

Decoherence of Fermions Subject to a Quantum Bath

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Abstract. The destruction of quantum-mechanical phase coherence by a fluctuating quantum bath has been investigated mostly for a single particle. However, for electronic transport through disordered samples and mesoscopic interference setups, we have to treat a many-fermion system subject to a quantum bath. Here, we review a novel technique for treating this situation in the case of ballistic interferometers, and discuss its application to the electronic Mach–Zehnder setup. We use the results to bring out the main features of decoherence in a many-fermion system and briefly discuss the same ideas in the context of weak localization.

1 Introduction

There are two main messages of this brief review. First, regarding the physics of decoherence, we will argue that decoherence processes depend strongly on the type of system (single particle vs. many particles) and the type of noise (classical vs. quantum). Electronic interference experiments at low temperatures require a treatment of a many-fermion system coupled to a quantum bath. In that case, true many-body features come into play. This includes, in particular, the influence of Pauli blocking (that tends to restrain decoherence), and the fact that both hole- and particle-scattering processes contribute equally to the full decoherence rate. Second, regarding theoretical methods, we review a novel technique for treating ballistic interferometers subject to a quantum bath, which is based on the ideas behind the quantum Langevin equation (as it is known for the Caldeira–Leggett model). This is more efficient than generic methods (like Keldysh diagrams), and we will discuss the physical meaning of its ingredients. We apply it to the interference contrast and the current noise in an electronic Mach–Zehnder interferometer. We will also mention how the same ideas (if not the same technical methods) help to understand decoherence in weak localization within a path-integral framework.

The reasons for studying decoherence range from fundamental aspects of quantum mechanics to possible applications. On the fundamental level, the

transition from the quantum world (with interference effects) to the classical world is due to the unavoidable fluctuations of the environment that tend to destroy macroscopic superpositions very rapidly (see, e.g., [1, 2]). These issues have been studied mostly with simplified single-particle models of decoherence.

In solid-state physics, decoherence first was investigated in the field of spin resonance, where straightforward Markoff master equation treatments are often sufficient [3, 4]. The quantum dissipative two-level system (spin-boson model) was studied in more detail at the beginning of the eighties [5]. During the preceding decade, interest in these questions has seen a revival due to the prospect of quantum information applications [6], where the decoherence time has to be at least ten thousand times longer than the time of elementary operations in order for error correction to work.

Regarding electronic transport phenomena, which will be our focus in the following, decoherence effects became important for the first time during the study of interference effects in disordered conductors, such as universal conductance fluctuations and weak localization (for a review see, e.g., [7, 10]). Later on, man-made interference structures were produced in metals and semiconductors, including Aharonov–Bohm rings, double quantum dot interferometers, and (most recently) Mach–Zehnder interferometers. These setups are also important in the quantum information context, both for generating, transporting, or detecting entanglement, and as highly sensitive measurement devices. The main nontrivial dependence of the interference contrast on temperature or transport voltage is produced by decoherence.

2 Single Particle Decoherence

Let us look at a single particle traversing a two-way interferometer (Fig. 1). Its wavepacket has been split into two packets $\psi_{L/R}$, going along the two arms (left/right). After these packets recombine, they form an interference pattern. This consists of a classical part (sum of probabilities) and an interference term, which is sensitive to a relative phase:

$$|\psi(x)|^2 = |\psi_L(x)|^2 + |\psi_R(x)|^2 + \psi_L^*(x)\psi_R(x)e^{i\varphi} + c.c. . \quad (1)$$

What happens once the particle is subjected to (classical) noise, i.e., a fluctuating potential $V(x, t)$? Even before acceleration/deceleration effects are noticeable, a random relative phase φ between the two paths is introduced. The actual pattern is obtained by averaging $|\psi|^2$ over many experimental runs. Since $|\langle e^{i\varphi} \rangle| \leq 1$, the interference term is suppressed.

At low temperatures ($k_B T < \hbar\omega$) we have to consider a quantum bath, for which there exists an alternative description of decoherence: The bath acts as a kind of which-way detector, with its initial state evolving towards either one of two states, $|\chi_R\rangle$ or $|\chi_L\rangle$, depending on the path of the particle [8–10].

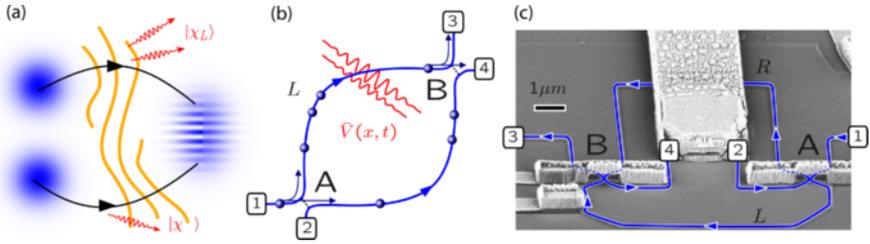


Fig. 1. (a) Decoherence in a two-way interferometer. The fluctuations (*wiggly lines*) introduce a random relative phase between the two paths. The quantum bath is left behind in two different states. This blurs the interference pattern (*right part of pattern*). (b) Schematic Mach-Zehnder setup. (c) SEM picture of electronic QHE Mach-Zehnder at the Weizmann institute (courtesy of I. Neder and M. Heiblum)

Now it is the overlap of these bath states that determines the suppression of the interference term. That overlap is nothing but the Feynman–Vernon influence functional [11].

Even if the particle is coupled to a quantum bath, decoherence may still be described using a classical noise spectrum, if the particle’s energy is high and its motion is semiclassical (Fig. 2). To understand this, consider a simple weak coupling situation, where the total decoherence rate is given by the sum of downward and upward scattering rates, calculated using Fermi’s Golden Rule. These rates can be related to the spectrum $\langle \hat{V}\hat{V} \rangle_{\omega} \equiv \int dt e^{i\omega t} \langle \hat{V}(t)\hat{V}(0) \rangle$ of the quantum noise potential \hat{V} . We have $\Gamma_{\downarrow} \propto \langle \hat{V}\hat{V} \rangle_{\omega} \propto n(\omega) + 1$ and $\Gamma_{\uparrow} \propto \langle \hat{V}\hat{V} \rangle_{-\omega} \propto n(\omega)$, where ω is the frequency transfer and $n(\omega)$ the thermal occupation. Obviously, the sum of these rates does not change if we replace \hat{V} by classical noise with a symmetrized correlator $\langle VV \rangle_{\omega} = \left(\langle \hat{V}\hat{V} \rangle_{\omega} + \langle \hat{V}\hat{V} \rangle_{-\omega} \right) / 2$ (red curve in Fig. 2b). This can also be seen in a more general treatment, using a semiclassical evaluation of the Feynman–Vernon influence functional. We note that such a replacement is impermissible near the ground state of the system, where downward transitions are blocked.

3 Many Particles

Up to now, we have considered a single particle subject to classical or quantum noise. This has been the mainstay of research in quantum dissipative systems for a long time, with paradigmatic models such as the spin-boson model or the Caldeira–Leggett model of a single particle coupled to a bath of harmonic oscillators. However, in electronic transport experiments (and

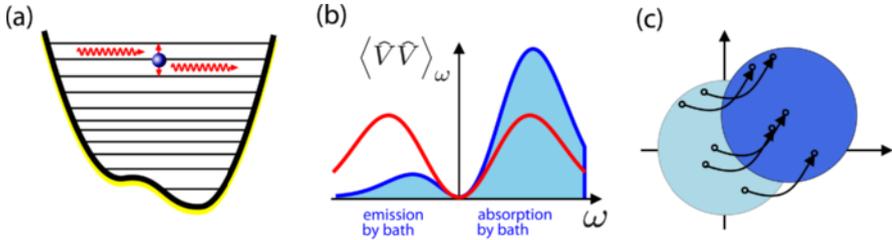


Fig. 2. (a) The scattering rates $\Gamma_{\downarrow/\uparrow}$ of a highly excited particle are related to the quantum noise spectrum (b). (c) A many-fermion system subject to classical noise

other setups, e.g., cold atom BEC interferometers) we are invariably dealing with a many-particle system. What are the new features arising in that case?

Everything remains straightforward if the noise is classical. Then, the many-particle problem reduces to the single-particle case: The wave function of each particle evolves according to the single-particle Schrödinger equation with a given noise field $V(x, t)$. In the case of many fermions, all the single-particle wave functions remain orthogonal, forming a Slater determinant (in the absence of intrinsic interactions). Pauli blocking is then completely unimportant.

We now discuss the one remaining combination: a many-fermion system coupled to a quantum bath. Unlike all the previous cases, this cannot be reduced to “single particle + classical noise”: True many-body effects come into play (and appropriate methods are needed). Up to now, comparatively few quantum-dissipative many-particle systems have been studied. Examples include open Luttinger liquids [12], many-electron Aharonov–Bohm rings subject to quantum charge [13] or flux [14] noise, many-fermion generalizations of the Caldeira–Leggett model [14, 16, 17], and double quantum dot interferometers coupled to a quantum bath [15]. Here, we are going to review a recently developed general method of solution for ballistic interferometers [18, 26], and then briefly discuss the same physics in the context of disordered systems (weak localization).

4 The Mach–Zehnder Interferometer

The Mach–Zehnder (MZ) interferometer arguably represents the simplest kind of two-way interference setup (Fig. 1). Tuning the relative phase (via the magnetic flux ϕ) yields sinusoidal interference fringes in the currents at the two output ports.

Recently, this model has been realized in electronic transport experiments. The group of *Moty Heiblum* at the Weizmann institute managed to employ edge channels of the integer quantum Hall effect in a two-dimensional electron gas to build an ideal MZ setup with single-channel transport and without

backscattering [19, 20]. The group measured the decrease of visibility (interference contrast in $I(\phi)$) as a function of rising temperature and transport voltage. No complete explanation for the results has been provided up to now, especially for the oscillations in the visibility [20]. Here, we will explore the possibility that at least part of the decrease in visibility is due to decoherence processes.

The effects of classical noise $V(x, t)$ onto a MZ setup have been studied intensively: The suppression of interference contrast to lowest order in the noise correlator was first calculated in the work of *Seelig* and *Büttiker* [21]. Building on this result, we treated the model to all orders in the interaction, calculating both the interference contrast and the effects on the shot noise in the output port of the interferometer [22, 23]. The shot noise has been measured and suggested as a tool to diagnose different sources of the loss in visibility [19]. These studies have recently been extended to the full counting statistics [24] and a renewed analysis of the dephasing terminal model [25]. However, as pointed out above, the situation is more involved for quantum noise, which is needed to account for the loss of visibility with rising bias voltage.

5 Equations of Motion Approach to Decoherence in Ballistic Interferometers

Recently, we have introduced a novel equations of motion technique for a many-particle system subject to a quantum bath [18], inside a ballistic interferometer (a detailed discussion may be found in [26]). It is similar in spirit to the quantum Langevin equation that can be employed to solve the Caldeira–Leggett model [4, 27]. Briefly, the idea of the latter is the following (when formulated on the level of Heisenberg equations). The total quantum force \hat{F} acting on the given particle, due to the bath particles, can be decomposed into two parts:

$$\hat{F}(t) = \hat{F}_{(0)}(t) + \int_{-\infty}^t D^R(t-t')\hat{x}(t')dt'. \quad (2)$$

The first describes the intrinsic fluctuations. It derives from the solution to the free equations of motion of the bath oscillators, with thermal and quantum (zero-point) fluctuations due to the stochastic initial conditions. The second part of the force is due to the response of the bath to the particle’s motion. We will call it the “back-action” term, and it gives rise to features such as mass renormalization and friction. Equation (2) is valid on the operator level (not only for averages). In this way, one has “integrated out” the bath by solving for its motion. Plugging the force \hat{F} into the right-hand-side (rhs) of the Heisenberg equation of motion for \hat{x} yields the quantum Langevin equation, which in practice can only be solved for a free particle or a harmonic oscillator (linear equations).

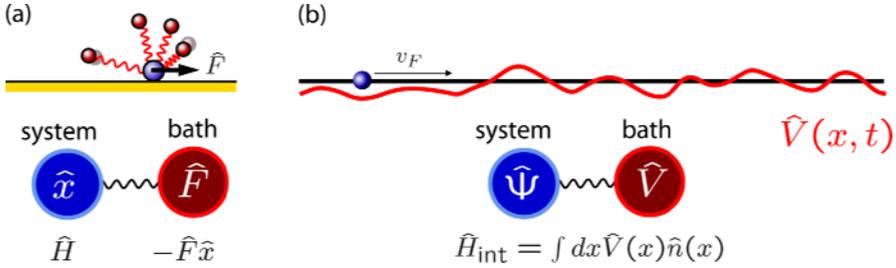


Fig. 3. (a) The Caldeira–Leggett model (single particle and oscillator bath) and (b) a ballistic many-particle system subject to a quantum noise potential $\hat{V}(x, t)$

In the case of a many-particle system, it is the density $\hat{n}(x) = \hat{\psi}^\dagger(x)\hat{\psi}(x)$ that couples to a scalar noise potential $\hat{V}(x)$. The place of \hat{x} and \hat{F} in the quantum Langevin equation for a single particle is thus taken by the particle field $\hat{\psi}$ and \hat{V} , respectively. Let us now specialize to the case of fermions traveling ballistically inside the arm of an interferometer. We will assume chiral motion and use a linearized dispersion relation, as this is sufficient to describe decoherence. Then the fermion field obeys the following equation (with a slight approximation [18, 26]; we set $\hbar = 1$):

$$i(\partial_t + v_F \partial_x) \hat{\psi}(x, t) = \hat{V}(x, t) \hat{\psi}(x, t). \quad (3)$$

The formal solution of this equation is straightforward and analogous to the version for classical noise $V(x, t)$. The particle picks up a fluctuating “quantum phase” inside a time-ordered exponential:

$$\begin{aligned} \hat{\psi}(x, t) &= \hat{T} \exp \left[-i \int_{t_0}^t dt_1 \hat{V}(x - v_F(t - t_1), t_1) \right] \\ &\times \hat{\psi}(x - v_F(t - t_0), t_0). \end{aligned} \quad (4)$$

In contrast to the case of classical noise, the field \hat{V} contains the response to the fermion density, in addition to the intrinsic fluctuations $\hat{V}_{(0)}$:

$$\hat{V}(x, t) = \hat{V}_{(0)}(x, t) + \int_{-\infty}^t dt' D^R(x, t, x', t') \hat{n}(x', t'). \quad (5)$$

Here D^R is the unperturbed retarded bath Green’s function, $D^R(1, 2) \equiv -i\theta(t_1 - t_2) \langle [\hat{V}(1), \hat{V}(2)] \rangle$, where \hat{V} -correlators refer to the free field. With these two equations, it becomes possible to calculate correlators of the fermion field (such as current and shot noise).

6 Decoherence Rate in a Many-Fermion System

Employing the formal solution from above (and using a lowest-order Markoff approximation [18, 26]), we find that the contribution of each electron to the interference term in the current is multiplied by a factor

$$1 - \Gamma_\varphi(\epsilon)\tau + i\delta\bar{\varphi}(\epsilon), \quad (6)$$

with a phase shift $\delta\bar{\varphi} \propto \tau$. We focus on the suppression brought about by a decoherence rate $\Gamma_\varphi(\epsilon)$ that depends on the energy $\epsilon(k)$ of the incoming electron:

$$\Gamma_\varphi(\epsilon) = \int_0^\infty \frac{d\omega}{v_F} \text{DOS}_q(\omega) \left[\underbrace{2n(\omega) + 1}_{\text{thermal \& zeropoint fluctuations}} - \underbrace{(\bar{f}(\epsilon - \omega) - \bar{f}(\epsilon + \omega))}_{\substack{\text{from "back - action"} \\ \Rightarrow \text{Pauli blocking}}} \right] \quad (7)$$

The rate is an integral over all possible energy transfers ω . They are weighted by the bath spectral “density of states” $\text{DOS}_q(\omega) = -\text{Im}D_q^R(\omega)/\pi$, where $q = \omega/v_F$ for ballistic motion. The first term in brackets stems from the $\hat{V}_{(0)}$ in the quantum phase. By itself, this would give rise to an energy-independent rate and a visibility independent of bias voltage (in contrast to experimental results [19, 20]). Thus, the second term is crucially important: It contains the average nonequilibrium distribution $\bar{f} = (f_L + f_R)/2$ inside the arms (for equal coupling to both arms) and implements the physics of Pauli blocking. At $T = 0$, it suppresses all transitions that would take the electron into an occupied state (when $\bar{f}(\epsilon - \omega) = 1$ and this cancels against the 1 from the zero-point fluctuations).

The other main difference (vs. the case of a single particle) is less obvious but equally important. The decoherence rate Γ_φ is not simply given by the particle-scattering rate, but contains a contribution from hole scattering processes, where a particle at another energy $\epsilon + \omega$ is scattered into the given state at ϵ (with a factor $f(\epsilon + \omega)$ associated). This is a generic feature for decoherence of fermionic systems coupled to a quantum bath, and we now discuss the physical reason. In a single-particle language, the first beam splitter creates a superposition of the form $t|R\rangle + r|L\rangle$, with t/r transmission and reflection amplitudes and R/L a packet inside the right/left arm. In the presence of a sea of other fermions, we should write instead a superposition of many-body states, for example:

$$t|1, 1, 0, \underline{0}, 0, 0; 1, 1, 1, \underline{1}, 1, 0\rangle + r|1, 1, 0, \underline{1}, 0, 0; 1, 1, 1, \underline{0}, 1, 0\rangle. \quad (8)$$

Here the occupations $|\text{left}; \text{right}\rangle$ of single-particle states in both arms are indicated, with a bar denoting the energy level ϵ of interest and the remaining particles (in the nonequilibrium distributions) playing the role of spectators. The interference is sensitive to the coherence $t|\dots, \underline{0}, \dots; \dots, \underline{1}, \dots\rangle +$

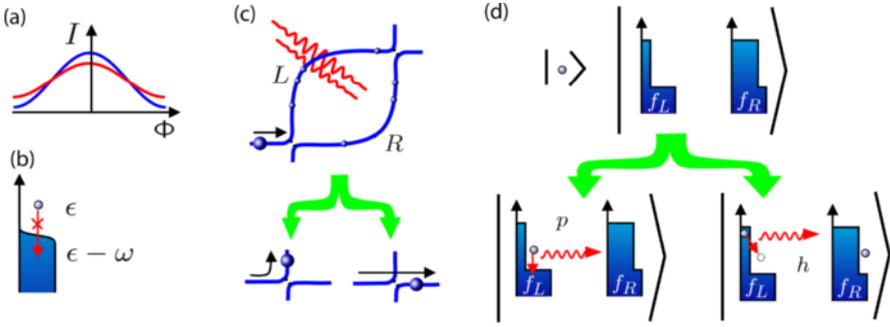


Fig. 4. (a) Decoherence suppresses the visibility of the interference pattern (*red curve (online color)*). (b) Pauli blocking restrains decoherence. (c) A particle arriving at the first beam splitter turns into a coherent superposition $t|R\rangle + r|L\rangle$. (d) For a many-fermion system, both particle- and hole-scattering contribute to decoherence: $\Gamma_\varphi = (\Gamma_p + \Gamma_h)/2$

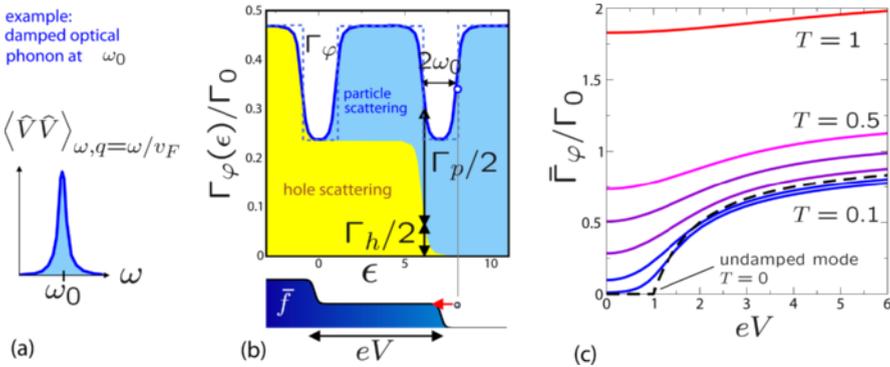


Fig. 5. The decoherence rate Γ_φ for the illustrative example of an optical phonon mode (a), as a function of energy of the incoming electron (b), and (c) the energy-averaged rate $\bar{\Gamma}_\varphi$ as a function of voltage V and temperature T (in units of ω_0). *Dashed curves in (b) and (c) refer to an ideal undamped mode at $T = 0$*

$r|\dots, \underline{1}, \dots; \dots, \underline{0}, \dots\rangle$ that requires not only the presence of a particle in one arm but also the absence of a particle in the other respective arm.

This is why the many-body superposition can equally be destroyed by particle- and hole-scattering (leading to states with $|\dots, \underline{0}, \dots; \dots, \underline{0}, \dots\rangle$ or $|\dots, \underline{1}, \dots; \dots, \underline{1}, \dots\rangle$, respectively).

We have illustrated this in Fig. 5, where we have chosen a simple model bath spectrum (a broadened optical phonon mode at ω_0). Let us focus on small temperatures $T \ll \omega_0$. If the electron is far above the Fermi sea, it

can easily undergo spontaneous emission and lose its coherence, thus the decoherence rate is maximal. For smaller energies, near the upper step in \bar{f} , it might end up in an occupied state, thus Γ_p is reduced, leading to a dip in Γ_φ . Inside the transport voltage window, Γ_p remains at 1/2 its previous value, but now Γ_h raises Γ_φ back to its maximal value. Finally, another dip is observed near the lower edge of the voltage window. The visibility is directly given by $1 - \bar{\Gamma}_\varphi$, with $\bar{\Gamma}_\varphi$ the energy-average of Γ_φ over the voltage window. At $T, V \rightarrow 0$, $\bar{\Gamma}_\varphi$ vanishes. Decoherence sets in only when the electron can emit phonons and thereby reveal its path through the MZ setup. At higher temperatures, the Fermi distribution becomes smeared, thereby easing the restrictions of Pauli blocking, and the thermal fluctuations of the bath grow, increasing Γ_φ . For $T \gg \omega_0$, the energy/voltage-dependence of Γ_φ becomes unimportant, and an approximate treatment becomes possible, replacing the quantum bath by classical noise.

7 Decoherence in Weak Localization

We now briefly discuss how the same concepts apply to weak localization [7, 29–31], where the constructive interference of time-reversed pairs of diffusive trajectories increases the electrical resistance of a disordered sample. One is interested in the linear response conductance, where the external perturbation (the electric field) induces a particle-hole excitation by lifting one of the particles above the Fermi sea. This creates a many-body state similar to the one above, $\sqrt{1 - \delta^2} |\dots, \underline{1}, \dots, \underline{0}, \dots\rangle + \delta |\dots, \underline{0}, \dots, \underline{1}, \dots\rangle$, where δ is the small amplitude of the excited state. Following arguments analogous to those above [33], we see again why both particle- and hole-scattering processes contribute to the decoherence rate.

Many discussions of decoherence in weak localization have focussed on the thermal (classical) part of the Nyquist noise [29, 30]. This leads to a single-particle problem that can be treated using path integrals [30]. Diagrammatic calculations of the decoherence rate in the presence of a quantum bath [31] yielded results that can be interpreted in the manner discussed above. It is obviously desirable, though difficult, to cast these as well into the powerful path-integral framework. *Golubev* and *Zaikin* were the first to present a formally exact influence functional approach for many-fermion systems [35]. Their semiclassical evaluation yielded a decoherence rate that does not vanish at $T = 0$ and is independent of electron energy, in contrast to diagrammatic calculations and the ideas about Pauli blocking discussed above.

Recently, we have revisited this problem [33, 34] and have shown that the results of much more complicated diagrammatic calculations [32, 34] can be exactly reproduced by a rather simple prescription. The case “many particles + quantum bath” may be reduced to “single particle + classical noise”, pro-

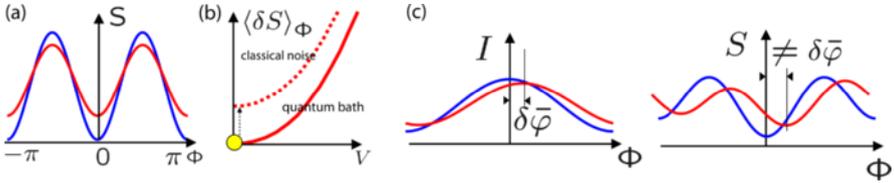


Fig. 6. Main effects of a quantum bath on shot noise in a MZ setup: (a) Suppression of visibility in $S(\phi)$. (b) No Nyquist noise correction, but classical conductance fluctuations $S \propto V^2$ at high voltages. (c) Different phase shifts in $I(\phi)$ and $S(\phi)$ for asymmetric setups

vided one uses an effective, modified noise spectrum of the following form [33]:

$$\langle \hat{V} \hat{V} \rangle_\omega \mapsto \langle VV \rangle_\omega^{\text{eff}} \equiv \frac{1}{2} \langle \{ \hat{V}, \hat{V} \} \rangle_\omega + \frac{1}{2} \langle [\hat{V}, \hat{V}] \rangle_\omega (f(\epsilon + \omega) - f(\epsilon - \omega)) \tag{9}$$

The first part is the symmetrized quantum correlator, containing the zero-point fluctuations. The second part incorporates Pauli blocking. These terms correspond to those in the equation of motion approach (7), with which this method is consistent. The resulting decoherence rate vanishes at $T = 0$ [35, 36]. Earlier similar ideas [30, 37, 38] represent approximations to the present approach.

8 Effects of a Quantum Bath on Shot Noise

We briefly return to the MZ setup. Using our approach, it is possible to discuss the influence of the quantum bath on the shot noise [28] power S in the output port [18, 26]. The visibility of the interference pattern $S(\phi)$ is reduced, although this cannot be described by the same decoherence rate as for the current. The most important feature refers to the phase shifts observed in asymmetric setups, for the current and the shot noise. These can become different: $I(\phi) = \bar{I} + \delta I \cos(\phi - \delta\bar{\phi})$ and $S(\phi) = \bar{S} + S_1 \cos(\phi - \delta\phi_1) + S_2 \cos(2(\phi - \delta\phi_2))$, with $\delta\bar{\phi} \neq \delta\phi_1 \neq \delta\phi_2$ in general. This prediction is in contrast to all simpler models (involving classical noise etc.), which usually do not give rise to phase shifts at all. Something like this seems to have been observed in recent experiments at the Weizmann institute. Another equally important avenue of current research is the application to nonlinear (non-Gaussian) environments.

9 Conclusions

Decoherence in transport interference situations can often be reduced to the case of a single particle subject to classical noise. However, for a many-fermion system subject to a quantum bath, true many body features remain and have to be taken into account via suitable techniques. The two main physical features are Pauli blocking and the importance of both hole- and particle-scattering processes. The technical innovation reviewed here is an equation of motion approach that is well suited to describe decoherence of many particles moving in ballistic interferometers. We have discussed its application to the MZ interferometer setup, the loss of visibility in the current and (briefly) the effects on shot noise. In addition, we have pointed out that the same kind of physics applies to decoherence in weak localization.

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References

- [1] E. Joos et al.: *Decoherence and the Appearance of a Classical World in Quantum Theory* (Springer, Heidelberg, 2003)
- [2] W. H. Zurek: Decoherence and the transition from quantum to classical, *Phys. Today* **44**, 36 (1991)
- [3] K. Blum, *Density Matrix Theory and Applications* (Plenum, New York, 1996)
- [4] U. Weiss: *Quantum Dissipative Systems* (World Scientific, Singapore, 2000)
- [5] A. J. Leggett, S. Chakravarty, A. T. Dorsey, M. P. A. Fisher, A. Garg, W. Zwerger: Dynamics of the dissipative two-state system, *Rev. Mod. Phys.* **59**, 1-85 (1987)
- [6] C. H. Bennett, D. P. DiVincenzo: Quantum information and computation, *Nature* **404**, 247-255 (2000)
- [7] C. W. J. Beenakker, H. van Houten: Quantum transport in semiconductor nanostructures, *Solid State Phys.* **44**, 1-228 (1991)
- [8] A. Stern, Y. Aharonov, Y. Imry: Phase uncertainty and loss of interference: A general picture, *Phys. Rev. A* **41**, 3436-3448 (1990)
- [9] D. Loss, K. Mullen: Dephasing by a dynamic asymmetric environment, *Phys. Rev. B* **43**, 13252 (1991)
- [10] Y. Imry: *Introduction to Mesoscopic Physics*, 2nd ed. (Oxford University Press, 2002)
- [11] R. P. Feynman, F. L. Vernon: The theory of a general quantum mechanical system interacting with a linear dissipative system, *Ann. Phys. (NY)* **24**, 118-173 (1963)
- [12] A. H. C. Neto, C. D. Chamon, C. Nayak: Open Luttinger liquids, *Phys. Rev. Lett.* **79**, 4629 (1997)

- [13] P. Cedraschi, V. V. Ponomarenko, M. Büttiker: Zero-point fluctuations and the quenching of the persistent current in normal metal rings, *Phys. Rev. Lett.* **84**, 346-349 (2000)
- [14] F. Marquardt, C. Bruder: Aharonov–Bohm ring with fluctuating flux, *Phys. Rev. B* **65**, 125315 (2002)
- [15] F. Marquardt, C. Bruder: Dephasing in sequential tunneling through a double-dot interferometer, *Phys. Rev. B* **68**, 195305 (2003)
- [16] F. Marquardt, D. S. Golubev: Relaxation and dephasing in a many-fermion generalization of the Caldeira–Leggett model, *Phys. Rev. Lett.* **93**, 130404 (2004)
- [17] F. Marquardt, D. S. Golubev: Many-fermion generalization of the Caldeira–Leggett model, *Phys. Rev. A* **72**, 022113 (2005)
- [18] F. Marquardt: Fermionic Mach–Zehnder interferometer subject to a quantum bath, *Europhys. Lett.* **72**, 788 (2005)
- [19] Y. Ji, Y. Chung, D. Sprinzak, M. Heiblum, D. Mahalu, H. Shtrikman: An electronic Mach–Zehnder interferometer, *Nature* **422**, 415 (2003)
- [20] I. Neder, M. Heiblum, Y. Levinson, D. Mahalu, V. Umansky: Unexpected behavior in a two-path electron interferometer, *Phys. Rev. Lett.* **96**, 016804 (2006)
- [21] G. Seelig, M. Büttiker: Charge-fluctuation-induced dephasing in a gated mesoscopic interferometer, *Phys. Rev. B* **64**, 245313 (2001)
- [22] F. Marquardt, C. Bruder: Influence of dephasing on shot noise in an electronic Mach–Zehnder interferometer, *Phys. Rev. Lett.* **92**, 056805 (2004)
- [23] F. Marquardt, C. Bruder: Effects of dephasing on shot noise in an electronic Mach–Zehnder interferometer, *Phys. Rev. B* **70**, 125305 (2004)
- [24] H. Förster, S. Pilgram, M. Büttiker: Decoherence and full counting statistics in a Mach–Zehnder interferometer, *Phys. Rev. B* **72**, 075301 (2005)
- [25] S. Pilgram, P. Samuelsson, H. Förster, M. Büttiker: Full counting statistics for voltage and dephasing probes, *cond-mat/0512276*
- [26] F. Marquardt: Equations of motion approach to decoherence and current noise in ballistic interferometers coupled to a quantum bath, *cond-mat/0604458* (2006)
- [27] A. O. Caldeira, A. J. Leggett: Path integral approach to quantum Brownian motion, *Physica* **121A**, 587 (1983)
- [28] Ya. M. Blanter, M. Büttiker: Shot noise in mesoscopic conductors, *Phys. Rep.* **336**, 1 (2000)
- [29] B. L. Altshuler, A. G. Aronov, D. E. Khmel'nitsky: Effects of electron–electron collisions with small energy transfers on quantum localization, *J. Phys. C Solid State* **15**, 7367 (1982)
- [30] S. Chakravarty, A. Schmid: Weak localization: The quasiclassical theory of electrons in a random potential, *Phys. Rep.* **140**, 195 (1986)
- [31] H. Fukuyama, E. Abrahams: Inelastic scattering time in two-dimensional disordered metals, *Phys. Rev. B* **27**, 5976 (1983)
- [32] I. Aleiner, B. L. Altshuler, M. E. Gershenson: Interaction effects and phase relaxation in disordered systems, *Waves in Random Media* **9**, 201 (1999) [*cond-mat/9808053*]
- [33] F. Marquardt, J. v. Delft, R. Smith, V. Ambegaokar: Decoherence in weak localization I: Pauli principle in influence functional, *cond-mat/0510556*

- [34] J. v. Delft, F. Marquardt, R. Smith, V. Ambegaokar: Decoherence in weak localization II: Bethe–Salpeter calculation of Cooperon, cond-mat/0510557; J. v. Delft: Influence functional for decoherence of interacting electrons in disordered conductors, cond-mat/0510563
- [35] D. S. Golubev, A. D. Zaikin: Quantum decoherence and weak localization at low temperatures, Phys. Rev. B **59**, 9195 (1999)
- [36] In contrast, the decoherence rate given in Ref. [35] is equal to that obtained for a single particle, with $f = 0$ in our formula, thus saturating at $T = 0$. We should note that the authors of Ref. [35] disagree with our conclusions, see their comment: D. S. Golubev and A. D. Zaikin, cond-mat/0512411.
- [37] W. Eiler: Electron–electron interaction and weak localization, J. Low Temp. Phys. **56**, 481 (1984)
- [38] D. Cohen, Y. Imry: Dephasing at low temperatures, Phys. Rev. B **59**, 11143 (1999)