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Single-barrier problem and Anderson localization in a one-dimensional interacting electron system

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Transport through a single barrier in a one-dimensional (1D) interacting electron system is studied theoretically. By using renormalization group and duality mapping, the phase diagram of the ground state is shown to be divided into four regions in terms of the zero-temperature limits of the charge and spin conductances. The conductances are calculated perturbatively for both limits of weak and strong potential. The results are applied to clarify the crossover and scaling of the Anderson localization in a 1D system with dilute impurities. It is shown that the temperature dependence of the resistivity of the system can change significantly around a characteristic temperature corresponding to discretization energy.

I. INTRODUCTION

Effects of electron correlation on quantum transport have attracted great attention.[1] In particular, in a one-dimensional (1D) electron system both the interaction and randomness play important roles, i.e., the former makes the system a Luttinger liquid[2] while the latter causes the localization.[3] The interplay of these two has been discussed by several authors, and a metal-insulator transition is predicted as the interaction changes.[4–6] Although the Luttinger liquid has been studied only from an academic point of view, it may actually be realized in very narrow single-channel quantum wires, which will be widely fabricated in laboratories in the near future. Another example of real Luttinger-liquid systems is an edge state in the two-dimensional fractional quantum Hall system.[7,8] If there are some impurities or roughness, electronic transport in quantum wires or edge states will be crucially affected by the presence of these defects. Thus, not only in the academic sense but also from the practical point of view, it will be quite interesting and important to investigate the electronic transport in Luttinger liquids with one or few impurities.

Kane and Fisher have very recently discussed the electronic transport through a *single* barrier in a Luttinger liquid,[9] inspired by recent studies on the electron transport through very narrow mesoscopic quantum wires which can be regarded as one dimensional.[10–12] They derived an effective Lagrangian for the phase field at the barrier site by integrating out the continuum degrees of freedom. Hence the problem becomes zero dimensional and is equivalent to that of a quantum mechanical particle moving in a periodic potential subject to the dissipation of Caldeira-Leggett type.[13–16] They showed that the system is classified into two phases: insulating phase and perfectly conducting phase. A similar conclusion has also been drawn by Glazman, Ruzin, and Shklovskii,[17] who have studied the tunneling of the Wigner crystal through a pinning potential barrier.

In the first part of this paper we generalize the model proposed by Kane and Fisher to include the spin degrees of freedom in order to make the model more realistic. There are several aspects coming out. The phase diagram is divided into four regions in the plane of coupling constants by the behavior of the charge and spin conductances at low temperatures. Phase boundaries depend on the strength of the potential in contrast to the spinless case. These zero-dimensional (0D) results are then applied to the localization problem. We propose the characteristic temperature $T_{\rm dis} = v/k_B R$ (R: average interval between neighboring impurities) in addition to $T_{\rm loc} = v/k_B L_{\rm loc}$ in Ref. 6 ($L_{\rm loc}$: localization length). Above $T_{\rm dis}$ the impurity potential acts as the assembly of independent barriers and the 0D results are applicable. Below $T_{\rm dis}$ the recursion formulas for renormalization-group (RG) flow change to those discussed in Ref. 6. This crossover, which manifests itself in the nonmonotonous temperature dependence of the resistivity, is described in a unified fashion. The issue of the resonant tunneling through a double-barrier structure in a Luttinger liquid is not discussed in this paper but will be reported elsewhere.[18] We set $\hbar = k_B = 1$ in this paper.

II. MODEL

We analyze the spin-dependent Tomonaga-Luttinger model[19,20] with a scattering potential at $x = 0$, in which only forward scatterings are included as electron-electron interactions. The partition function of the system at temperature T can be written in terms of phase fields, $\theta(x,\tau)$ and $\phi(x,\tau)$, as

$$Z = \int \mathcal{D}\theta \int \mathcal{D}\phi \exp\left(- \int_0^\beta d\tau [L_0(\tau) + L_1(\tau)] \right), \quad (2.1)$$

where $\beta = 1/T$, $\theta(x,\beta) = \theta(x,0)$, and $\phi(x,\beta) = \phi(x,0)$. $L_0(\tau)$ is the (imaginary-time) Lagrangian of a pure system *with no impurity*,

$$L_0 = \frac{1}{4\pi} \int dx \left\{ \frac{1}{v_\rho \eta_\rho} (\partial_\tau \theta)^2 + \frac{v_\rho}{\eta_\rho} (\partial_x \theta)^2 + \frac{1}{v_\sigma \eta_\sigma} (\partial_\tau \phi)^2 + \frac{v_\sigma}{\eta_\sigma} (\partial_x \phi)^2 \right\}, \quad (2.2)$$

and $L_1(\tau)$ represents the impurity (barrier) potential and is given by

$$L_1 = -V_0 \sum_s \left[\Psi_{1s}^\dagger(0,\tau)\Psi_{2s}(0,\tau) + \Psi_{2s}^\dagger(0,\tau)\Psi_{1s}(0,\tau)\right] = -\frac{2V_0}{\pi\alpha}\cos\theta(0,\tau)\cos\phi(0,\tau), \tag{2.3}$$

where V_0 is the strength of the scattering potential, $\Psi_{1(2)s}$ is the field operator for an electron with velocity v_F ($-v_F$) and spin s, and α is a cutoff parameter of the order of the lattice spacing. We assume $V_0 > 0$, but actually the results do not depend on the sign of V_0. Interactions between electrons are parametrized by coupling constants, $g_{2\parallel}$ and $g_{2\perp}$, or $\tilde{g}_{2\parallel} = g_{2\parallel}/\pi v_F$ and $\tilde{g}_{2\perp} = g_{2\perp}/\pi v_F$, where v_F is the Fermi velocity. The phase fields θ and ϕ represent charge and spin density fluctuations, respectively, and these fluctuations with wave number k cost excitation energy $\epsilon = v_\rho k$ and $v_\sigma k$, where v_ρ and v_σ are defined by

$$v_\rho = v_F\left\{1 - \left(\frac{\tilde{g}_{2\parallel} + \tilde{g}_{2\perp}}{2}\right)^2\right\}^{1/2}, \quad v_\sigma = v_F\left\{1 - \left(\frac{\tilde{g}_{2\parallel} - \tilde{g}_{2\perp}}{2}\right)^2\right\}^{1/2}. \tag{2.4}$$

The other coefficients in Eq. (2.2) are given by

$$\eta_\rho = \left(\frac{1 - (\tilde{g}_{2\parallel} + \tilde{g}_{2\perp})/2}{1 + (\tilde{g}_{2\parallel} + \tilde{g}_{2\perp})/2}\right)^{1/2}, \quad \eta_\sigma = \left(\frac{1 - (\tilde{g}_{2\parallel} - \tilde{g}_{2\perp})/2}{1 + (\tilde{g}_{2\parallel} - \tilde{g}_{2\perp})/2}\right)^{1/2}, \tag{2.5}$$

which determine the exponents of various correlation functions at zero temperature.[2] Note that η_σ is fixed to be 1 when the system is invariant with respect to the SU(2) rotation in the spin space.

We integrate out the phase fields except $\theta_0 = \theta(x=0)$ and $\phi_0 = \phi(x=0)$. Introducing auxiliary fields $\lambda_1(\tau)$ and $\lambda_2(\tau)$, we first rewrite the partition function as

$$Z = \int \mathcal{D}\theta_0 \int \mathcal{D}\phi_0 \int \mathcal{D}\lambda_1 \int \mathcal{D}\lambda_2 \int \mathcal{D}\theta \int \mathcal{D}\phi \exp\Big(- \int_0^\beta d\tau \{ L_0(\tau) + L_1(\tau) + i\lambda_1(\tau)[\theta_0(\tau) - \theta(0,\tau)] + i\lambda_2(\tau)[\phi_0(\tau) - \phi(0,\tau)]\}\Big), \tag{2.6}$$

and then integrate out $\theta(x,\tau)$ and $\phi(x,\tau)$ to obtain

$$Z = \int \mathcal{D}\theta_0 \int \mathcal{D}\phi_0 \int \mathcal{D}\lambda_1 \int \mathcal{D}\lambda_2 \exp\Big(-\frac{\pi\eta_\rho v_\rho}{\beta}\sum_{\omega_n}\lambda_1(\omega_n)\lambda_1(-\omega_n)\int_{-\infty}^{\infty}\frac{dq}{2\pi}\frac{1}{\omega_n^2 + v_\sigma^2 q^2} -\frac{\pi\eta_\sigma v_\sigma}{\beta}\sum_{\omega_n}\lambda_2(\omega_n)\lambda_2(-\omega_n)\int_{-\infty}^{\infty}\frac{dq}{2\pi}\frac{1}{\omega_n^2 + v_\rho^2 q^2} + \frac{i}{\beta}\sum_{\omega_n}[\lambda_1(\omega_n)\theta_0(-\omega_n) + \lambda_2(\omega_n)\phi_0(-\omega_n)] + \frac{2V_0}{\pi\alpha}\int_0^\beta d\tau \cos\theta_0(\tau)\cos\phi_0(\tau)\Big), \tag{2.7}$$

where $\omega_n = 2\pi n/\beta$ ($n = 0, \pm1, \pm2, \ldots$) and we have neglected an unimportant numerical factor. Here the Fourier transforms are defined as

$$\theta_0(\omega_n) = \int_0^\beta d\tau\theta_0(\tau)e^{i\omega_n\tau}, \quad \phi_0(\omega_n) = \int_0^\beta d\tau\phi_0(\tau)e^{i\omega_n\tau}, \lambda_j(\omega) = \int_0^\beta d\tau\lambda_j(\tau)e^{i\omega_n\tau} \quad (j = 1,2). \tag{2.8}$$

Integrating out λ_1 and λ_2, we finally obtain

$$Z = \int \mathcal{D}\theta_0 \int \mathcal{D}\phi_0 \exp\left(-\frac{1}{2\pi\eta_\rho\beta}\sum_{\omega_n}|\omega_n||\theta_0(\omega_n)|^2 - \frac{1}{2\pi\eta_\sigma\beta}\sum_{\omega_n}|\omega_n||\phi_0(\omega_n)|^2 + \frac{2V_0}{\pi\alpha}\int_0^\beta d\tau\cos\theta_0(\tau)\cos\phi_0(\tau)\right). \tag{2.9}$$

Note that this partition function is similar to that of a quantum Brownian particle [coordinate (θ_0, ϕ_0)] moving in the periodic cosine potential (2.3) and coupled to a dissipative environment;[13–16] in our model the low-lying charge and spin excitations cause the dissipation. Hence our 1D problem is now reduced to quantum mechanics of a particle, i.e., a 0D field theory. To avoid ultraviolet divergences, we introduce a high-frequency cutoff, $\Lambda \sim v_F/\alpha$, which may also serve as a mass of the particle $m \sim 1/\Lambda$.[16]

III. WEAK BARRIER POTENTIAL

In this section we consider the limit where the barrier potential is very weak. We thus perform RG transformations and calculate charge and spin conductances perturbatively with respect to V_0.

A. Scaling equations

Following Fisher and Zwerger,[16] we derive scaling equations for the barrier potential by recursively integrating out high-frequency modes. At zero temperature the partition function (2.9) is written as

$$Z = \int \mathcal{D}\theta_0 \int \mathcal{D}\phi_0 \exp(-S_0 - S_1), \tag{3.1}$$

where

$$S_0 = \frac{1}{2\pi\eta_\rho}\int_{-\Lambda}^{\Lambda}\frac{d\omega}{2\pi}|\omega||\theta_0(\omega)|^2 + \frac{1}{2\pi\eta_\sigma}\int_{-\Lambda}^{\Lambda}\frac{d\omega}{2\pi}|\omega||\phi_0(\omega)|^2, \tag{3.2}$$

$$S_1 = -\frac{2V_0}{\pi\alpha}\int d\tau \cos\theta_0(\tau)\cos\phi_0(\tau).$$

We first divide the phase fields into slow and fast modes,

$$\theta_0(\tau) = \theta_{0s}(\tau) + \theta_{0f}(\tau), \qquad \phi_0(\tau) = \phi_{0s}(\tau) + \phi_{0f}(\tau), \tag{3.3}$$

such that

$$\theta_0(\omega) \approx \begin{cases} \theta_{0s}(\omega) & \text{for } |\omega| \le \mu \\ \theta_{0f}(\omega) & \text{for } \mu \le |\omega| \le \Lambda, \end{cases}$$
$$\phi_0(\omega) \approx \begin{cases} \phi_{0s}(\omega) & \text{for } |\omega| \le \mu \\ \phi_{0f}(\omega) & \text{for } \mu \le |\omega| \le \Lambda. \end{cases} \tag{3.4}$$

Integrating out the fast modes θ_{0f} and ϕ_{0f}, we then get an effective action $\widetilde{S}$ for the slow modes in powers of V_0:

$$\widetilde{S} = \frac{1}{2\pi\eta_\rho}\int_{-\mu}^{\mu}\frac{d\omega}{2\pi}|\omega||\theta_{0s}(\omega)|^2 + \frac{1}{2\pi\eta_\sigma}\int_{-\mu}^{\mu}\frac{d\omega}{2\pi}|\omega||\phi_{0s}(\omega)|^2 + \langle F_1\rangle - \tfrac{1}{2}\left(\langle F_1^2\rangle - \langle F_1\rangle^2\right) + O(V_0^3). \tag{3.5}$$

Averaging over the fast modes, we get

$$\langle F_1\rangle = \frac{2V_0}{\pi\alpha}\left(\frac{\mu}{\Lambda}\right)^{\frac{1}{2}(\eta_\rho+\eta_\sigma)}\int d\tau \cos\theta_{0s}(\tau)\cos\phi_{0s}(\tau). \tag{3.6}$$

Since $G_\theta(\tau) = \langle\theta_{0f}(\tau)\theta_{0f}(0)\rangle$ and $G_\phi(\tau) = \langle\phi_{0f}(\tau)\phi_{0f}(0)\rangle$ are short ranged and fall off exponentially for $\tau \gg 1/\mu$ (Ref. 16), the second-order cumulant can be approximated as

$$\langle F_1^2\rangle - \langle F_1\rangle^2 \approx \left(\frac{2V_0}{\pi\alpha}\right)^2\left(\frac{\mu}{\Lambda}\right)^{\eta_\rho+\eta_\sigma}\int d\tau\Bigg\{a_1\cos[2\theta_{0s}(\tau)]\cos[2\phi_{0s}(\tau)] + a_2\left[1 - \frac{1}{2}\left(\frac{d\theta_{0s}}{d\tau}\right)^2 - \frac{1}{2}\left(\frac{d\phi_{0s}}{d\tau}\right)^2\right] + a_3\cos[2\theta_{0s}(\tau)] + a_4\cos[2\phi_{0s}(\tau)]\Bigg\}, \tag{3.7}$$

where a_i ($i = 1$–4) is some constant. Finally we must rescale the imaginary time as $\tau \to (\Lambda/\mu)\tau$ to complete the RG transformation. Note that it is not necessary to rescale θ_0 and ϕ_0 due to the fact that the theory has underlying symmetries, $\theta_0(\tau) \to \theta_0(\tau) + 2\pi$ and $\phi_0(\tau) \to \phi_0(\tau) + 2\pi$.[16] In addition, since the second term of the integrand in Eq. (3.7), $\int d\tau[(d\theta_{0s}/d\tau)^2 + (d\phi_{0s}/d\tau)^2]$, is irrelevant compared with $\int d\omega[|\omega||\theta_{0s}(\omega)|^2 + |\omega||\phi_{0s}(\omega)|^2]$, both η_ρ and η_σ are not renormalized. Hence the quantities left to be renormalized are the barrier potential, $V_0\cos\theta_0\cos\phi_0$, and its descendants in the second-order perturbation, i.e., $V_{2,0}\cos 2\theta_0$, $V_{0,2}\cos 2\phi_0$, and $V_{2,2}\cos 2\theta_0\cos 2\phi_0$. We note that the potentials, $V_{2,0}\cos 2\theta_0$ and $V_{0,2}\cos 2\phi_0$, can also be written in terms of the fermion field operators at $x = 0$ as $\frac{1}{2}V_{2,0}(\Psi_{2\uparrow}^\dagger\Psi_{1\uparrow}\Psi_{2\downarrow}^\dagger\Psi_{1\downarrow} + \text{H.c.})$ and $\frac{1}{2}V_{0,2}(\Psi_{2\uparrow}^\dagger\Psi_{1\uparrow}\Psi_{1\downarrow}^\dagger\Psi_{2\downarrow} + \text{H.c.})$.

From Eq. (3.6) we get

$$V_0(\mu) = V_0(\Lambda)\left(\frac{\mu}{\Lambda}\right)^{\frac{1}{2}(\eta_\rho+\eta_\sigma)-1}, \tag{3.8}$$

or in differential form

$$\frac{dV_0}{dl} = [1 - \tfrac{1}{2}(\eta_\rho + \eta_\sigma)]V_0(l) + O(V_0^3), \tag{3.9}$$

where $dl = -d\mu/\mu$. Thus, if $\eta_\rho + \eta_\sigma > 2$ the potential scales to zero whereas for $\eta_\rho + \eta_\sigma < 2$ it grows as the cutoff μ is reduced.

The scaling equations for $V_{2,0}$ and $V_{0,2}$ can be derived in a similar way, and the lowest-order RG equations are

$$\frac{dV_{2,0}}{dl} = (1 - 2\eta_\rho)V_{2,0}(l), \tag{3.10a}$$

$$\frac{dV_{0,2}}{dl} = (1 - 2\eta_\sigma)V_{0,2}(l), \tag{3.10b}$$

which show that $V_{2,0}\cos 2\theta_0$ ($V_{0,2}\cos 2\phi_0$) is relevant when $\eta_\rho < 1/2$ ($\eta_\sigma < 1/2$). These three RG equations, (3.9), (3.10a), and (3.10b), suffice for determining the phase diagram at zero temperature. Other higher-order terms, $V_{m,n}\cos m\theta_0\cos n\phi_0$ ($m + n \ge 4$), generated by higher-order expansions are not important, since at least one of the above three pinning potentials is always relevant in parameter regions in which the higher-order terms become relevant, $m^2\eta_\rho + n^2\eta_\sigma < 2$.

From the RG equations we can deduce the phase diagram at $T = 0$ as shown in Fig. 1 where the phase boundaries are $\eta_\rho + \eta_\sigma = 2$, $\eta_\rho = 1/2$, and $\eta_\sigma = 1/2$. In region I, $V_0 \cos\theta_0 \cos\phi_0$ is relevant so that both θ_0 and ϕ_0 are pinned around the potential minima, which means that electrons are perfectly reflected by the barrier at $T = 0\,\mathrm{K}$. In region II, only $V_{0,2} \cos 2\phi_0$ is a relevant perturbation, and therefore spin phase, ϕ_0, is pinned whereas charge phase, θ_0, is not pinned. In region III, on the other hand, $V_{2,0} \cos 2\theta_0$ is relevant; the charge phase is pinned, while the spin phase is not pinned. The physical implications of these phenomena will be discussed in Sec. IV. Finally in region IV all the pinning potentials are irrelevant, so electrons can freely go through the barrier. It is interesting to note that the phase boundary $\eta_\rho + \eta_\sigma = 2$ obtained above is different from that of the Anderson localization transition studied before in the weak-pinning limit.[5,6] This difference will be discussed in detail in Sec. V.

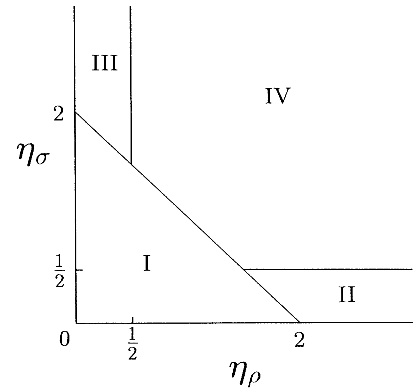

FIG. 1. The phase diagram of the ground state in the $\eta_\rho - \eta_\sigma$ plane for the weak potential limit. The phase boundaries are $\eta_\rho + \eta_\sigma = 2$, $\eta_\rho = 1/2$, and $\eta_\sigma = 1/2$.

B. Conductance

Next we shall calculate the charge (spin) conductance G_ρ (G_σ) in powers of V_0 by using the influence-functional formalism.[21] Since the method is described in detail in Ref. 16, where the mobility of a quantum Brownian particle is calculated, we simply apply their results to our problem. We refer the reader to Ref. 16 for details.

When a voltage V is applied across the potential barrier, an additional term, $eV\theta_0/\pi$, should be added to L_1 in Eq. (2.3). The (charge) current J_ρ induced by the voltage difference is given by $J_\rho = -(e/\pi)(d\theta_0/dt)$ where t is a real time, and the charge conductance is defined by $G_\rho \equiv J_\rho/V$ with $V \to 0$. On the other hand, when there is a magnetic field difference H between the two sides of the potential barrier, another term, $\mu_B H(\phi_0/2\pi)$ (μ_B: Bohr magneton), must be included in L_1, resulting in a spin current $J_\sigma = (1/2\pi)(d\phi_0/dt)$; the spin conductance is defined as $G_\sigma \equiv J_\sigma/H$ with $H \to 0$. In the absence of the potential $V_0 \cos\theta_0 \cos\phi_0$, G_ρ (G_σ) is $e^2\eta_\rho/\pi$ ($\mu_B\eta_\sigma/2\pi$).

First we evaluate the charge conductance in powers of V_0, assuming $V \neq 0$ but $H = 0$. Following the same path as in Ref. 16, we arrive at the following expression of the charge current J_ρ:

$$
\begin{aligned}
J_\rho = \frac{e^2}{\pi}\eta_\rho V - e\eta_\rho \lim_{t\to\infty} \frac{1}{t} \sum_{n>0}^{\infty} \sum_{n'>0}^{\infty} \sum_{\{e_{1j},s_{1j}\}} \sum_{\{e_{2j},s_{2j}\}} \left(\frac{iV_0}{2\pi\alpha}\right)^n \left(-\frac{iV_0}{2\pi\alpha}\right)^{n'} \\
\times \int_0^t dt_1 \int_0^{t_1} dt_2 \cdots \int_0^{t_{n-1}} dt_n \int_0^t dt'_1 \int_0^{t'_1} dt'_2 \cdots \int_0^{t'_{n'-1}} dt'_{n'} \\
\times \tfrac{1}{2} \int_0^t dt' [\rho(t') + \rho'(t')] \\
\times \exp\!\Bigg(i \int_0^t dt_0 \Big\{ \frac{1}{\pi} eV\tilde{\theta}(t_0) - \tfrac{1}{2}\tilde{\theta}(t_0)[\rho(t_0) + \rho'(t_0)] \\
- \tfrac{1}{2}\tilde{\phi}(t_0)[\sigma(t_0) + \sigma'(t_0)] \Big\} - S_2(\tilde{\theta}, \tilde{\phi}) \Bigg),
\end{aligned}
\tag{3.11}
$$

where

$$
\rho(t') = \sum_{j=1}^{n} e_{1j}\delta(t' - t_j), \qquad \rho'(t') = \sum_{j=1}^{n'} e_{2j}\delta(t' - t'_j) \qquad (e_{ij} = \pm 1), \tag{3.12a}
$$

$$
\sigma(t') = \sum_{j=1}^{n} s_{1j}\delta(t' - t_j), \qquad \sigma'(t') = \sum_{j=1}^{n'} s_{2j}\delta(t' - t'_j) \qquad (s_{ij} = \pm 1), \tag{3.12b}
$$

$$
\tilde{\theta}(t') = \pi\eta_\rho \left(\sum_{j=1}^{n} e_{1j}\Theta(t' - t_j) - \sum_{j=1}^{n'} e_{2j}\Theta(t' - t'_j) \right), \tag{3.12c}
$$

$$\tilde{\phi}(t') = \pi\eta_\sigma \left(\sum_{j=1}^{n} s_{1j}\Theta(t'-t_j) - \sum_{j=1}^{n'} s_{2j}\Theta(t'-t'_j) \right) , \tag{3.12d}$$

$$S_2(\tilde{\theta},\tilde{\phi}) = \frac{1}{\pi\eta_\rho}\int_0^t du_1 \int_0^{u_1} du_2 \tilde{\theta}(u_1)\alpha_R(u_1-u_2)\tilde{\theta}(u_2) + \frac{1}{\pi\eta_\sigma}\int_0^t du_1 \int_0^{u_1} du_2 \tilde{\phi}(u_1)\alpha_R(u_1-u_2)\tilde{\phi}(u_2), \tag{3.12e}$$

$$\alpha_R(u) = \int_0^\infty \frac{d\omega}{\pi}\omega\cos(\omega u)\coth\left(\frac{\beta\omega}{2}\right) , \tag{3.12f}$$

and $\Theta(t')$ is the step function. The summations in Eq. (3.11) are performed under the charge and spin neutrality conditions,

$$\sum_{j=1}^{n} e_{1j} = \sum_{j=1}^{n'} e_{2j}, \qquad \sum_{j=1}^{n} s_{1j} = \sum_{j=1}^{n'} s_{2j}, \tag{3.13}$$

which imply that $n+n'$ must be even. In the lowest order, Eq. (3.11) is evaluated as

$$J_\rho = \frac{e^2}{\pi}\eta_\rho V - e\eta_\rho \left(\frac{V_0}{\pi\alpha}\right)^2 \tanh(\tfrac{1}{2}\eta_\rho\beta eV) \times \int_{-\infty}^{\infty} dt\cos(\eta_\rho eVt)\exp\left[-(\eta_\rho+\eta_\sigma)\int_0^\infty d\omega \frac{e^{-\omega/\Lambda}}{\omega}\left((1-\cos\omega t)\coth\frac{\beta\omega}{2} + i\sin\omega t\right)\right] , \tag{3.14}$$

where we have adopted an exponential cutoff, $\exp(-\omega/\Lambda)$. The charge conductance G_ρ is then obtained as

$$G_\rho = \frac{e^2\eta_\rho}{\pi}\left(1 - \frac{\sqrt{\pi}\eta_\rho}{2}\frac{\Gamma((\eta_\rho+\eta_\sigma)/2)}{\Gamma((\eta_\rho+\eta_\sigma+1)/2)}\left(\frac{V_0}{\alpha\Lambda}\right)^2\left(\frac{\pi T}{\Lambda}\right)^{\eta_\rho+\eta_\sigma-2}\right) , \tag{3.15}$$

where $\Gamma(x)$ is the Γ function. The next leading term in the expansion can also be obtained from Eq. (3.11). Here we evaluate it by a simpler method: we regard the barrier potential as $(2V_{2,0}/\pi\alpha)\cos 2\theta_0$ instead of $(2V_0/\pi\alpha)\cos\theta_0\cos\phi_0$, and calculate the charge conductance in the lowest order. By so doing, together with Eq. (3.15), we get G_ρ up to the order $(V_0/\alpha\Lambda)^4$ as

$$G_\rho = \frac{e^2\eta_\rho}{\pi}\left[1 - c_0\left(\frac{V_0}{\alpha\Lambda}\right)^2\left(\frac{\pi T}{\Lambda}\right)^{\eta_\rho+\eta_\sigma-2} - c_1\left(\frac{V_{2,0}}{\alpha\Lambda}\right)^2\left(\frac{\pi T}{\Lambda}\right)^{4\eta_\rho-2}\right] , \tag{3.16}$$

where c_0 and c_1 are dimensionless numbers which depend on η_ρ, η_σ, and the cutoff procedure. Note that $V_{2,0}$ is of the order of $V_0^2/\alpha\Lambda$.

The current-voltage characteristic of the single barrier is obtained from Eq. (3.14). At zero temperature it becomes

$$\begin{aligned} J_\rho &= \frac{e^2}{\pi}\eta_\rho V - e\eta_\rho\left(\frac{V_0}{\pi\alpha}\right)^2\int_{-\infty}^{\infty} dt\cos(\eta_\rho eVt)\exp\left[-(\eta_\rho+\eta_\sigma)\int_0^\infty d\omega\frac{e^{-\omega/\Lambda}}{\omega}(1-e^{-i\omega t})\right] \\ &\quad - e\eta_\rho\left(\frac{V_{2,0}}{\pi\alpha}\right)^2\int_{-\infty}^{\infty} dt\cos(2\eta_\rho eVt)\exp\left[-4\eta_\rho\int_0^\infty d\omega\frac{e^{-\omega/\Lambda}}{\omega}(1-e^{-i\omega t})\right] \\ &= \frac{e^2}{\pi}\eta_\rho V - \frac{\pi e\eta_\rho}{\Lambda\Gamma(\eta_\rho+\eta_\sigma)}\left(\frac{V_0}{\pi\alpha}\right)^2\left(\frac{\eta_\rho eV}{\Lambda}\right)^{\eta_\rho+\eta_\sigma-1} - \frac{\pi e\eta_\rho}{\Lambda\Gamma(4\eta_\rho)}\left(\frac{V_{2,0}}{\pi\alpha}\right)^2\left(\frac{2\eta_\rho eV}{\Lambda}\right)^{4\eta_\rho-1} , \end{aligned} \tag{3.17}$$

where we have included the contribution from $V_{2,0}\cos 2\theta_0$ and neglected unimportant exponential factors such as $\exp(-\eta_\rho eV/\Lambda)$. For temperatures $T \ll eV$ the current-voltage relation deviates from Ohm's law, as described in Eq. (3.17). If $eV \ll T \ll \Lambda$, on the other hand, the current-voltage characteristic obeys Ohm's law, $J_\rho = G_\rho V$, where G_ρ is given by Eq. (3.16).

The spin conductance G_σ can also be evaluated in a similar manner by taking $H \neq 0$ and $V = 0$. Here we show only the final expression which is valid up to the order $(V_0/\alpha\Lambda)^4$:

$$G_\sigma = \frac{\mu_B\eta_\sigma}{2\pi}\left[1 - c'_0\left(\frac{V_0}{\alpha\Lambda}\right)^2\left(\frac{\pi T}{\Lambda}\right)^{\eta_\rho+\eta_\sigma-2} - c'_1\left(\frac{V_{0,2}}{\alpha\Lambda}\right)^2\left(\frac{\pi T}{\Lambda}\right)^{4\eta_\sigma-2}\right], \tag{3.18}$$

where c'_0 and c'_1 are dimensionless numbers.

The temperature dependences of G_ρ and G_σ are naturally understood from the scaling equations (3.9), (3.10a), and (3.10b). For example, integrating the RG equation (3.9) from $\mu = \Lambda$ to $\mu = T$ yields the renormalized potential, $(2V_0/\pi\alpha)(T/\Lambda)^{\frac{1}{2}(\eta_\rho+\eta_\sigma)-1}$. Then, the reduction of the conductances due to the potential scattering is proportional to the square of the renormalized one, giving the power-law dependence of $T^{\eta_\rho+\eta_\sigma-2}$ in Eqs. (3.16) and (3.18). The same reasoning can be applied also to the reductions due to $V_{2,0}\cos 2\theta_0$ and $V_{0,2}\cos 2\phi_0$.

The above perturbative calculations are valid if the reductions of the conductances due to the potential scattering are much smaller than the zeroth-order term, $e^2\eta_\rho/\pi$ or $\mu_B\eta_\sigma/2\pi$. Thus Eqs. (3.16) and (3.17) are valid down to $T = 0\,\mathrm{K}$ in regions II and IV of Fig. 1, where $G_\rho(T = 0\,\mathrm{K}) = e^2\eta_\rho/\pi$. In the other regions (I and III), however, the expansion is valid only for high temperatures. At low temperatures expansions with respect to the tunneling matrix elements become appropriate, giving $G_\rho(T = 0\,\mathrm{K}) = 0$. Similarly, the expansion of G_σ is justified only in regions III and IV down to $T = 0\,\mathrm{K}$, and $G_\sigma(T = 0\,\mathrm{K}) = \mu_B\eta_\sigma/2\pi$. In regions I and II, on the other hand, the expansions fail at low temperatures, which suggests $G_\sigma(T = 0\,\mathrm{K}) = 0$. In summary, at low temperatures the perturbative calculations are justified in the regions where the pinning potential is irrelevant and the phase field of interest is not pinned. Finally we note that for the noninteracting case ($\eta_\rho = \eta_\sigma = 1$) the leading-order corrections proportional to $(V_0/\alpha\Lambda)^2$ in Eqs. (3.16) and (3.18) are independent of T, which is consistent with what the Landauer formula tells; the conductance can take any value from 0 to e^2/π at zero temperature.

IV. STRONG BARRIER POTENTIAL

In this section we consider the opposite limit in which the barrier potential is very strong, $V_0/\alpha\Lambda \gg 1$. In this limit the electron transport can be viewed as the tunneling from a potential minimum to an adjacent minimum, and tunneling matrix elements are natural expansion parameters. The cosine potential (2.3) has minima at $(\theta_0, \phi_0) = ((m+n)\pi, (m-n)\pi)$ and maxima at $(\theta_0, \phi_0) = ((m+n+1)\pi, (m-n)\pi)$, where m and n are integers (Fig. 2). Thus a particle initially at $(\theta_0, \phi_0) = (0, 0)$ can tunnel to $(\pm\pi, \pm\pi)$ through a lower tunnel barrier and to $(\pm 2\pi, 0)$ and $(0, \pm 2\pi)$ through a higher barrier. Physically these processes correspond to the tunneling of one electron or hole $[(\pm\pi, \pm\pi)]$, the singlet pair of two electrons or holes $[(\pm 2\pi, 0)]$, and the triplet electron-hole pair $[(0, \pm 2\pi)]$, respectively.

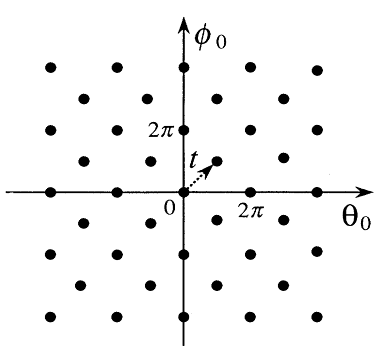

FIG. 2. Minima of the pinning potential $-(2V_0/\pi\alpha)\cos\theta_0\cos\phi_0$. The matrix element for the $(0,0) \to (\pi,\pi)$ tunneling is t.

A. Duality mapping and scaling equations

Generalizing the duality argument by Schmid[14] and using the dilute instanton gas approximation (DIGA), we show below that the partition function in the strong potential limit is mapped to that in the weak potential limit discussed in the preceding section.

Remembering that the high-frequency cutoff Λ serves as the mass m of the Brownian particle, we may write the partition function (2.9) as

$$Z = \int \mathcal{D}\theta_0 \int \mathcal{D}\phi_0 \exp(-S_0 - S_1), \tag{4.1a}$$

$$S_0 = \int_0^\beta d\tau \left\{ \frac{1}{2} m \left[\left(\frac{d\theta_0}{d\tau}\right)^2 + \left(\frac{d\phi_0}{d\tau}\right)^2 \right] + \frac{2V_0}{\pi\alpha}(1 - \cos\theta_0\cos\phi_0) \right\}, \tag{4.1b}$$

$$S_1 = \frac{1}{2\pi\eta_\rho}\sum_{\omega_n} |\omega_n||\theta_0(\omega_n)|^2 + \frac{1}{2\pi\eta_\sigma}\sum_{\omega_n} |\omega_n||\phi_0(\omega_n)|^2. \tag{4.1c}$$

Note that for simplicity we have assumed that the mass is isotropic in the (θ_0, ϕ_0) plane. We evaluate the partition function in the semiclassical limit, in which the functional integral is dominated by the stationary path of S_0+S_1. It is important to notice that S_0 describes the physics in the short time scale, i.e., tunneling of an electron (instanton), whereas S_1 describes the physics in the long time scale, i.e., interaction between instantons. We, therefore, first construct the stationary paths of S_0, denoted by $\bar{\theta}_0$ and $\bar{\phi}_0$, and then we substitute them into S_1. $\bar{\theta}_0(\tau)$ and $\bar{\phi}_0(\tau)$ are determined from

$$\frac{\delta S_0}{\delta\bar{\theta}_0} = -m\frac{d^2\bar{\theta}_0}{d\tau^2} + \frac{2V_0}{\pi\alpha}\sin\bar{\theta}_0\cos\bar{\phi}_0 = 0, \tag{4.2a}$$

$$\frac{\delta S_0}{\delta\bar{\phi}_0} = -m\frac{d^2\bar{\phi}_0}{d\tau^2} + \frac{2V_0}{\pi\alpha}\cos\bar{\theta}_0\sin\bar{\phi}_0 = 0, \tag{4.2b}$$

or equivalently,

$$\frac{d^2}{d\tau^2}(\bar{\theta}_0 + \bar{\phi}_0) = \frac{2V_0}{\pi\alpha m}\sin(\bar{\theta}_0 + \bar{\phi}_0), \tag{4.3a}$$

$$\frac{d^2}{d\tau^2}(\bar{\theta}_0 - \bar{\phi}_0) = \frac{2V_0}{\pi\alpha m}\sin(\bar{\theta}_0 - \bar{\phi}_0). \tag{4.3b}$$

A solution of $d^2X/d\tau^2 = (2V_0/\pi\alpha m)\sin X$ describing one instanton at $x = 0$ is given by

$$X(\tau) = 2\arccos\left(-\tanh[\tau(2V_0/\pi\alpha m)^{1/2}]\right), \tag{4.4}$$

which satisfies $X(-\infty) = 0$ and $X(\infty) = 2\pi$. From this we see that the width of the instanton is of the order of $(\pi\alpha m/2V_0)^{1/2}$. In the DIGA, we neglect the overlaps of instantons assuming that β is much larger than the width and that fugacity of instantons is very small. Thus we write $\bar{\theta}_0$ and $\bar{\phi}_0$ as linear combinations of the one-instanton solution $X(\tau)$:

$$\bar{\theta}_0(\tau)+\bar{\phi}_0(\tau)=\sum_{j=1}^{n_1} e_{1j}X(\tau-\tau_{1j}), \tag{4.5a}$$

$$\bar{\theta}_0(\tau)-\bar{\phi}_0(\tau)=\sum_{j=1}^{n_2} e_{2j}X(\tau-\tau_{2j}), \tag{4.5b}$$

where $e_{ij}=1$ (instanton) or -1 (anti-instanton) and τ_{ij}'s specify the locations of instantons or anti-instantons. It follows from $\bar{\theta}_0(0)=\bar{\theta}_0(\beta)$ and $\bar{\phi}_0(0)=\bar{\phi}_0(\beta)$ that $\sum_j e_{1j}=\sum_j e_{2j}=0$ (neutrality condition). We may write the Fourier transform of $\bar{\theta}_0(\tau)$ and $\bar{\phi}_0(\tau)$ as

$$\bar{\theta}_0(\omega_n)=\int_0^\beta d\tau\,\bar{\theta}_0(\tau)e^{i\omega_n\tau}=\frac{i\pi}{\omega_n}\sum_{j=1}^{n_1}e_{1j}\exp(i\omega_n\tau_{1j})+\frac{i\pi}{\omega_n}\sum_{j=1}^{n_2}e_{2j}\exp(i\omega_n\tau_{2j}), \tag{4.6a}$$

$$\bar{\phi}_0(\omega_n)=\int_0^\beta d\tau\,\bar{\phi}_0(\tau)e^{i\omega_n\tau}=\frac{i\pi}{\omega_n}\sum_{j=1}^{n_1}e_{1j}\exp(i\omega_n\tau_{1j})-\frac{i\pi}{\omega_n}\sum_{j=1}^{n_2}e_{2j}\exp(i\omega_n\tau_{2j}), \tag{4.6b}$$

where we have used an approximation,

$$\int_0^\beta d\tau\,e^{i\omega_n\tau}\frac{d}{d\tau}X(\tau-\tau_{ij})\approx e^{i\omega_n\tau_{ij}}\int_{-\infty}^{\infty}d\tau\frac{dX(\tau)}{d\tau}=2\pi e^{i\omega_n\tau_{ij}}. \tag{4.7}$$

By substituting Eqs. (4.6a) and (4.6b) into Eq. (4.1c), the partition function can be calculated, within the DIGA, as

$$\begin{aligned}Z=&\sum_{n_1=0}^{\infty}\sum_{n_2=0}^{\infty}\sum_{\{e_{1j}\}}\sum_{\{e_{2j}\}}\frac{1}{n_1!n_2!}\int_0^\beta d\tau_{11}\cdots\int_0^\beta d\tau_{1n_1}\int_0^\beta d\tau_{21}\cdots\int_0^\beta d\tau_{2n_2}\\&\times y_0^{n_1+n_2}\exp\left(-\frac{\pi}{2\eta_\rho\beta}\sum_{\omega_n}\frac{1}{|\omega_n|}\left|\sum_{j=1}^{n_1}e_{1j}\exp(i\omega_n\tau_{1j})+\sum_{j=1}^{n_2}e_{2j}\exp(i\omega_n\tau_{2j})\right|^2\right.\\&\left.\qquad-\frac{\pi}{2\eta_\sigma\beta}\sum_{\omega_n}\frac{1}{|\omega_n|}\left|\sum_{j=1}^{n_1}e_{1j}\exp(i\omega_n\tau_{1j})-\sum_{j=1}^{n_2}e_{2j}\exp(i\omega_n\tau_{2j})\right|^2\right),\end{aligned} \tag{4.8}$$

where $\sum_{\{e_{j\pm}\}}$ represents summation over possible configurations of e_{ij}'s under the neutrality conditions, and y_0 is an instanton fugacity, i.e., tunneling matrix element t corresponding to $(\theta_0,\phi_0)=(0,0)\rightarrow(\pm\pi,\pm\pi)$. Equation (4.8) can be simplified by introducing the dual fields $\tilde{\theta}_0$ and $\tilde{\phi}_0$ as

$$\begin{aligned}Z\propto&\sum_{n_1=0}^{\infty}\sum_{n_2=0}^{\infty}\sum_{\{e_{1j}\}}\sum_{\{e_{2j}\}}\frac{1}{n_1!n_2!}\int_0^\beta d\tau_{11}\cdots\int_0^\beta d\tau_{1n_1}\int_0^\beta d\tau_{21}\cdots\int_0^\beta d\tau_{2n_2}\int\mathcal{D}\tilde{\theta}_0\int\mathcal{D}\tilde{\phi}_0\\&\times y_0^{n_1+n_2}\exp\left(-\frac{\eta_\rho}{2\pi\beta}\sum_{\omega_n}|\omega_n||\tilde{\theta}_0(\omega_n)|^2-\frac{\eta_\sigma}{2\pi\beta}\sum_{\omega_n}|\omega_n||\tilde{\phi}_0(\omega_n)|^2\right.\\&\left.\qquad+i\sum_{j=1}^{n_1}e_{1j}[\tilde{\theta}_0(\tau_{1j})+\tilde{\phi}_0(\tau_{1j})]+i\sum_{j=1}^{n_2}e_{2j}[\tilde{\theta}_0(\tau_{2j})-\tilde{\phi}_0(\tau_{2j})]\right)\\=&\int\mathcal{D}\tilde{\theta}_0\int\mathcal{D}\tilde{\phi}_0\exp\left(-\frac{\eta_\rho}{2\pi\beta}\sum_{\omega_n}|\omega_n||\tilde{\theta}_0(\omega_n)|^2-\frac{\eta_\sigma}{2\pi\beta}\sum_{\omega_n}|\omega_n||\tilde{\phi}_0(\omega_n)|^2\right.\\&\left.\qquad+2y_0\int_0^\beta d\tau\cos[\tilde{\theta}_0(\tau)+\tilde{\phi}_0(\tau)]+2y_0\int_0^\beta d\tau\cos[\tilde{\theta}_0(\tau)-\tilde{\phi}_0(\tau)]\right)\\=&\int\mathcal{D}\tilde{\theta}_0\int\mathcal{D}\tilde{\phi}_0\exp\left(-\frac{\eta_\rho}{2\pi\beta}\sum_{\omega_n}|\omega_n||\tilde{\theta}_0(\omega_n)|^2-\frac{\eta_\sigma}{2\pi\beta}\sum_{\omega_n}|\omega_n||\tilde{\phi}_0(\omega_n)|^2+4y_0\int_0^\beta d\tau\cos\tilde{\theta}_0(\tau)\cos\tilde{\phi}_0(\tau)\right),\end{aligned} \tag{4.9}$$

which is identical to the original partition function (2.9) with correspondences $\eta_{\rho(\sigma)}\leftrightarrow 1/\eta_{\rho(\sigma)}$, $2y_0\leftrightarrow V_0/\pi\alpha$, $\tilde{\theta}_0\leftrightarrow\theta_0$, and $\tilde{\phi}_0\leftrightarrow\phi_0$. It is of interest to note that $\tilde{\theta}_0$ represents the Josephson phase whereas θ_0 corresponds to the charge. $\tilde{\phi}_0$ is also the conjugate variable of ϕ_0 in the same sense.

Because the partition function in the strong potential limit is found to be identical to that in the weak poten-

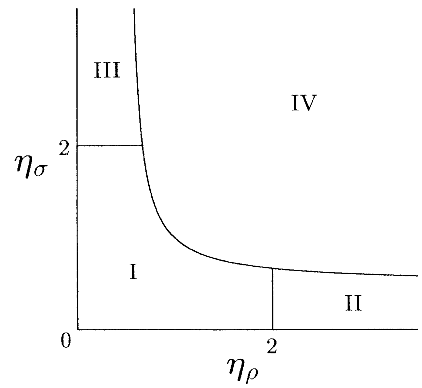

FIG. 3. The phase diagram of the ground state for the strong potential limit. The phase boundaries are $\frac{1}{\eta_\rho}+\frac{1}{\eta_\sigma}=2$, $\eta_\rho=2$, and $\eta_\sigma=2$. The charge conductance G_ρ is 0 in regions I and III, and $e^2\eta_\rho/\pi$ in regions II and IV at zero temperature. The spin conductance G_σ is 0 in regions I and II, and $\mu_B\eta_\sigma/2\pi$ in regions III and IV at zero temperature.

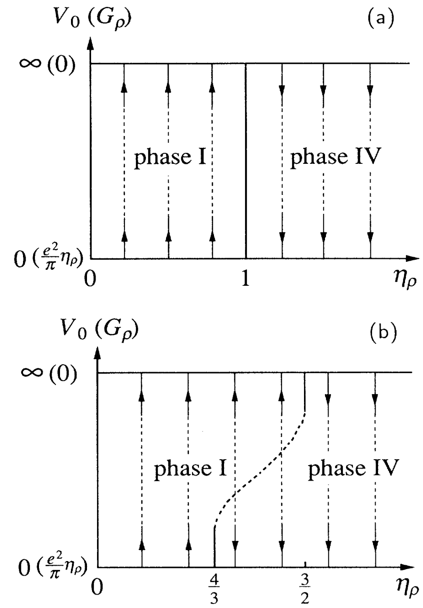

FIG. 4. The RG flow diagram for $\eta_\rho=\eta_\sigma$ (a) and $\eta_\rho=2\eta_\sigma$ (b).

tial limit, we can readily write down the scaling equations applying the analysis in the preceding section. As is shown in Sec. III, the second-order cumulant expansion of $y_0\cos\tilde{\theta}_0\cos\tilde{\phi}_0$ yields $y_{2,0}\cos 2\tilde{\theta}_0$ and $y_{0,2}\cos 2\tilde{\phi}_0$. By analogy with the fact that $y_0\cos\tilde{\theta}_0\cos\tilde{\phi}_0$ represents the tunneling from $(\theta_0,\phi_0)=(0,0)$ to $(\pm\pi,\pm\pi)$, $y_{2,0}\cos 2\tilde{\theta}_0$ and $y_{0,2}\cos 2\tilde{\phi}_0$ correspond to tunnelings from $(0,0)$ to $(\pm 2\pi,0)$ and to $(0,\pm 2\pi)$, respectively. We note here that if $(d\theta_0/d\tau)^2$ and $(d\phi_0/d\tau)^2$ in Eq. (4.1b) do not have the same coefficient m, then the effective action in Eq. (4.9) will have $y_{2,0}\cos 2\theta_0$, $y_{0,2}\cos 2\phi_0$, etc.

From Eqs. (3.9), (3.10a), and (3.10b), we obtain

$$\frac{dy_0}{dl}=\left[1-\frac{1}{2}\left(\frac{1}{\eta_\rho}+\frac{1}{\eta_\sigma}\right)\right]y_0(l), \tag{4.10a}$$

$$\frac{dy_{2,0}}{dl}=\left(1-\frac{2}{\eta_\rho}\right)y_{2,0}(l), \tag{4.10b}$$

$$\frac{dy_{0,2}}{dl}=\left(1-\frac{2}{\eta_\sigma}\right)y_{0,2}(l), \tag{4.10c}$$

from which we deduce the phase diagram at $T=0\,\mathrm{K}$ (Fig. 3). The ground state is classified into four regions, and phase boundaries are $\frac{1}{\eta_\rho}+\frac{1}{\eta_\sigma}=2$, $\eta_\rho=2$, and $\eta_\sigma=2$. In region I all the fugacities (tunneling matrix elements) scale to zero, which means that no tunneling occurs at $T=0\,\mathrm{K}$. In region II (III) only $y_{2,0}$ ($y_{0,2}$) scales to a larger value, which means that only the singlet electron pair (triplet electron-hole pair) can tunnel although the individual electron is perfectly reflected by the barrier at $T=0\,\mathrm{K}$. This corresponds to the fact that in this region II (III) the singlet superconductivity (spin density wave) instability is the most enhanced one for the 1D system without impurities.[2] Lastly, in region IV y_0 scales to be larger so that the barrier transmits electrons perfectly at $T=0\,\mathrm{K}$. The phase diagram is qualitatively the same as that in the weak potential limit (Fig. 1). In contrast to the spinless model,[9] the phase boundaries change as V_0 increases from Fig. 1 to Fig. 3; the pinning regions (I, II, and III) expand as the potential barrier becomes higher. From Figs. 1 and 3 we can deduce the RG flows (Fig. 4). Here the essential point is that η_ρ and η_σ are not renormalized so that the RG flows are all vertical.[16] When $\eta_\rho=\eta_\sigma$ [Fig. 4(a)], the phase boundary is vertical at $\eta_\rho=\eta_\sigma=1$, and the RG flows are reminiscent of those of a quantum Brownian particle in a cosine potential[15,16] as well as of the spinless fermion model.[9] The same flow diagram is obtained for the case of $\eta_\sigma=1$, i.e., the case where the system has an SU(2) spin symmetry. The noninteracting Fermi liquid ($\eta_\rho=\eta_\sigma=1$) is just on the vertical phase boundary, where the barrier potential is a marginal perturbation. In the general case, however, the phase boundary is not vertical and looks like an unstable fixed line. In Fig. 4(b) we show the RG flows along the line $\eta_\rho=2\eta_\sigma$.

B. Conductance

In this section we calculate the charge and spin conductances perturbatively in powers of the tunneling matrix element t from the golden rule.

As shown by Caldeira and Leggett,[13] the dissipation suffered by the particle of coordinate (θ_0,ϕ_0) in the partition function (2.9) can be expressed as the linear coupling with harmonic oscillators:

$$Z\propto\int\prod_j\mathcal{D}x_{1j}\int\prod_k\mathcal{D}x_{2k}\int\mathcal{D}\theta_0\int\mathcal{D}\phi_0\exp\left(-\int_0^\beta d\tau L(\{x_{1j}\},\{x_{2k}\};\theta_0,\phi_0)\right), \tag{4.11}$$

where

$$L(\{x_{1j}\},\{x_{2k}\};\theta_0,\phi_0)=\sum_j\left(\frac{1}{2}m_{1j}\left(\frac{dx_{1j}}{d\tau}\right)^2+\frac{1}{2}m_{1j}\omega_{1j}^2x_{1j}^2+g_{1j}x_{1j}\theta_0+\frac{g_{1j}^2}{2m_{1j}\omega_{1j}^2}\theta_0^2\right)$$
$$+\sum_k\left(\frac{1}{2}m_{2k}\left(\frac{dx_{2k}}{d\tau}\right)^2+\frac{1}{2}m_{2j}\omega_{2k}^2x_{2k}^2+g_{2j}x_{2k}\phi_0+\frac{g_{2k}^2}{2m_{2k}\omega_{2k}^2}\phi_0^2\right)+\frac{2V_0}{\pi\alpha}(1-\cos\theta_0\cos\phi_0) \tag{4.12}$$

with spectral functions for the harmonic oscillators $\{x_{1j}\}$ and $\{x_{2k}\}$,

$$J_1(\omega)=\sum_j\frac{\pi g_{1j}^2}{2m_{1j}\omega_{1j}}\delta(\omega-\omega_{1j})=\frac{\omega}{\pi\eta_\rho}\Theta(\omega), \tag{4.13a}$$

$$J_2(\omega)=\sum_k\frac{\pi g_{2k}^2}{2m_{2k}\omega_{2k}}\delta(\omega-\omega_{2k})=\frac{\omega}{\pi\eta_\sigma}\Theta(\omega). \tag{4.13b}$$

The tunneling probability through the potential barrier is calculated from the overlap between the initial state and the final state. The probability of the tunneling from $(\theta_0,\phi_0)=(0,0)$ to (π,π) is thus given by

$$P_{(0,0)\to(\pi,\pi)}=2\pi t^2\sum_{i,f}|\langle f|i\rangle|^2e^{-\beta E_i}\delta(E_f-E_i-eV)\Big/\sum_i e^{-\beta E_i}$$
$$=t^2\int_{-\infty}^{\infty}dt_0\langle e^{-iH_ft_0}e^{iH_it_0}\rangle_i e^{ieVt_0}, \tag{4.14}$$

where V is the applied voltage and $|i\rangle$ ($|f\rangle$) represents eigenstates of H_i (H_f) with energy E_i (E_f). The initial- and final-state Hamiltonian are obtained from $L(\{x_{1j}\},\{x_{2k}\};\theta_0,\phi_0)$ by setting $(\theta_0,\phi_0)=(0,0)$ and (π,π), respectively:

$$H_i=\sum_j\omega_{1j}(a_j^\dagger a_j+\tfrac{1}{2})+\sum_k\omega_{2k}(b_k^\dagger b_k+\tfrac{1}{2}), \tag{4.15a}$$

$$H_f=\sum_j\left[\omega_{1j}(a_j^\dagger a_j+\tfrac{1}{2})+\frac{\pi g_{1j}}{\sqrt{2m_{1j}\omega_{1j}}}(a_j+a_j^\dagger)+\frac{\pi^2g_{1j}^2}{2m_{1j}\omega_{1j}^2}\right]$$
$$+\sum_k\left[\omega_{2k}(b_k^\dagger b_k+\tfrac{1}{2})+\frac{\pi g_{2k}}{\sqrt{2m_{2k}\omega_{2k}}}(b_k+b_k^\dagger)+\frac{\pi^2g_{2k}^2}{2m_{2k}\omega_{2k}^2}\right], \tag{4.15b}$$

where a_j and b_k ($a_j^\dagger$ and $b_k^\dagger$) are the annihilation (creation) operator for the mode j and k, respectively. The thermal average in Eq. (4.14) is performed with respect to H_i as $\langle X\rangle_i=\mathrm{Tr}(Xe^{-\beta H_i})/\mathrm{Tr}(e^{-\beta H_i})$. The above two Hamiltonians are related to each other by $H_f=U^\dagger H_iU$, where the unitary operator U is given by

$$U=\exp\left[\sum_j\frac{\pi g_{1j}}{\sqrt{2m_{1j}\omega_{1j}^3}}(a_j^\dagger-a_j)+\sum_k\frac{\pi g_{2k}}{\sqrt{2m_{2k}\omega_{2k}^3}}(b_k^\dagger-b_k)\right]. \tag{4.16}$$

With the help of this relation, Eq. (4.14) is evaluated as

$$P_{(0,0)\to(\pi,\pi)}=t^2\int_{-\infty}^{\infty}dt_0\exp\left[ieVt_0-\pi\int_0^\infty\frac{d\omega}{\omega^2}[J_1(\omega)+J_2(\omega)]\left((1-\cos\omega t_0)\coth\frac{\beta\omega}{2}+i\sin\omega t_0\right)\right]. \tag{4.17}$$

In the same way, the probability of the reverse process, $(\theta_0,\phi_0)=(\pi,\pi)\to(0,0)$, is obtained as

$$P_{(\pi,\pi)\to(0,0)}=t^2\int_{-\infty}^{\infty}dt_0\langle e^{-iH_it_0}e^{iH_ft_0}\rangle_f e^{-ieVt_0}$$
$$=e^{-\beta eV}P_{(0,0)\to(\pi,\pi)}, \tag{4.18}$$

where the last line represents the detailed balance. The difference between $P_{(0,0)\to(\pi,\pi)}$ and $P_{(\pi,\pi)\to(0,0)}$ amounts to the net charge current J_ρ:

$$J_\rho=2e\,(P_{(0,0)\to(\pi,\pi)}-P_{(\pi,\pi)\to(0,0)})$$
$$=2et^2(1-e^{-\beta eV})\int_{-\infty}^{\infty}dt_0\exp\left[ieVt_0-\left(\frac{1}{\eta_\rho}+\frac{1}{\eta_\sigma}\right)\int_0^\infty\frac{d\omega}{\omega}\left((1-\cos\omega t_0)\coth\frac{\beta\omega}{2}+i\sin\omega t_0\right)\right], \tag{4.19}$$

where the prefactor 2 comes from the spin degeneracy. Hence, in the lowest order, the charge conductance G_ρ is given by

$$G_\rho = 2e^2t^2\beta \int_{-\infty}^{\infty} dt_0 \exp\left[-\left(\frac{1}{\eta_\rho}+\frac{1}{\eta_\sigma}\right)\int_0^\infty d\omega \frac{e^{-\omega/\Lambda}}{\omega}\left((1-\cos\omega t_0)\coth\frac{\beta\omega}{2}+i\sin\omega t_0\right)\right], \quad (4.20)$$

where we have introduced an exponential cutoff $e^{-\omega/\Lambda}$ to avoid ultraviolet divergences. At low temperatures Eq. (4.20) is estimated as

$$G_\rho = d_1 e^2 \left(\frac{t}{\Lambda}\right)^2 \left(\frac{\pi T}{\Lambda}\right)^{\frac{1}{\eta_\rho}+\frac{1}{\eta_\sigma}-2}, \quad (4.21)$$

where d_1 is given by $2\pi^{3/2}\Gamma(\frac{1}{2\eta_\rho}+\frac{1}{2\eta_\sigma})/\Gamma(\frac{1}{2\eta_\rho}+\frac{1}{2\eta_\sigma}+\frac{1}{2})$. The next-order term of the charge conductance is due to the tunneling from $(\theta_0,\phi_0)=(0,0)$ to $(2\pi,0)$, whose tunneling matrix element, t_2, is of the order of t^2/Λ. The probability for this tunneling process is then obtained as

$$P_{(0,0)\to(2\pi,0)} = t_2^2\int_{-\infty}^{\infty} dt_0 \exp\left[2ieVt_0 - 4\pi\int_0^\infty \frac{d\omega}{\omega^2}J_1(\omega)\left((1-\cos\omega t_0)\coth\frac{\beta\omega}{2}+i\sin\omega t_0\right)\right]. \quad (4.22)$$

The probability for the reverse process is obtained from the relation $P_{(2\pi,0)\to(0,0)} = e^{-2\beta eV}P_{(0,0)\to(2\pi,0)}$. Thus the charge conductance due to these tunneling processes becomes

$$2e^2t_2^2\beta\int_0^\infty dt_0 \exp\left[-\frac{4}{\eta_\rho}\int_0^\infty d\omega\frac{e^{-\omega/\Lambda}}{\omega}\left((1-\cos\omega t_0)\coth\frac{\beta\omega}{2}+i\sin\omega t_0\right)\right]. \quad (4.23)$$

Combining Eqs. (4.21) and (4.23), we get the charge conductance, up to the order of $(t/\Lambda)^4$, as

$$G_\rho = d_1 e^2\left(\frac{t}{\Lambda}\right)^2\left(\frac{\pi T}{\Lambda}\right)^{\frac{1}{\eta_\rho}+\frac{1}{\eta_\sigma}-2} + d_2 e^2\left(\frac{t_2}{\Lambda}\right)^2\left(\frac{\pi T}{\Lambda}\right)^{\frac{4}{\eta_\rho}-2}, \quad (4.24)$$

where d_2 is a dimensionless number of order unity.

The current-voltage characteristic at zero temperature is also obtained from Eqs. (4.19) and (4.22) as

$$\begin{aligned} J_\rho &= 2et^2\int_{-\infty}^{\infty} dt_0 \exp\left[ieVt_0 - \left(\frac{1}{\eta_\rho}+\frac{1}{\eta_\sigma}\right)\int_0^\infty d\omega\frac{e^{-\omega/\Lambda}}{\omega}(1-e^{-i\omega t_0})\right] \\ &\quad + et_2^2\int_{-\infty}^{\infty} dt_0\exp\left[2ieVt_0 - \frac{4}{\eta_\rho}\int_0^\infty d\omega\frac{e^{-\omega/\Lambda}}{\omega}(1-e^{-i\omega t_0})\right] \\ &= \frac{4\pi e t^2}{\Lambda\Gamma(\frac{1}{\eta_\rho}+\frac{1}{\eta_\sigma})}\left(\frac{eV}{\Lambda}\right)^{\frac{1}{\eta_\rho}+\frac{1}{\eta_\sigma}-1} + \frac{2\pi e t_2^2}{\Lambda\Gamma(\frac{4}{\eta_\rho})}\left(\frac{2eV}{\Lambda}\right)^{\frac{4}{\eta_\rho}-1}, \end{aligned} \quad (4.25)$$

where we have neglected exponential factors that reduce to unity as $eV/\Lambda \to 0$. Equation (4.25) shows that the tunneling is suppressed in the charge-pinning regions I and III: $J_\rho \propto V^g$ $(g>1)$. This result is reminiscent of recent theories on the effect of electromagnetic environment on the Coulomb blockade in a single tunnel junction.[22,23] In our model the many-body correlations suppress the tunneling. In regions II and IV, on the other hand, Eq. (4.25) tells that the tunneling is enhanced to give $J_\rho \propto V^g$ with $g<1$. This is, however, not the case; the enhancement suggests that the expansion in powers of t is not valid, and rather the expansion in powers of V_0 described in Sec. III becomes appropriate.

We can also evaluate the spin conductance G_σ in the same way. The lowest-order conductance is obtained again from $P_{(0,0)\to(\pi,\pi)}$ and $P_{(\pi,\pi)\to(0,0)}$, but the relation between the two probabilities is now given by $P_{(\pi,\pi)\to(0,0)} = e^{-\beta\mu_B H/2}P_{(\pi,\pi)\to(0,0)}$, where H is the magnetic field difference across the barrier. The next-order term is obtained by examining the tunneling from $(\theta_0,\phi_0)=(0,0)$ to $(0,2\pi)$. Hence the spin conductance is calculated up to the order of $(t/\Lambda)^4$ as

$$\begin{aligned} G_\sigma &= d_1'\mu_B\left(\frac{t}{\Lambda}\right)^2\left(\frac{\pi T}{\Lambda}\right)^{\frac{1}{\eta_\rho}+\frac{1}{\eta_\sigma}-2} \\ &\quad + d_2'\mu_B\left(\frac{t_2}{\Lambda}\right)^2\left(\frac{\pi T}{\Lambda}\right)^{\frac{4}{\eta_\sigma}-2}, \end{aligned} \quad (4.26)$$

where d_1' and d_2' are dimensionless numbers.

Equations (4.24) and (4.26) are correct low-temperature expansions for the conductances in the parameter regions where the corresponding phase field is pinned at zero temperature: the expansion is valid in regions I and III of Fig. 3 for G_ρ and in I and II for G_σ. In the other regions, II and IV for G_ρ and III and IV for G_σ, as the temperature is lowered, the tunneling probabilities scale to infinity while the potential V_0 scales to zero. Thus, in this case the perturbative calculations in powers of V_0 become appropriate for low temperatures.

V. ANDERSON LOCALIZATION

Now we discuss the relevance of the above 0D results to the Anderson localization in 1D electron systems.[4–6] We have shown in Sec. III that in the weak-potential limit the phase boundary of the single-barrier problem is $\eta_\rho + \eta_\sigma = 2$. In the previous studies on the Anderson localization,[4–6] on the other hand, the phase boundary of the localization transition is shown to be $\eta_\rho + \eta_\sigma = 3$. We will show below that our strong-pinning picture holds true for high temperatures but gives way to the weak-pinning picture at low temperatures, and that this crossover occurs around a temperature comparable with the discretization energy.

Suppose that N impurities are distributed dilutely at $x = x_j$ ($j = 1, 2, \ldots, N$) with average interval R. Then the partition function of the system is given by

$$Z_0 = \int \mathcal{D}\theta \int \mathcal{D}\phi \exp\left(- \int_0^\beta d\tau [L_0(\tau) + L'(\tau)] \right), \quad (5.1)$$

where L_0 is the pure Lagrangian (2.2) and L' is given by

$$L' = -\sum_{x_j} \frac{2V_j}{\pi\alpha} \cos[\theta(x_j,\tau) + 2k_F x_j] \cos[\phi(x_j,\tau)]. \quad (5.2)$$

In the same way as in Sec. II, we introduce the phase fields at impurity sites, $\theta_j(\tau)$ and $\phi_j(\tau)$, and auxiliary fields, $\lambda_{1j}(\tau)$ and $\lambda_{2j}(\tau)$, and then integrate out $\theta(x,\tau)$ and $\phi(x,\tau)$. The result is a generalization of Eq. (2.7),

$$\begin{aligned} Z_0 = \int \prod_j \mathcal{D}\theta_j \mathcal{D}\phi_j \mathcal{D}\lambda_{1j} \mathcal{D}\lambda_{2j} \exp\Bigg(& -\frac{\pi\eta_\rho}{2\beta} \sum_{\omega_n} \sum_{x_j} \sum_{x_k} \frac{1}{|\omega_n|} \exp[-|\omega_n(x_j - x_k)|/v_\rho] \lambda_{1j}(\omega_n)\lambda_{1k}(-\omega_n) \\ & -\frac{\pi\eta_\sigma}{2\beta} \sum_{\omega_n} \sum_{x_j} \sum_{x_k} \frac{1}{|\omega_n|} \exp[-|\omega_n(x_j - x_k)|/v_\sigma] \lambda_{2j}(\omega_n)\lambda_{2k}(-\omega_n) \\ & +\frac{i}{\beta} \sum_{\omega_n} \sum_{x_j} [\lambda_{1j}(\omega_n)\theta_j(-\omega_n) + \lambda_{2j}(\omega_n)\phi_j(-\omega_n)] \\ & +\sum_{x_j} \frac{2V_j}{\pi\alpha} \int_0^\beta d\tau \cos[\theta_j(\tau) + 2k_F x_j] \cos\phi_j(\tau) \Bigg), \end{aligned} \quad (5.3)$$

where we have used the relation

$$\int_{-\infty}^{\infty} \frac{dq}{2\pi} \frac{e^{-iqx}}{\omega_n^2 + v^2 q^2} = \frac{1}{2v|\omega_n|} \exp(-|\omega_n x|/v). \quad (5.4)$$

Assuming that the randomness of impurity distribution affects the electronic transport mainly through the random distribution of the phase $2k_F x_j$, we now approximate the x_j's in $\exp[-|\omega_n(x_j - x_k)|/v_{\rho(\sigma)}]$ by $x_j = jR$ ($j = 1, 2, \ldots, N$) while keeping $2k_F x_j$ to be a random variable. Then we introduce the Fourier transforms,

$$\theta_j(\omega_n) = \frac{1}{\sqrt{N}} \sum_q e^{ijqR} \theta(q,\omega_n), \qquad \phi_j(\omega_n) = \frac{1}{\sqrt{N}} \sum_q e^{ijqR} \phi(q,\omega_n), \quad (5.5a)$$

$$\lambda_{1j}(\omega_n) = \frac{1}{\sqrt{N}} \sum_q e^{ijqR} \lambda_1(q,\omega_n), \qquad \lambda_{2j}(\omega_n) = \frac{1}{\sqrt{N}} \sum_q e^{ijqR} \lambda_2(q,\omega_n), \quad (5.5b)$$

where q belongs to the first Brillouin zone ($-\pi/R \le q \le \pi/R$). Substituting Eqs. (5.5a) and (5.5b) into Eq. (5.3) and integrating out $\lambda_1(q,\omega_n)$ and $\lambda_2(q,\omega_n)$, we get the effective Euclidean action for $\theta(q,\omega_n)$ and $\phi(q,\omega_n)$ as

$$S_{\text{eff}} = \frac{1}{\beta} \sum_{\omega_n} \sum_q [K_\rho(q,\omega_n)|\theta(q,\omega_n)|^2 + K_\sigma(q,\omega_n)|\phi(q,\omega_n)|^2] - \sum_{x_j} \frac{2V_j}{\pi\alpha} \int_0^\beta d\tau \cos[\theta_j(\tau) + 2k_F x_j] \cos[\phi_j(\tau)] \quad (5.6)$$

with

$$K_\gamma(q, i\omega_n) = \frac{|\omega_n|}{2\pi\eta_\gamma} \frac{1 - 2\exp(-R|\omega_n|/v_\gamma)\cos qR + \exp(-2R|\omega_n|/v_\gamma)}{1 - \exp(-2R|\omega_n|/v_\gamma)} \qquad (\gamma = \rho, \sigma). \quad (5.7)$$

It is easily seen that the kernel $K_\gamma(q,\omega_n)$ is approximated as

$$K_\gamma(q,\omega_n) \approx \begin{cases} \dfrac{|\omega_n|}{2\pi\eta_\gamma} & \text{for } |\omega_n| \gg v_\gamma/R, \\ \dfrac{R}{4\pi v_\gamma \eta_\gamma}(\omega_n^2 + v_\gamma^2 q^2) & \text{for } |\omega_n| \ll v_\gamma/R \text{ and } |qR| \ll \pi. \end{cases} \quad (5.8)$$

Now we shall assume that v_ρ and v_σ are of the same order of magnitude ($\sim v$). Then v/R is the discretization energy within the interval between two neighboring impurities. For $|\omega_n| \gg v/R$ the correlations between the different impurities are unimportant, and the effective action is just the sum of the action in Eq. (2.9) with respect to the impurity sites with a trivial modification, $V_0 \cos\theta_0(\tau)\cos\phi_0(\tau) \to V_j \cos[\theta_j(\tau) + 2k_F x_j]\cos\phi_j(\tau)$. On the other hand, if $|\omega_n| \ll v/R$, the correlation between the impurity sites must be properly treated. The effective action describing the long-wavelength ($|q| < \pi/R$) and low-frequency ($|\omega_n| < v/R$) phenomena is obtained by integrating out the high-frequency ($|\omega_n| > v/R$) components by the RG method for the single-impurity problem discussed above. When the impurity potential V_j is weak enough, it is renormalized to $\tilde{V}_j = V_j(R/a)^{1-\frac{1}{2}(\eta_\rho+\eta_\sigma)}$ with a being the lattice spacing. This $\tilde{V}_j$ exists at every site j in this coarse grained system, and we can now apply the previous analysis[4–6] assuming a random potential expressed as a continuous function of x. Thus we can expect to see a crossover from the single-impurity behavior to the dense-impurity behavior, when the relevant energy scale or, equivalently, the temperature is varied across $T_{\rm dis} = v/k_B R$. We note that this crossover temperature $T_{\rm dis}$ is higher than $T_{\rm loc}$ because the localization length $L_{\rm loc}$ is always longer than R. It may be instructive to show that this RG procedure is compatible with the previous theories on the localization length. According to the weak-pinning analysis,[5] the localization length is given by

$$L_{\rm loc} \sim n_i^{-1}(n_i\alpha)^{\frac{2-\eta}{3-\eta}}\left(\frac{V_0}{v}\right)^{-\frac{2}{3-\eta}}, \tag{5.9}$$

where $\eta = \eta_\rho + \eta_\sigma$, $n_i = 1/R$, and α is the short-distance cutoff. By integrating out the high-frequency modes $|\omega_n| > v/R$, the potential is renormalized to $\tilde{V}_j = V_0(R/a)^{1-\frac{1}{2}\eta}$. Substituting $\tilde{V}_j$ for V_0 and R for α in Eq. (5.9), we get

$$L_{\rm loc} \sim R\left[\frac{V_0}{v}\left(\frac{R}{a}\right)^{1-\frac{1}{2}\eta}\right]^{-\frac{2}{3-\eta}} = R\left(\frac{R}{a}\right)^{-\frac{2-\eta}{3-\eta}}\left(\frac{V_0}{v}\right)^{-\frac{2}{3-\eta}}, \tag{5.10}$$

which is exactly the same as Eq. (5.9) with $\alpha = a$. In this way our RG procedure matches the previous result.

For homogeneous 1D systems a physically observed quantity is resistivity $\rho(T)$, which is related to the conductance by $\rho(T) \sim (LG_\rho)^{-1}$ where L is some characteristic length scale. L is estimated as R for $T > T_{\rm dis}$, v_F/T for $T_{\rm loc} < T < T_{\rm dis}$, and $L_{\rm loc}$ for $T < T_{\rm loc}$, while for $T < T_{\rm dis}$ G_ρ^{-1} is proportional to $V(T)^2$ with $V(T)$ being the renormalized potential strength down to the energy scale of the order of $k_B T$. As a particular example, suppose that interaction parameters lie in the range $2 < \eta_\rho + \eta_\sigma < 3$. At high temperature $T > T_{\rm dis}$ both charge and spin phase fields are not pinned, and the resistivity $\rho(T)$ is proportional to the inverse of the conductance of a single impurity. As discussed by Fisher and Zwerger,[16] when $\eta_\rho + \eta_\sigma > 2$, the resistivity is nonmonotonous as a function of the temperature, showing a maximum at some temperature T^*. In the high temperature limit, it approaches $\rho_\infty \equiv (Re^2\eta_\rho/\pi)^{-1}$. Below T^* the resistivity decreases again toward ρ_∞ with decreasing temperature. When the temperature is further reduced below $T_{\rm dis}$, the resistivity changes to decrease to zero as $\rho(T) \propto T^{\eta_\rho+\eta_\sigma-2}$ and then turns to increase around $T_{\rm loc}$ as $\rho(T) \propto T^{\eta_\rho+\eta_\sigma-3}$, as the phase fields begin to be pinned.[6] Thus the temperature dependence of the resistivity has a fairly complicated structure with two crossover temperatures, $T_{\rm loc}$ and $T_{\rm dis}$, and this scenario can be checked experimentally by changing the concentration of the impurities, i.e., R^{-1}. Schematic temperature dependence of the resistivity for general cases is shown in Fig. 5.

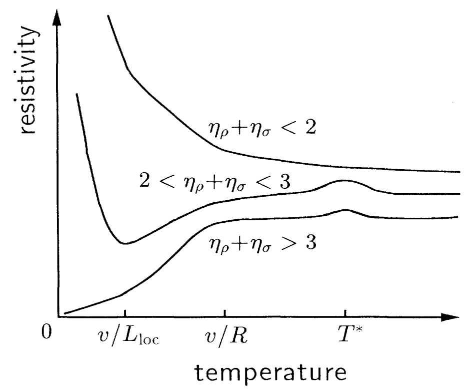

FIG. 5. The schematic temperature dependence of the resistivity. In the high temperature limit it approaches $\rho_\infty = (Re^2\eta_\rho/\pi)^{-1}$. For $\eta_\rho + \eta_\sigma < 3$, it has a maximum at some temperature $T^*(> T_{\rm dis})$. Below $T_{\rm dis}$, $\rho(T) \propto T^{\eta_\rho+\eta_\sigma-2}$ for $T_{\rm loc} < T < T_{\rm dis}$ and $\rho(T) \propto T^{\eta_\rho+\eta_\sigma-3}$ for $T < T_{\rm loc}$.

VI. CONCLUSIONS

In this paper we have investigated a model for the transport through a single barrier in a 1D spin-$\frac{1}{2}$ interacting electron system. We have obtained the phase diagram of the ground state for both weak and strong potential barrier, classifying in terms of the charge and spin conductances at zero temperature. We find regions where only the charge or the spin can transport coherently, in addition to the insulating region ($G_\rho = G_\sigma = 0$) and the perfect-conductor region ($G_\rho = e^2\eta_\rho/\pi$ and $G_\sigma = \mu_B\eta_\sigma/2\pi$). We have obtained correct low-temperature expansions of the charge and spin conductances. The results are applied to the problem of the Anderson localization for the energy and/or temperature higher than the discretization energy v/R. We have shown that the transport in a 1D system with dilute impurities changes qualitatively around $T_{\rm dis} = v/R$ in addition to $T_{\rm loc} = v/L_{\rm loc}$. We have also pointed out a possible anomalous temperature dependence of the resistivity due to this crossover, which may be observed experimentally, probably with narrow 1D quantum wires.

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