Electronic dephasing in wires due to metallic gates

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The dephasing effect of metallic gates on electrons moving in one quasi-one-dimensional diffusive wire is analyzed. The incomplete screening in this geometry implies that the effect of the gate can be described, at high energies or temperatures, as an electric field fluctuating in time. The resulting system can be considered a realization of the Caldeira-Leggett model of an environment coupled to many particles. Within the range of temperatures where this approximation is valid, a simple estimation of the inverse dephasing time gives $\tau_{\mathrm{G}}^{-1} \propto T^{1/2}$.

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I. INTRODUCTION

The low temperature dephasing time of electrons in diffusive metals has attracted a great deal of attention.[1–4] Different mechanisms have been proposed to explain the anomalous dephasing properties reported in these experiments. Some of them are extrinsic, like dynamical defects,[5] two level systems,[6] or magnetic impurities.[7] Alternatively, intrinsic effects have also been proposed.[8] Screening effects in a quasi-one-dimensional geometry are significantly reduced, leading to a breakdown of Fermi liquid theory at low temperatures[9,10] (see also Ref. 11). Decoherence can be induced by inelastic processes within the system under study, or by fluctuating potentials induced by external sources. Metallic gates are perhaps the most ubiquitous source of external dephasing, as their gapless spectrum leads to fluctuations at all time or energy scales, and they are electromagnetically coupled to other metallic systems. In the following we study the dephasing induced by metallic gates on quasi-one-dimensional diffusive wires. The study follows the analysis in Ref. 12, where dephasing effects in ballistic quantum dots was considered.

Using electrostatic arguments, it is easy to show that the fields induced by a fluctuation of charge in the gate decay exponentially outside the gate, with a decay length proportional to the wavelength of the fluctuation. Then, the presence of the gate implies the existence of two regimes, which depend on the relative value of the distance over which the electrons in the wire diffuse and the wire-gate distance. We can also define the two regimes in terms of the energy or temperature required to cover a distance of the order of the separation between the wire and the gate.

(i) For distances along the wire $L \ll z$, or time scales larger than $\mathcal{D}_w^{-1} z^2$, where $\mathcal{D}_w$ is the diffusion coefficient of the wire, and z is the distance to the gate, the dephasing time is the sum of a contribution from charge fluctuations within the wire and another due to the fluctuations at the gate. The contribution from the gate comes from fluctuations in the charge density of the gate of wavelength comparable or larger than z. The fields induced by fluctuations of shorter wavelengths decay at distances shorter than the wire-gate separation, and their influence on the wire can be neglected (see below). The electric fields induced by the gate can be calculated within the dipolar approximation. Because of the one-dimensional geometry of the wire, this field is not screened by the charge fluctuations of the wire. As discussed in Ref. 12, this coupling can be considered a generalization to a many particle system of the Caldeira-Leggett model of Ohmic dissipation.[13] This model shows anomalous dephasing in many situations.[14–17]

(ii) At distances $L \gg z$ or time scales lower than $\mathcal{D}_w^{-1} z^2$ the distance between the wire and the gate can be neglected. The screening by the gate leads to an effective short range potential along the wire.[18]

Section II describes mathematically the model to be studied. The different regimes mentioned above are discussed in Sec. III. Section IV generalizes the results to gates with geometries which differ from that depicted in Fig. 1. Finally, Sec. V contains a discussion of the most relevant results. The units are such that $\hbar = 1$.

II. THE MODEL

We study the setup sketched in Fig. 1. A wire of width w is located at height z over a two-dimensional metallic gate of width w'. The effects of the finite width of both systems is included through the densities of states, ν_{w} and ν_{G}, defined as number of states per unit length and per unit area, respectively. We study the contribution to dephasing from the gate using the scheme proposed in Ref. 11. The probability of transition of a particle at the Fermi level to other states, after time t, using second order perturbation theory, at temperature $T = \beta^{-1}$, is

$$\mathcal{P}^{(2)}(t) \simeq \int_0^t d\tau \int_0^t d\tau' \int d\mathbf{q}_{\parallel} \int_{|\omega|>1/t} d\omega e^{i\mathbf{q}_{\parallel}[\mathbf{r}(\tau)-\mathbf{r}(\tau')]+i\omega(\tau-\tau')} \times \frac{|\omega|}{1-e^{-\beta\omega}} \operatorname{Im}[v_{\mathrm{scr}}(\mathbf{q},\omega)], \tag{1}$$

where $\mathbf{q}_{\parallel}$ is the momentum in the direction parallel to the wire, and $v_{\mathrm{scr}}(\mathbf{q}_{\parallel},\omega)$ is the screened potential between points within the wire. In general, we can write the screened potential as $v_{\mathrm{scr}}(\bar{z},\mathbf{q}_{\parallel},\mathbf{q}_{\perp},\omega)$, where $\bar{z}$ is the vertical coordinate, and $\mathbf{q}_{\perp}$ is the momentum in the direction perpendicular to the wire and parallel to the gate. We assume that the wire is

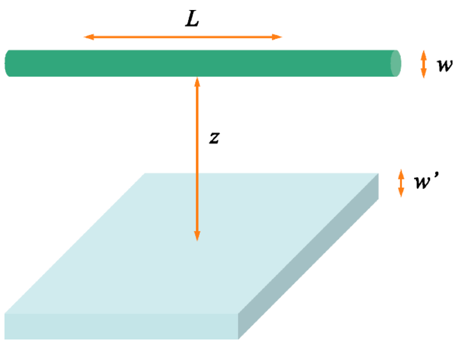

FIG. 1. Sketch of the system considered in the text. L denotes a distance along the wire, as discussed in the text.

at $\bar{z}=z$, and the gate is at $\bar{z}=0$. We define the bare, unscreened potential as $v_0(\bar{z},\mathbf{q}_\parallel,\mathbf{q}_\perp,\omega)$. The screened potentials at the position of the gate and at the position of the wire can be written as

$$v_{\rm scr}(z,\mathbf{q}_\parallel,\mathbf{q}_\perp,\omega) = v_0(z,\mathbf{q}_\parallel,\mathbf{q}_\perp,\omega) + \frac{2\pi e^2}{|\mathbf{q}|}\chi_w(\mathbf{q}_\parallel,\omega) \times \int d\mathbf{q}_\perp v_{\rm scr}(z,\mathbf{q}_\parallel,\mathbf{q}_\perp,\omega) + \frac{2\pi e^2 e^{-|\mathbf{q}|z}}{|\mathbf{q}|}\chi_G(\mathbf{q}_\parallel,\mathbf{q}_\perp,\omega)v_{\rm scr}(0,\mathbf{q}_\parallel,\mathbf{q}_\perp,\omega),$$

$$v_{\rm scr}(0,\mathbf{q}_\parallel,\mathbf{q}_\perp,\omega) = v_0(0,\mathbf{q}_\parallel,\mathbf{q}_\perp,\omega) + \frac{2\pi e^2 e^{-|\mathbf{q}|z}}{|\mathbf{q}|}\chi_w(\mathbf{q}_\parallel,\omega) \times \int d\mathbf{q}_\perp v_{\rm scr}(z,\mathbf{q}_\parallel,\mathbf{q}_\perp,\omega) + \frac{2\pi e^2}{|\mathbf{q}|}\chi_G(\mathbf{q}_\parallel,\mathbf{q}_\perp,\omega)v_{\rm scr}(0,\mathbf{q}_\parallel,\mathbf{q}_\perp,\omega), \quad (2)$$

where $\mathbf{q}_\parallel$ and $\mathbf{q}_\perp$ are the momenta parallel and perpendicular to the wire, and χ_w and χ_G are the polarizabilities of the wire and the gate. They are given by

$$\chi_w(\mathbf{q}_\parallel,\omega) = -\frac{\nu_w \mathcal{D}_w \mathbf{q}_\parallel^2}{i\omega + \mathcal{D}_w \mathbf{q}_\parallel^2},$$

$$\chi_G(\mathbf{q},\omega) = -\frac{\nu_G \mathcal{D}_G \mathbf{q}^2}{i\omega + \mathcal{D}_G \mathbf{q}^2}, \quad (3)$$

where ν_w, ν_G, $\mathcal{D}_w$, and $\mathcal{D}_G$ are the densities of states and diffusion coefficients of the wire and the gate, respectively.

From the second of the equations in (2), we can write

$$v_{\rm scr}(0,\mathbf{q}_\parallel,\mathbf{q}_\perp,\omega) = \frac{v_0(0,\mathbf{q}_\parallel,\mathbf{q}_\perp,\omega) + \frac{2\pi e^2}{|\mathbf{q}|}\chi_w(\mathbf{q}_\parallel,\omega)\int d\mathbf{q}_\perp v_{\rm scr}(z,\mathbf{q}_\parallel,\mathbf{q}_\perp,\omega)}{1 - \frac{2\pi e^2}{|\mathbf{q}|}\chi G(\mathbf{q}_\parallel,\mathbf{q}_\perp,\omega)}. \quad (4)$$

A point charge at a point in the wire leads to a bare potential such that

$$v_0(z,\mathbf{q}_\parallel,\mathbf{q}_\perp,\omega) = \frac{2\pi e^2}{|\mathbf{q}|},$$

$$v_0(0,\mathbf{q}_\parallel,\mathbf{q}_\perp,\omega) = \frac{2\pi e^2 e^{-|\mathbf{q}|z}}{|\mathbf{q}|}. \quad (5)$$

Using this expression, and inserting Eq. (4) into the first equation in (2), we find

$$v_{\rm scr}(\mathbf{q}_\parallel,\omega) = \frac{e^2\left[\log\left(\frac{q_c}{\mathbf{q}_\parallel}\right) + \mathcal{F}(\mathbf{q}_\parallel,\omega)\right]}{1 + e^2\chi_w(\mathbf{q}_\parallel,\omega)\left[\log\left(\frac{q_c}{\mathbf{q}_\parallel}\right) + \mathcal{F}(\mathbf{q}_\parallel,\omega)\right]}, \quad (6)$$

where q_c is a high momentum cutoff proportional to the inverse of the width of the wire. The function $\mathcal{F}$ can be written as

$$\mathcal{F}(\mathbf{q}_\parallel,\omega) = \int d_{\mathbf{q}\perp}\frac{4\pi^2 e^2 e^{-2|\mathbf{q}|z}}{|\mathbf{q}|^2}\frac{\chi_G(\mathbf{q},\omega)}{1 - \frac{2\pi e^2}{|\mathbf{q}|}\chi_G(\mathbf{q},\omega)}. \quad (7)$$

For sufficiently low momenta, $|\mathbf{q}| \ll e^2\nu_G$, we can write

$$\mathcal{F}(\mathbf{q}_\parallel,\omega) \approx \mathcal{F}_1(\mathbf{q}_\parallel) - i\omega\mathcal{F}_2(\mathbf{q}_\parallel)$$

$$\mathcal{F}_1(\mathbf{q}_\parallel) = \int d_{\mathbf{q}\perp}\frac{2\pi^2 e^{-2|\mathbf{q}|z}}{|\mathbf{q}|} \sim \log(\mathbf{q}_\parallel z)$$

$$\mathcal{F}_2(\mathbf{q}_\parallel) \approx \frac{1}{e^2\nu_G\mathcal{D}_G}\int d_{\mathbf{q}\perp}\frac{e^{-2|\mathbf{q}|z}}{|\mathbf{q}|^2} \sim \frac{1}{e^2\nu_G\mathcal{D}_G}\frac{1}{\mathbf{q}_\parallel}. \tag{8}$$

For small frequencies, $\omega \leqslant \mathcal{D}_w\mathbf{q}_\parallel^2$, we find

$$\mathrm{Im}[v_{\mathrm{scr}}(\mathbf{q}_\parallel,\omega)] \approx \frac{\omega e^4\nu_w}{\mathcal{D}w\mathbf{q}_\parallel^2}\frac{\left[\log\left(\frac{q_c}{\mathbf{q}_\parallel}\right)+\mathcal{F}_1(\mathbf{q}_\parallel)\right]^2 + \mathcal{D}_w\mathbf{q}_\parallel^2\mathcal{F}_2(\mathbf{q}_\parallel)}{\left\{1+\nu_w e^2\left[\log\left(\frac{q_c}{\mathbf{q}_\parallel}\right)+\mathcal{F}_1(\mathbf{q}_\parallel)\right]\right\}^2}. \tag{9}$$

III. RESULTS

A. Dipolar approximation $[\tau^{-1}(T) \gg \mathcal{D}_w/z^2]$

The integral over the time difference $\tau-\tau'$ in Eq. (1) is bounded by the inverse of the temperature, β. Both z and the value of $\langle[\mathbf{r}(\tau)-\mathbf{r}(\tau')]^2\rangle$ for $\tau-\tau' \simeq \beta$ act as a lower cutoff in the integrals over $\mathbf{q}_\parallel$, so that $|\mathbf{q}_\parallel|^{-2} \ll \min\{\langle[\mathbf{r}(\tau)-\mathbf{r}(\tau')]^2\rangle, z^2\}$. When $\langle[\mathbf{r}(\tau)-\mathbf{r}(\tau')]^2\rangle \ll z^2$, we can limit the integral over $\mathbf{q}_\parallel$ to $|\mathbf{q}_\parallel| \leqslant z^{-1}$ and substitute

$$e^{i\mathbf{q}_\parallel[\mathbf{r}(\tau)-\mathbf{r}(\tau')]} \approx -\{\mathbf{q}_\parallel[\mathbf{r}(\tau)-\mathbf{r}(\tau')]\}^2. \tag{10}$$

Within these approximations, and setting $q_c^{-1} \approx w$, Eq. (9) becomes

$$\mathrm{Im}[v_{\mathrm{scr}}(\mathbf{q}_\parallel,\omega)]e^{i\mathbf{q}_\parallel[\mathbf{r}(\tau)-\mathbf{r}(\tau')]} \approx v_{1\mathrm{D}}(\mathbf{q}_\parallel,\omega)e^{i\mathbf{q}_\parallel[\mathbf{r}(\tau)-\mathbf{r}(\tau')]} + v_{2\mathrm{D}}(\mathbf{q}_\parallel,\omega),$$

$$v_{1\mathrm{D}}(\mathbf{q}_\parallel,\omega) = \frac{\omega}{\nu_w\mathcal{D}_w\mathbf{q}_\parallel^2},$$

$$v_{2\mathrm{D}}(\mathbf{q}_\parallel,\omega) = \frac{\omega|\mathbf{q}_\parallel|[\mathbf{r}(\tau)-\mathbf{r}(\tau')]^2}{e^4\nu_w^2\nu_G\mathcal{D}_G\log^2(z/w)}. \tag{11}$$

Inserting Eq. (11) in Eq. (1) one obtains that $\mathcal{P}^{(2)}(t)$ can be written as the sum of two contributions, $\mathcal{P}_w^{(2)}(t)$ and $\mathcal{P}_G^{(2)}(t)$ arising from $v_{1\mathrm{D}}$ and $v_{2\mathrm{D}}$ [note that screening effects from the gate are included in the denominator of $v_{1\mathrm{D}}$ in Eq. (12) through the term $\log(\mathbf{q}_c z)$]. The first term when inserted in Eq. (1), give the contribution to the function $\mathcal{P}_G^{(2)}(t)$ calculated in Ref. 11, leading to the standard expression for the dephasing in one-dimensional wires.

The contribution from $v_{2\mathrm{D}}$ in Eq. (11) to Eq. (1) can be written as

$$\mathcal{P}_G^{(2)}(t) \simeq \frac{T}{e^4\nu_w^2\nu_G\mathcal{D}_G z^2\log^2(z/w)}\int_0^t d\tau \times\int_0^t d\tau'\int_{1/t}^T d\omega\{[\mathbf{r}(\tau)-\mathbf{r}(\tau')]\}^2 e^{i\omega(\tau-\tau')}. \tag{12}$$

Using $\langle[\mathbf{r}(\tau)-\mathbf{r}(\tau')]^2\rangle = \mathcal{D}_w(\tau-\tau')$, we finally obtain

$$\mathcal{P}_G^{(2)}(t) \simeq \frac{T\mathcal{D}_w t^2}{e^4\nu_w^2\nu_G\mathcal{D}_G\log^2(z/w)z^2}. \tag{13}$$

From this equation we can define a dephasing time due to the presence of the gate, $\mathcal{P}_G^{(2)}(\tau_G) \approx 1$, as

$$\hbar\tau_G^{-1} \simeq \sqrt{\frac{T\mathcal{D}_w}{e^4\nu_w^2\nu_G\mathcal{D}_G\log^2(z/w)z^2}}. \tag{14}$$

On the other hand, the dephasing time due to intrinsic processes can be written as (see also Ref. 19):

$$\hbar\tau_w^{-1} \simeq \frac{T^{2/3}}{\mathcal{D}_w^{1/3}\nu_w^{2/3}}. \tag{15}$$

Because of the different temperature dependence, τ_G^{-1} is greater than τ_w^{-1} at temperatures below a value T' given by

$$T' \simeq \frac{\mathcal{D}_w^5}{\mathcal{D}_G^3\nu_G^3 e^{12}\nu_w^2\log^6(z/w)z^6}. \tag{16}$$

The approximations leading to this result are valid provided that $\tau_G^{-1} \geqslant \mathcal{D}_w/z^2$. This condition breaks down below a temperature T'' given by

$$T'' \simeq \frac{\mathcal{D}_w}{z^2}e^4\nu_w^2\nu_G\mathcal{D}_G\log^2(z/w). \tag{17}$$

The dephasing due to the gate dominates if a temperature range $T'' \leqslant T \leqslant T'$. The inequality $T'' \leqslant T'$ implies

$$1 \leqslant f = \frac{\mathcal{D}_w}{e^4\nu_w\nu_G\mathcal{D}_G z\log^2(z/w)}. \tag{18}$$

We can write this expression using as parameters the Fermi wave vectors in the wire and gate, k_F^w and k_F^G, the width of the wire and of the gate, w_w and w_G, and the mean free paths in the wire and the gate, l_w and l_G. In terms of the dimensionless Bohr radius of the wire, $r_{sw} = (m_w e^2)/(k_F^w\hbar^2)$, we can write

$$\frac{\mathcal{D}_w}{e^2} \approx \frac{l_w}{r_{sw}},$$

$$e^2\nu_w \approx r_{sw}(k_F^w w_w)^2,$$

$$\mathcal{D}_G\nu_G \approx (k_F^G l_G)(k_F^G w_G) \tag{19}$$

so that

$$\frac{\mathcal{D}_w}{e^4\nu_w\nu_G\mathcal{D}_G z} \approx \frac{1}{r_{sw}^2(k_F^w w_w)^2(k_F^G l_G)(k_F^G w_G)}\left(\frac{l_w}{z}\right). \tag{20}$$

In the presence of a dielectric between the wire and the gate with dielectric constant ϵ_0, one has to replace the electric charge e^2 by e^2/ϵ_0 in all expressions, or, alternatively, r_{sw} by $r_{sw}\epsilon_0$. Hence, Eq. (18) can only be satisfied for very clean wires, such that $z \ll l_w$, or in the presence of a large dielectric constant, ϵ_0.

B. Low temperature regime $[\tau^{-1}(T) \ll \mathcal{D}_w/z^2]$

At low temperatures the electrons diffuse coherently over distances much larger than z. The dipolar approximation, Eq. (10) cannot be made, and the cutoff in the integrals over $\mathbf{q}_{\parallel}$ is $[\mathcal{D}_w(\tau-\tau')]^{-1}$. Then, using Eq. (9) we obtain

$$\mathcal{P}_G^{(2)}(t) \simeq \frac{Tt}{e^4\nu_w^2\nu_G\mathcal{D}_G\log^2(q_c z)}\log\left(\frac{\mathcal{D}_w}{z^2 t}\right). \tag{21}$$

Neglecting logarithmic corrections, this result leads to

$$\hbar\tau_G^{-1} \simeq \frac{T}{e^4\nu_w^2\nu_G\mathcal{D}_G\log^2(q_c z)}. \tag{22}$$

In this regime there is also a contribution from the intrinsic processes, given by Eq. (15). These processes will dominate at sufficiently low temperatures, $T \leqslant T'''$ where

$$T''' = \frac{e^{12}\nu_w^4\nu_G^3\mathcal{D}_G^3\log^6(q_c z)}{\mathcal{D}_w}. \tag{23}$$

Note that when $T'' \leqslant T'''$, where T'' is given in Eq. (17) the contributions from the gate are always smaller than those coming from fluctuations in the wire. The condition $T''' \leqslant T''$ reduces to Eq. (18).

Combining the results in this section and in the preceding one, we can write

$$\mathcal{P}_G^{(2)}(t) \approx \begin{cases} \dfrac{T}{T''}\left(\dfrac{\mathcal{D}_w t}{z^2}\right)^2, & t^{-1} \ll \dfrac{\mathcal{D}_w}{z^2} \\ \dfrac{T}{T''}\dfrac{\mathcal{D}_w t}{z^2}, & t^{-1} \gg \dfrac{\mathcal{D}_w}{z^2}, \end{cases}$$

$$\mathcal{P}_{\text{int}}^{(2)}(t) \approx f^{-1}\frac{T}{T''}\left(\frac{\mathcal{D}_w t}{z^2}\right)^{3/2}, \tag{24}$$

where f is defined in Eq. (18). These expressions lead to

$$\tau_G^{-1} \approx \begin{cases} \dfrac{\mathcal{D}_w}{z^2}\left(\dfrac{T}{T''}\right)^{1/2}, & T \geqslant T'' \\ \dfrac{\mathcal{D}_w}{z^2}\dfrac{T}{T''}, & T \leqslant T'', \end{cases}$$

$$\tau_{\text{int}}^{-1} \approx \frac{\mathcal{D}_w}{z^2}\left(\frac{T}{fT''}\right)^{2/3}. \tag{25}$$

IV. EXTENSIONS TO OTHER GEOMETRIES

A. One-dimensional gate

The analysis in the two previous sections can be extended, in a straightforward way, to the case where the gate is another quasi-one-dimensional wire. The gate polarizability, defined in Eq. (3) depends only on the momentum parallel to the wire. The function $\mathcal{F}$ in Eq. (7) becomes

$$\mathcal{F}(\mathbf{q}_{\parallel},\omega) \approx e^4K_0^2(\mathbf{q}_{\parallel}z)\frac{\chi_G(\mathbf{q}_{\parallel},\omega)}{1-e^2K_0(\mathbf{q}_{\parallel}z)\chi_G(\mathbf{q}_{\parallel},\omega)}, \tag{26}$$

where $K_0(\mathbf{q}_{\parallel}z)$ is a modified Bessel function:

$$K_0(\mathbf{q}_{\parallel}z) = \int_0^{\infty} d\mathbf{q}_{\perp}\frac{\cos(\mathbf{q}_{\perp}z)}{\sqrt{\mathbf{q}_{\parallel}^2+\mathbf{q}_{\perp}^2}}. \tag{27}$$

We can, as in the preceding section, study separately the dipolar regime, $\tau^{-1} \ll \mathcal{D}_w/z^2$, and the long time regime, $\tau^{-1} \gg \mathcal{D}_w/z^2$. In the dipolar regime, using Eq. (10), we obtain

$$\tau_G^{-1} \approx \sqrt{\frac{T\nu_w\mathcal{D}_w}{e^2\nu_G\mathcal{D}_G z(\nu_w+\nu_G)^2}} \tag{28}$$

(note that now ν_w is a quasi-one-dimensional density of states). The restriction $\tau^{-1} \leqslant \mathcal{D}_w/z^2$ implies that this result is only valid for temperatures $T \geqslant T''$ where

$$T'' = \frac{\mathcal{D}_w e^2\nu_G\mathcal{D}_G(\nu_w+\nu_G)^2}{z^3\nu_w}. \tag{29}$$

At high temperatures, $T \geqslant T'$, we find $\tau_G^{-1} \leqslant \tau_{\text{int}}^{-1}$, where

$$T' = \frac{\mathcal{D}_w^5\nu_w^7}{e^6(\nu_G\mathcal{D}_G)^3(\nu_w+\nu_G)^6z^3}. \tag{30}$$

The condition required for the relevance of the dephasing due to the gate, $T'' \leqslant T'$ implies

$$1 \leqslant \frac{\mathcal{D}_w\nu_w^2}{e^2\nu_G\mathcal{D}_G(\nu_w+\nu_G)^2}. \tag{31}$$

It is interesting to note that this condition does not depend on the distance between the wire and the gate. At temperatures below T'', such that $\tau^{-1} \leqslant \mathcal{D}_w/z^2$, the combined system acts like an effective one-dimensional wire, leading to a $\tau^{-1} \propto T^{2/3}$ dependence.

B. Three-dimensional gate

We assume that quasiparticles at the gate are specularly reflected at the boundary. Then, the screening properties of the system can be calculated from the fluctuations of surface charges at the top of the gate.[20] The function $\mathcal{F}$ in Eq. (7) can be written as

$$\mathcal{F}(\mathbf{q}_{\parallel},\omega) = \int d\mathbf{q}\perp \frac{2\pi e^2 e^{-2\mathbf{q}z}}{|\mathbf{q}|}\frac{\mathcal{B}(\mathbf{q},\omega)-1}{\mathcal{B}(\mathbf{q},\omega)+1},$$

$$\mathcal{B}(\mathbf{q},\omega) = \frac{|\mathbf{q}|}{\pi}\int d\mathbf{q}_z\frac{1}{(|\mathbf{q}|^2+\mathbf{q}_z^2)\epsilon_{\text{G}}(\mathbf{q},\mathbf{q}_z,\omega)},$$

$$\epsilon_{\text{G}}(\mathbf{q},\mathbf{q}_z,\omega) = 1 + \frac{4\pi e^2}{|\mathbf{q}|^2+\mathbf{q}_z^2}\chi_{\text{G}}(\mathbf{q},\mathbf{q}_z,\omega). \tag{32}$$

Using this expression, the parameters τ_G, T', T'', and T''' which characterize the dephasing induced by the gate can be calculated. One finds that they show the same dependence as for a quasi-two-dimensional gate, with the only replacement $\nu_G^{2D} \to \nu_G^{3D} \times z$, similarly to the results discussed in Ref. 12. Qualitatively, a three-dimensional gate behaves as a two-dimensional gate of width z. The constraint which needs to be satisfied for the gate induced dephasing to be dominant is

$$1 \leqslant \frac{\mathcal{D}_w}{e^4 \nu_w \nu_G \mathcal{D}_G z^2 \log^2(z/w)}. \tag{33}$$

C. Granular gate

We now analyze the dephasing induced by a gate made up of disconnected metallic grains. We assume that each grain has volume V and that these grains have a diffusion coefficient $\mathcal{D}_G$ and a density of states ν_G, leading to an intrinsic dc conductivity $\sigma = e^2 \nu_G \mathcal{D}_G$. Their response to an applied field is characterized by their polarizability, $\mathcal{P}$ and absorption coefficient, $\gamma(\omega) \approx V\omega^2/\sigma$.[21] The dielectric constant of the granular system can be written as

$$\epsilon(\omega) = V^{-1}\left(\mathcal{P} + \frac{i\omega V}{e^2 \nu_G \mathcal{D}_G}\right). \tag{34}$$

We can insert this expression into Eqs. (32), and carry out the following steps in order to obtain the dephasing effects of the gate. The inequality which needs to be satisfied for the gate induced dephasing to prevail is

$$1 \leqslant \frac{\mathcal{D}_w}{e^4 \nu_w \nu_G \mathcal{D}_G z^2 (1 + V^{-1} P)^2}. \tag{35}$$

This expression, valid for grains larger than the mean free path, is very similar to the corresponding one for a three-dimensional gate, Eq. (33). The effects of a granular gate, however, can be greatly enhanced by the surface roughness of the grains[22] (see also Ref. 23), or for grains much smaller than the mean free path.

V. CONCLUSIONS

We have analyzed, within the standard approach to dephasing in metals[9–11] the effects of a two-dimensional diffusive gate on the coherence properties of electrons in quasi-one-dimensional wires. The method used can be easily generalized to other types of gates, like granular metals or ballistic systems. Note that the analysis presented here deals only with dephasing due to thermal fluctuations in the gate, and it does not consider other possible sources of dephasing.

At short distances, or low temperatures, the gate induces at the position of the wire an electric field, which fluctuates in time but is approximately constant in space. The resulting model can be considered an extension of the Caldeira-Leggett model[13] to a many particle system. Using the self-consistent perturbation theory to calculate the dephasing time, we find a $\tau_G^{-1} \propto T^{1/2}$ dependence. The existence of this regime requires only the validity of the dipolar approximation, see Sec. III A. Hence, it is not restricted to the diffusive 1D wire.[24] At very low temperatures, the separation between the gate and the wire becomes irrelevant, and the contribution from the gate changes to a $\tau_G^{-1} \propto T$ dependence, as for a two-dimensional metal.

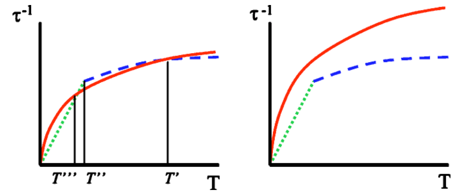

FIG. 2. Schematic representation of the intrinsic and gate contributions to the inverse dephasing time, Eq. (25). Left: constraint in Eq. (18) is satisfied. Full line, τ_{int}, Eq. (15). Broken line, τ_G, Eq. (14). Dotted line, τ_G, Eq. (22). The values of the temperatures T', T'', and T''' are given in Eqs. (16), (17), and (23). Right: constraint in Eq. (18) is not satisfied. The intrinsic contribution to the inverse dephasing time, Eq. (15), dominates.

The presence of the gate does not change qualitatively the contribution to the dephasing rate due to processes intrinsic to the wire, which has the same dependence on temperature as for an isolated wire, $\tau_{\text{int}}^{-1} \propto T^{2/3}$ (see Refs. 9–11). Note, however, that the long range part of the potential of an electron in the wire is suppressed by the screening by the gate.

The main results are schematically depicted in Fig. 2. The contribution from intrinsic processes dominates both at low and high temperatures. The existence of an intermediate range of temperatures where the gate determines the inverse dephasing time depends on the inequalities in Eqs. (18), (31), and (33) or (35), depending on whether the gate is a two-dimensional metal, a one-dimensional wire, a three-dimensional metal, or a granular metal. Note that the effective dimensionality of the gate depends on the value of the distance z to the wire where dephasing is studied. A planar gate whose width w_G is smaller than z will behave as a two-dimensional metal, and a conducting wire whose diameter is smaller than z can be described as a one-dimensional gate.

In current experiments on metallic wires, the observed mean free path is comparable to the width of the wire.[1–4] In this situation it is unlikely the regime where the dephasing is determined by the gate can be observed. Note, however, that the screening of the electrostatic interactions within the wire and between the wire and the gate tends to enhance the relative contribution from the gate.

On the other hand, in multiwall carbon nanotubes Fabry-Perot interferences have been found in samples of lengths ≈ 500 nm,[25] so that the elastic mean free path is, at least, comparable to this length. The number of channels in these nanotubes, $k_F^w w_w$ is not large. Using Eq. (18), we find that a quasi-two-dimensional metallic gate can reduce significantly the observed coherence effects[25] if it is sufficiently dirty (so that $k_F^G l_G$ is not too large) and it is at distances of order 10^3 Å. A similar situation may arise in conducting channels made from doped semiconductors,[26,27] where high mobilities and small lateral dimensions have been achieved. In these systems, the elastic mean free path can exceed 1 μm, and the lateral dimensions of the wire are comparable to the Fermi wavelength. Using Eq. (18) we find that a metallic gate with a short mean free path can contribute significantly to the dephasing in the wire when its distance to the wire is one micron or less.

It is also worth noting that a nonperturbative treatment shows that the coupling to an environment modelled by the Caldeira-Leggett model strongly suppresses quantum coherence in a variety of situations.[14–17]

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[1] P. Mohanty, E. M. Q. Jariwala, and R. A. Webb, Phys. Rev. Lett. **78**, 3366 (1997).
[2] F. Pierre and N. O. Birge, Phys. Rev. Lett. **89**, 206804 (2002).
[3] P. Mohanty and R. A. Webb, Phys. Rev. Lett. **91**, 066604 (2003).
[4] F. Pierre, A. B. Gougam, A. Anthore, H. Pothier, D. Esteve, and N. O. Birge, Phys. Rev. B **68**, 085413 (2003).
[5] Y. Imry, J. Fukuyama, and C. Schwab, Europhys. Lett. **47**, 608 (1999).
[6] A. Zawadovski, J. von Delft, and D. C. Ralph, Phys. Rev. Lett. **83**, 2632 (1999).
[7] A. Kaminski and L. I. Glazman, Phys. Rev. Lett. **86**, 2400 (2001).
[8] D. S. Golubev and A. D. Zaikin, Phys. Rev. Lett. **81**, 1074 (1998).
[9] B. Altshuler, A. Aronov, and D. Khmelnitskii, J. Phys. C **15**, 7367 (1982).
[10] A. Stern, Y. Aharonov, and Y. Imry, Phys. Rev. A **41**, 3436 (1990).
[11] G. Montambaux and E. Akkermans, cond-mat/0404361 (2004).
[12] F. Guinea, R. A. Jalabert, and F. Sols, Phys. Rev. B **70**, 085310 (2004).
[13] A. O. Caldeira and A. J. Leggett, Phys. Rev. Lett. **46**, 211 (1981).
[14] D. S. Golubev and A. D. Zaikin, Physica B **255**, 164 (1998).
[15] G.-L. Ingold, R. A. Jalabert, and K. Richter, Am. J. Phys. **69**, 201 (2001).
[16] F. Guinea, Phys. Rev. B **65**, 205317 (2002).
[17] F. Guinea, Phys. Rev. B **67**, 045103 (2003).
[18] This screening effect, however, only leads to logarithmic corrections in the standard expression for the inverse dephasing time in wires (Refs. 9–11).
[19] Y. M. Blanter, Phys. Rev. B **54**, 12 807 (1996).
[20] R. H. Ritchie and A. L. Marusak, Surf. Sci. **4**, 234 (1966).
[21] B. Mehlig and M. Wilkinson, J. Phys.: Condens. Matter **9**, 3277 (1997).
[22] E. J. Austin and M. Wilkinson, J. Phys.: Condens. Matter **5**, 8461 (1993).
[23] W. P. Halperin, Rev. Mod. Phys. **58**, 533 (1986).
[24] It is interesting to note that, for a ballistic system, $\langle[\mathbf{r}(\tau) - \mathbf{r}(\tau')]^2\rangle \approx v_F^2(\tau-\tau')^2$. This result implies that $\tau_G^{-1} \propto T^{1/3}$ in this case.
[25] W. Liang, M. Bockrath, D. Bozovic, J. H. Hafner, M. Tinkham, and H. Park, Nature (London) **411**, 665 (2001).
[26] A. Yacoby, H. L. Störmer, K. W. Baldwin, L. N. Pfeiffer, and K. W. West, Solid State Commun. **101**, 77 (1997).
[27] O. M. Auslaender, A. Yacoby, R. de Picciotto, K. W. Baldwin, L. N. Pfeiffer, and K. W. West, Science **295**, 825 (2002).