Effects of Magnetic Field and Flux on Perfect Channel in Metallic Carbon Nanotubes

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Numerical calculations of the conductance are performed to study effects of a weak magnetic field or flux on a perfectly conducting channel in metallic carbon nanotubes in the presence of scatterers with potential range larger than the lattice constant. The perfect channel can easily be destroyed even by a very small magnetic field or flux when there are several bands at the energy, while it is much more robust in the energy range of metallic linear bands.

Keywords: graphite, carbon nanotube, impurity, effective-mass theory, localization, magnetic field, magnetic flux

§1. Introduction

Carbon nanotubes (CN’s) are quasi-one-dimensional materials made of $sp^2$-hybridized carbon networks and have been a subject of an extensive study. In particular, the electronic structure of a single CN has been studied theoretically, which predicted that CN becomes either metallic or semiconducting depending on its chiral vector, i.e., boundary conditions in the circumference direction. In this paper we shall investigate effects of a weak magnetic field or flux on a perfect conducting channel in metallic carbon nanotubes in the presence of scatterers with potential range larger than the lattice constant.

Transport properties are particularly interesting. For impurity scattering, it was shown that there is no backscattering for potentials with a range larger than the lattice constant in metallic CN’s. The absence of backscattering was related to the presence of a topological singularity giving rise to Berry’s phase due to a rotation around the origin in the wave-vector space. It has been confirmed by numerical calculations in a tight-binding model. There have been some reports on experiments which support this theoretical prediction.

When the Fermi level moves away from the energy range where only linear bands are present, interband scattering appears because of the presence of several bands at the Fermi level. Even in such a case, a perfectly conducting channel transmitting through the system without being scattered back has been shown to exist and the conductance is always larger than $2e^2/\pi h$. It has been shown that the perfect channel is sensitive to phase breaking scattering. The Fermi energy is known to be controlled by a gate voltage for a wide range.

In understanding transport properties of nanotubes, a $k\cdot p$ method or an effective-mass approximation has been quite useful. It has been used successfully in the study of wide varieties of electronic properties of CN’s. Some of such examples are magnetic properties including the Aharonov-Bohm effect on the band gap, optical absorption spectra, exciton effects, interaction effects on band gaps, lattice instabilities in the absence and presence of a magnetic field, magnetic properties of ensembles of nanotubes, effects of spin-orbit interaction, and electronic properties of nanotube caps. Long wavelength phonons and electron-phonon scattering have also been studied.

In this paper, we perform numerical calculations of the conductance of a metallic nanotube as a function of the length in the presence of a weak magnetic field or flux at different energies. We shall consider the case that a nanotube contains scatterers with potential range larger than or comparable to the lattice constant and several traveling modes are present at the Fermi energy. We shall show that the perfect channel can easily be destroyed even by a very small magnetic field or flux when there are several bands at the energy, while it is much more robust in the energy range of metallic linear bands.

The paper is organized as follows: In §2 the treatment of electronic states in a $k\cdot p$ scheme is reviewed briefly. Numerical results on the length dependence of the conductance are presented in §3. Effects of a nonzero curvature and lattice distortion are discussed in relation to a magnetic flux in §4, which contains also a short summary and conclusion.

§2. Formulation of Problem

2.1 Effective-Mass Scheme

A graphite sheet is a zero-gap semiconductor in the sense that the conduction and valence bands consisting of $\pi$ states cross at K and K’ points of the Brillouin zone. Electronic states near a K point of 2D graphite are described by the $k\cdot p$ equation:

$$\gamma \langle \sigma \cdot \hat{k} \rangle F(r) = \varepsilon F(r),$$

(2.1)

where $\gamma$ is the band parameter, $\sigma = (\sigma_x, \sigma_y)$ is the Pauli spin matrix, $\hat{k} = (\hat{k}_x, \hat{k}_y)$ is a wave vector operator, and $F(r)$ is the envelope function with two components representing the amplitude at two carbon atoms in a unit cell.

In a metallic nanotube, states can be obtained by imposing periodic boundary conditions in the circumference direction, i.e.,

$$F(r + L) = F(r) \exp(2\pi i \varphi).$$

(2.2)

where $L$ is the lattice translation vector called the chiral
vector and
\[ \varphi = \frac{\phi}{\phi_0}, \]  
(2.3)
with \( \phi_0 = ch/e \) being the magnetic flux quantum and \( \phi \) an Aharonov-Bohm flux passing through the cross section. In the following we shall choose the \( x \) axis in the circumference direction and the \( y \) axis in the axis direction, i.e., \( \mathbf{L} = (L, 0) \), where \( L \) is the circumference.

In the presence of a magnetic field \( H \) perpendicular to the \( x \) axis, the wave vector operator is written as \( \hat{k} = -i \nabla + (e/\hbar) \mathbf{A} \), where \( \mathbf{A} \) is the vector potential given by
\[ \mathbf{A} = \left(0, \frac{LH}{2\pi}, \sin \frac{2\pi x}{L}\right). \]  
(2.4)

The wave function is written as
\[ \mathbf{F}(r) = \mathbf{F}_k(x) \frac{1}{\sqrt{A}} \exp(iky), \]  
(2.5)
where \( A \) is the length of the nanotube. The energy levels and corresponding wave functions vary sensitively as a function of the parameter \( (L/2\pi)^2 \) proportional to the field strength, where \( L \) is the magnetic length defined by \( l = \sqrt{\hbar/cH} \). In the weak-field regime \( ((L/2\pi)^2 \ll 1) \) the magnetic field can be treated as a small perturbation. In the high-field regime \( ((L/2\pi)^2 \gg 1) \), on the other hand, Landau levels are formed near the top and bottom side of the cylinder surface. It is convenient for a numerical calculation to expand the wave function as
\[ \mathbf{F}_k(x) = \frac{1}{\sqrt{L}} \sum_{n=-n_c}^{n_c} G_n(k) \exp \left[ \frac{2\pi i x}{L} (n + \varphi) \right], \]  
(2.6)
with coefficients \( G_n \) and a cutoff \( n_c \). These coefficients can be calculated for the wave functions of a given energy in a manner discussed previously.\(^{42}\)

The \( k-p \) equation for the \( K' \) point is obtained by the replacement of \( \hat{k}_y \) by \( -\hat{k}_y \) in eq. (2.1). In the present case, the \( K \) and \( K' \) points are not mixed to each other due to scattering and therefore remain completely independent. This means that the consideration only of the \( K \) point is sufficient and the actual conductance or conductivity should be obtained by multiplication of factor four corresponding to the electron spin and the presence of \( K \) and \( K' \) points.

2.2 Impurity Scattering

In order to see the dependence of the conductance on the length of CN explicitly, we shall consider a model system described by scatterers with potential
\[ V(r) = \sum_i u_i \delta(r-r_i), \]  
(2.7)
where \( r_i \) denotes the position of the \( i \)th scatterer and \( u_i \) its strength. The position \( r_i \) distributes randomly and equal amount of attractive and repulsive scatterers are assumed, i.e., \( u_i = \pm u \), for the purpose of suppressing an average energy shift. In the lowest Born approximation the scattering strength is characterized by the dimensionless quantity
\[ W = \frac{n_i (|u|^2)}{4\pi \gamma^2} = \frac{n_i u^2}{4\pi \gamma^2}, \]  
(2.8)
where \( n_i \) is the concentration of impurities in a unit area and \( (\cdots) \) means the average over scatterers.\(^{43}\) The effective strength of the potential is characterized by the dimensionless parameter \( \bar{u} = u/2\gamma L \). These delta-function impurities can describe scatterers with more general potential as long as their potential range is much smaller than the circumference.

It is straightforward to calculate a scattering matrix for an impurity given by eq. (2.7) and a conductance of a finite-length nanotube containing many impurities, combining scattering matrices as discussed previously.\(^{12,42}\) The results in the absence of a magnetic field show that the conductance becomes ideal, i.e., \( G = 2e^2/\pi \hbar \) independent of the length when the Fermi level lies in linear bands because of the absence of backward scattering. Further, the conductance remains larger than \( G = 2e^2/\pi \hbar \) and approaches \( G = 2e^2/\pi \hbar \) with the increase in the length corresponding to the presence of a perfectly conducting channel transmitting through the system without being scattered back.

When a conventional Boltzmann transport equation is used, the conductivity is obtained as
\[ \sigma = \int d\varepsilon \left( -\frac{\partial f}{\partial \varepsilon} \right) \sigma(\varepsilon), \]  
(2.9)
with \( f \) being the Fermi distribution function and
\[ \sigma(\varepsilon) = \frac{e^2}{\pi \hbar} \sum_m \Lambda_m(\varepsilon), \]  
(2.10)
where \( \Lambda_m(\varepsilon) \) is the mean free path. It satisfies the transport equation
\[ \sum_{m'} (K_{m-m'+} - K_{m+m'+}) \Lambda_{m'}(\varepsilon) = 1, \]  
(2.11)
where \( m \) and \( m' \) denote the bands crossing the energy \( \varepsilon \), and \( +(-) \) the wave vector corresponding to the positive (negative) velocity in the \( y \) direction. The kernel for the transport equation is given by
\[ K_{\nu\mu} = \frac{A|\langle V_{\nu}|^2\rangle|}{\hbar^2 |V_{\mu} V_{\nu}|}, \]  
(2.12)
for \( \nu \neq \mu \), where \( V_{\mu} \) is the velocity of mode \( \mu \equiv (m \pm) \). The diagonal elements are defined by
\[ K_{\mu\mu} = -\sum_{\nu \neq \mu} K_{\nu\mu}. \]  
(2.13)

As has been shown previously,\(^{19,43}\) the conductivity in the absence of a magnetic field becomes infinite when the Fermi level lies in the energy range where only metallic linear bands are present, while it becomes finite when the Fermi level lies in excited conduction bands. This Boltzmann result, in particular the finite conductivity when the Fermi level lies in higher conduction bands, is quite in contrast to the presence of a perfectly transmitting channel which should make the conductivity infinite.
The origin of such a striking discrepancy has been shown to arise because of the lack of the phase coherence in the approach based on the transport equation. In fact, the perfect channel is sensitive to the presence of phase-breaking processes such as inelastic scattering. When the Fermi level lies in the linear bands, however, the conductivity becomes infinite in both transmission and Boltzmann approaches.

The absence of backward scattering and the presence of a perfect channel are a direct consequence of the time reversal symmetry. In the following we shall consider effects of a magnetic field perpendicular to the axis and a magnetic flux, both of which destroy this symmetry, and study their effects on the transport for varying Fermi level position. For the actual calculations we shall use $\tilde{u} = 0.01$ and $n_c = 10$ for which $W$ can be regarded as a single relevant parameter characterizing the effective strength of impurity scattering. Further, we shall mainly show results for the weak scattering case of $W^{-1} = 1000$ for which results become same if the length being scaled by $W$. Effects of stronger scattering (small $W^{-1}$) will be discussed later very briefly.

2.3 Conductance and Fluctuation

We shall consider two different averages of the conductance, an arithmetic average $\langle G(A) \rangle$ and a geometric average $\exp(\ln G(A))$, as a function of length $A$. The localization length $\xi$ and the inverse localization length $\alpha = \xi^{-1}$ are defined as

$$\exp(\ln G(A)) \propto \exp(-2\alpha A) = \exp(-2A/\xi),$$

for sufficiently large $A$. We shall also calculate the conductance fluctuation $\Delta G(A) = (\langle G(A) - \langle G(A) \rangle \rangle)^{1/2}$. The number of different samples is chosen to be several hundreds, which are sufficient to suppress statistical fluctuations.

3. Numberical Results

3.1 Energy Bands

Figures 1 and 2 show some examples of the band structure of a metallic nanotube in magnetic field perpendicular to the axis and magnetic flux passing through the cross section, respectively. In the following we shall confine ourselves to the case of weak magnetic field or flux because we are interested in effects of the change in the symmetry of the Hamiltonian rather than those of change of the band structure itself. For the purpose, further, we shall consider four different energies, $\varepsilon(2\pi\gamma/L)^{-1} = 0.5, 1.5, 2.5,$ and $3.5$, corresponding to the channel number 1, 3, 5, and 7, respectively.

Figures 3 and 4 show the wave number and the group velocity of the bands at the Fermi level as a function of the magnetic field and flux, respectively. As shown in the figures, these quantities do not vary so much in the field or flux range to be considered, showing that the change in the wave function and the band structure is not so important.

3.2 Excess Conductance

In the absence of magnetic field and flux the conductance is always larger than the ideal conductance of a single channel case as has previously been shown both analytically and numerically. Figure 5 shows the excess conductance defined by $\Delta G = G - (2e^2/\pi h)$ as a function of the length of the nanotube in such cases. The excess conductance decreases exponentially with the length and the decay rate is nearly independent of the position of the Fermi level.

The upward arrows show the mean free path $\Lambda_n$, which decreases with $|n|$ and satisfies $\Lambda_{-n} = \Lambda_n$. Roughly speaking, the length dependence of the excess conductance becomes exponential when the length exceeds $2N\Lambda_0 \simeq 2N\bar{\Lambda}$, where $\bar{\Lambda}$ is the average of $\Lambda_n$ and $N$ is the number of the bands. Although explicit results are not presented, for sufficiently large $W^{-1}$ the length scaling by $W$ is valid quite accurately, but shows slight deviation with decreasing $W^{-1}$, presumably due to the importance of evanescent modes, i.e., because of virtual excitations to other bands during multiple scattering on different scatterers.

3.3 Conductance and Localization

Figure 6 shows the arithmetic and geometric averages of the conductivity as a function of the length in weak magnetic fields for $W^{-1} = 1000$. The upward arrows indicate the localization length $\xi$. For $\varepsilon(2\pi\gamma/L)^{-1} = 0.5$, i.e., when only the metallic linear bands exist at the energy, the decay rate of the conductance with the increase of the length remains small in weak magnetic fields ($L/2\pi l)^2 \lesssim 0.1$. The decrease of the conductance with the length becomes appreciable in stronger magnetic fields ($L/2\pi l)^2 \gtrsim 0.1$ for which the modification of the wave function itself starts to be important. Further, the difference between two different averages becomes apparent roughly when the length exceeds the localization length. This behavior is similar to that in usual quantum wires and is a result of the fact that the conductance follows a log-normal distribution instead of a normal distribution in the localization regime. When the energy lies in higher bands, on the other hand, localization effects set in much weaker magnetic fields.

Figure 7 shows the inverse localization length $\alpha$ as a function of the magnetic field for $W^{-1} = 1000$. The inverse localization length increases only slowly as a function of the field for $\varepsilon(2\pi\gamma/L)^{-1} = 0.5$. This dependence follows the appearance of backward scattering in the matrix element discussed previously. When the energy lies in higher bands, the inverse localization length increases for much smaller value of $(L/2\pi l)^2$ and exhibits a plateau like behavior. The crossover magnetic field becomes smaller with the increase of $\varepsilon(2\pi\gamma/L)^{-1}$ or the channel number.

In the presence of magnetic flux the conductance exhibits a length dependence quite similar to that in the case of a magnetic field as long as the flux is small $\phi/\phi_0 \lesssim 0.2$. Therefore the explicit results will not be presented and only the resulting localization length is discussed in the following. Figure 8 shows the inverse localization length as a function of the magnetic flux for $W^{-1} = 1000$. The dependence on the flux is completely different between the case that only metallic linear bands
are present and that several bands coexist. The ‘plateau’ value of the inverse localization length is essentially same as that in the presence of a magnetic field. However, the threshold magnetic flux does not show appreciable dependence on the number of channels at the energy.

The inverse localization length at the ‘plateau’ increases slightly with the channel number. This is likely to be a direct consequence of the fact that the mean free path decreases with the channel number as shown in Fig. 5. In fact the localization length \( \xi \) shows the dependence \( \xi \propto 2N\Lambda_0 \gtrsim 2N\Lambda \), where \( N \) is the channel number, \( \Lambda_0 \) is the largest mean free path among the bands at the energy, and \( \Lambda \) is the average. The absolute value of \( \xi \) is of the same order as \( 2N\Lambda_0 \).

The flux dependence of the inverse localization length for \( \varepsilon(2\pi\gamma/L) = 1.5 \), 2.5, and 3.5 is likely to be converted into a common curve if its absolute value is scaled properly. In the presence of magnetic field perpendicular to the axis, another scaling in the effective magnetic-field strength is necessary. An Aharonov-Bohm magnetic flux causes only an additional phase proportional to the flux in the wave function and therefore its effect on the symmetry breaking is essentially independent of the energy. In a perpendicular magnetic field, on the other hand, it causes mixing between the subbands and the mixing increases with the wave vector along the axis direction, i.e., with the increase of the energy. Therefore, modifications of the wave functions due to the field become stronger with the increase of the energy.

### 3.4 Conductance Fluctuations

Figure 9 shows the conductance fluctuation as a function of the length in weak magnetic fields for \( W^{-1} = 1000 \). When the energy lies in the region of metallic linear bands, the conductance is quantized and therefore does not exhibit fluctuations in the absence of a magnetic field. With the increase of a magnetic field, the fluctuation appears, i.e., first increases with the increase of the length and then starts to decrease. The maximum fluctuation appears at a length slightly smaller than the localization length. The overall length dependence is similar to that of a one-dimensional system known exactly.\(^{48}\)

When there are several channels at the energy, the fluctuation takes a sharp maximum roughly when the length is about the mean free path. It has also a broad maximum at a length slightly smaller than the localization length. The maximum fluctuation around the localization length is similar to that in the case of a single-channel case and is quite analogous to the behavior in quantum wires with several current carrying subbands.\(^{44,45}\) Further, the absolute value of the fluctuation is almost same as that in quantum wires in magnetic fields. This can be regarded as a manifestation of universal conductance fluctuations\(^{49,50}\) and shows that the symmetry of the Hamiltonian has changed from a symplectic class into unitary in the presence of a weak magnetic field.\(^{51}\)

### 3.5 Dependence on \( W \)

In the limit of weak scatterers, effects of scattering are expected to be characterized by a single dimensionless quantity \( W^{-1} \). In fact, similar calculations for \( W^{-1} = 10000 \) show that the results are completely same if the length is scaled by \( W^{-1} \). With the increase of scattering strength, higher-order multiple scattering becomes important and the results become different from those obtained by the length scaling. In fact, the numerical results shown in Fig. 10 for \( W^{-1} = 100 \) show some deviation from this expectation, which increases with the number of bands. A part of the reason lies in the fact that effective scattering strength increases with the energy for a given \( W \) in two-dimensional graphite.\(^{52,53}\)

### §4. Discussion and Conclusion

The nonzero curvature present in actual carbon nanotubes causes a shift in the origin of \( k_x \) and \( k_y \) in the \( k-p \) Hamiltonian.\(^{37,54}\) The shift in the axis \( y \) direction does not cause any effect on the conductance. However, the shift in the circumference \( x \) direction can be replaced by an effective magnetic flux passing through the cross section, i.e., \( \Delta \gamma = (2\pi/L)(\phi/\phi_0) \), although the effective flux has a different signature between the K and K’ points. This difference is irrelevant in the present case that the K and K’ points are completely independent. For the K point, for example, it was estimated as

\[
\frac{\phi}{\phi_0} = -\frac{2\pi a}{4\sqrt{3}L}\cos 3\eta, \quad (4.1)
\]

with \( p = 1 - (3/8)(\gamma/\gamma') \), \( \gamma' = -(\sqrt{3}/2)V_{pp}a \), and \( \gamma = -(\sqrt{3}/2)(V_{pp} - V_{pp}^*)a \), where \( a \) is the lattice constant of the two-dimensional graphite, \( \eta \) is the chiral angle (\( \eta = 0 \) for a zigzag nanotube and \( \eta = \pi/6 \) for an armchair nanotube), and \( V_{pp} \) and \( V_{pp}^* \) are the conventional tight-binding parameters for neighboring \( p \) orbitals.\(^{57}\)

For usual parameters, we have \( \gamma'/\gamma \sim 8/3 \) and therefore it is very difficult to make a reliable estimation of \( p \) although \( p \ll 1. \)\(^{57}\) If we assume \( p = 0.1 \), for example, the effective flux is \( \phi/\phi_0 \sim 0.005 \) for a typical nanotube with \( L \sim 5 \) nm except in the armchair case. Figure 8 shows that this small flux can give rise to a large effect on the conductance when there are several bands at the Fermi level, while it causes essentially no effect when the Fermi level lies in the region of metallic linear bands.

Several first-principles calculations were reported on the band structure of nanotubes, which seems to show that the shifts \( \Delta k_x \) and \( \Delta k_y \) are quite sensitive to details of methods. In fact, Hamada et al\(^{47}\) gave \( \Delta k_y < 0 \) for \( (n_x, n_y) = (12, 6) \) (so-called (6,6) armchair nanotube) and Mintmire et al\(^{48}\) gave \( \Delta k_y > 0 \) for \( (n_x, n_y) = (10, 5) \) (so-called (5,5) tube).

The presence of a lattice distortion also causes a shift in \( k_x \) and \( k_y \) in the \( k-p \) Hamiltonian.\(^{39,54,55}\) The shift in the \( x \) direction can again be replaced by an effective magnetic flux (its signature is opposite between the K and K’ points).\(^{56-58}\) The flux for the K point is written as

\[
\frac{\phi}{\phi_0} = \frac{Lg_2}{2\pi\gamma}[(uxx - uyy)\cos 3\eta - 2uxy\sin 3\eta], \quad (4.2)
\]

where \( u_{\mu\nu} (\mu, \nu = x, y) \) denotes the lattice strain and \( g_2 \) is the electron phonon interaction energy given by
$g_2 = -(3\kappa/4)V_{pp}^\ast$ with $\beta \sim 2$ and $\kappa \sim 1/3$. This shows that twist and stretch deformation give rise to flux in armchair and zigzag nanotubes, respectively. It does not seem so difficult to cause appreciable effect on the conductance by external forces in particular when there are several bands at the Fermi level.

When the range of scattering potential becomes smaller than the lattice constant of the 2D graphite, the effective potential for two sublattice sites in a unit cell can be different and causes backscattering due to “pseudo-spin-flip scattering.” Further, it causes intervalley scattering between K and K’ points, which causes also backscattering. As a result the resistance becomes nonzero. Effects of such short-range scattering can vary strongly as a function of the Fermi level.

In nanotubes with small radius, the presence of a trigonal warping can also play important roles. In fact, the warping was shown to cause backward scattering even in the energy range of linear bands although effects are extremely small and likely to be neglected completely. However, it may give rise to a huge effect through the change in the symmetry when the Fermi level lies in higher bands. Calculations of effects of such short-range scatterers and trigonal warping are currently under way.

In this paper numerical calculations of the conductance have been performed to study effects of a weak magnetic field or flux on a perfectly conducting channel in metallic carbon nanotubes in the presence of scatterers with potential range larger than the lattice constant. The perfect channel can easily be destroyed even by a very small magnetic field or flux when there are several bands at the energy, while it is much more robust in the energy range of metallic linear bands. This result is quite analogous to that in the presence of phase breaking scattering discussed previously.

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References


Figure Captions
Fig. 1 The band structure of a metallic nanotube in the presence of a magnetic field perpendicular to the axis. The strength of the field is characterized by the parameter $(L/2\pi l)^2$ proportional to the magnetic field.
Fig. 2 The band structure of a metallic nanotube in the presence of a magnetic field perpendicular to the axis. The strength of the field is characterized by the parameter $(L/2\pi l)^2$ proportional to the magnetic field.
Fig. 3 The number waves and the group velocities of the bands at the Fermi level as a function of the magnetic field. The Fermi energy is at $\varepsilon(2\pi\gamma/L)^{-1} = 0.5, 1.5, 2.5, and 3.5$ from the left to right.
Fig. 4 The number waves and the group velocities of the bands at the Fermi level as a function of the magnetic field.
Fig. 5 The exuce conductance $\Delta G = G - (2e^2/\pi h)$ versus the length in the absence of magnetic field and flux for $\varepsilon(2\pi\gamma/L)^{-1} = 1.5, 2.5, and 3.5$. The arrows indicate the mean free path of the bands at these energies. The bands with same $|n|$ are degenerate and have a same mean free path.
Fig. 6 The arithmetic and geometric averages of the conductivity as a function of the length in weak magnetic fields $(L/2\pi l)^2 \varepsilon 0.5$ for $W^{-1} = 1000$. The upward arrows indicate the localization length $\xi_n$. (a) $\varepsilon(2\pi\gamma/L)^{-1} = 0.5$. (b) 1.5. (c) 2.5. (d) 3.5.
Fig. 7 The inverse localization length as a function of the magnetic field for $W^{-1} = 1000$ in a logarithmic scale. When several bands coexist at the Fermi level, the field-dependence can be expressed approximately by a single common curve by shifting each curve in both directions.
Fig. 8 The inverse localization length as a function of the magnetic field for $W^{-1} = 1000$ in a logarithmic scale. When several bands coexist at the Fermi level, the field-dependence can be expressed approximately by a single common curve by shifting each curve in the vertical direction.
Fig. 9 The conductance fluctuation as a function of the length in weak magnetic fields. The downward arrows indicate the localization length $\xi_n$. (a) $\varepsilon(2\pi\gamma/L)^{-1} = 0.5$. (b) 1.5. (c) 2.5. (d) 3.5.
Fig. 10 The inverse localization length as a function of the magnetic field for $W^{-1} = 100$ in a logarithmic scale. The deviation from the results in the weak-scattering case increases with the Fermi level.
Fig. 1

Fig. 2
Wave Number (units of $2\pi/L$)

Magnetic Field: $(L/2\pi l)^2$

Magnetic Flux (units of $\phi_0$)

Fig. 3

Fig. 4
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Fig. 5

Fig. 6(a)

Fig. 6(b)

Fig. 6(c)
Fig. 6(d)

Fig. 7

Fig. 8

Fig. 9(a)
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Fig. 9(b) and Fig. 9(c)

Fig. 9(d)

Fig. 10