.

tion has been performed by Doucot and Rammal[21,22] for honeycomb networks of channels with a finite width. For details, see Refs. 21 and 22. Although the calculation has been done for the diffusive regime where the mean free path is much smaller than the channel width, we have applied it to our data ignoring the contradiction between the fitted parameters and the assumptions of the calculation. In the fitting procedure we used three parameters: the phase coherence length, the mean free path, and the channel width. From the magnetoresistance at $T = 1.5$ K in Fig. 2(a), we obtained a good fitting of the experimental curve with the phase coherence length of 610 nm, the mean free path of 330 nm, and the channel width of 20 nm. The mean free path is larger than the channel width, which is inconsistent with the assumption of the calculation. This result shows our $h/2e$ oscillations are in the ballistic regime rather than in the diffusive regime and a calculation for coherent backscattering in the ballistic regime is required.

In conclusion, we have observed large-amplitude $h/2e$ oscillations in semiconductor antidot lattices. The amplitude is two orders of magnitude larger than those of metal microstructures. The $h/2e$ oscillations are definitely visible in the magnetoresistance of hexagonal lattices but are almost invisible in that of the lattices. We expect that electrons return to their starting points and contribute to coherent backscattering with the help of the hexagonal lattice potential. We have also shown that the $h/2e$ oscillations in the antidot lattices cannot be described by the AAS theory for the diffusive regime.

The authors would like to thank A. Okamoto and K. Kasahara for preparing the heterostructure material. The authors would also like to thank S. Ishizaka and T. Sakamoto for their valuable discussions, and R. Lang and J. Sone for their encouragement throughout this work.

---

[1] B. L. Al'tshuler, A. G. Aronov, and B. Z. Spivak, Pis'ma Zh. Eksp. Teor. Fiz. **33**, 101 (1981) [JETP Lett. **33**, 94 (1981)].
[2] D. Y. Sharvin and Y. V. Sharvin, Pis'ma Zh. Eksp. Teor. Fiz. **34**, 285 (1981) [JETP Lett. **34**, 272 (1981)].
[3] B. L. Al'tshuler *et al.*, Pis'ma Zh. Eksp. Teor. Fiz. **35**, 476 (1982) [JETP Lett. **35**, 588 (1982)].
[4] F. R. Ladan and J. Maurer, C. R. Acad. Sci. (Paris) Ser. II **297**, 227 (1983).
[5] M. Gijs, C. V. Haesendonck, and Y. Bruynseraede, Phys. Rev. Lett. **52**, 2069 (1984).
[6] B. Pannetier, J. Chaussy, R. Rammal, and P. Gandit, Phys. Rev. Lett. **53**, 718 (1984).
[7] B. Pannetier, J. Chaussy, R. Rammal, and P. Gandit, Phys. Rev. B **31**, 3209 (1985).
[8] G. J. Dolan, J. C. Licini, and D. J. Bishop, Phys. Rev. Lett. **56**, 1493 (1986).
[9] C. M. Marcus *et al.*, Phys. Rev. Lett. **69**, 506 (1992).
[10] H. U. Baranger, R. A. Jalabert, and A. D. Stone, Phys. Rev. Lett. **70**, 3876 (1993).
[11] K. Ensslin and P. M. Petroff, Phys. Rev. B **41**, 12 307 (1990).
[12] M. L. Roukes and A. Scherer, Bull. Am. Phys. Soc. **34**, 633 (1989).
[13] D. Weiss *et al.*, Phys. Rev. Lett. **66**, 2790 (1991).
[14] F. Nihey and K. Nakamura, Physica B **184**, 398 (1993).
[15] D. Weiss *et al.*, Phys. Rev. Lett. **70**, 4118 (1993).
[16] G. M. Gusev *et al.*, Pis'ma Zh. Eksp. Teor. Fiz. **55**, 129 (1992) [JETP Lett. **55**, 123 (1992)].
[17] C. W. J. Beenakker and H. van Houten, in *Solid State Physics*, edited by H. Ehrenreich and D. Turnbull (Academic Press, San Diego, 1991), Vol. 44, p. 1.
[18] J. Takahara *et al.*, Jpn. J. Appl. Phys. **31**, 3786 (1992).
[19] J. Machta and R. Zwanzig, Phys. Rev. Lett. **50**, 1959 (1983).
[20] R. Fleischmann, T. Geisel, and R. Ketzmerick, Phys. Rev. Lett. **68**, 1367 (1992).
[21] B. Doucot and R. Rammal, Phys. Rev. Lett. **55**, 1148 (1985).
[22] B. Doucot and R. Rammal, J. Phys. (Paris) **47**, 973 (1986).