

Nonequilibrium transport through an interacting monitored quantum dot

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We study the interplay between strong correlations and Markovian dephasing, resulting from monitoring the charge or spin degrees of freedom of a quantum dot described by a dissipative Anderson impurity model. Using the auxiliary master equation approach, we compute the steady-state spectral function and occupation of the dot and discuss the role of dephasing on Kondo physics. Furthermore, we consider a two-lead setup that allows to compute the steady-state current and conductance. We show that the Kondo steady state is robust to moderate charge dephasing but not to spin dephasing, which we interpret in terms of dephasing-induced heating of low-energy excitations. Finally, we show universal scaling collapse of the nonlinear conductance with a dephasing-dependent Kondo scale.

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I. INTRODUCTION

Traditionally considered as detrimental for quantum coherence, dissipation due to coupling to an external environment has recently emerged as a powerful tool to control and prepare quantum systems [1,2]. For example, tailored dissipative processes are currently used to extend the lifetime of superconducting qubits [3,4]. Similarly, continuous monitoring and quantum measurements offer a nonunitary way to prepare and manipulate quantum systems [5,6]. These breakthroughs in the field of quantum information and computation have raised interest on the possibility of manipulating many-body quantum systems via dissipative processes [7–12].

Mesoscopic systems provide a natural setting to explore these ideas. Quantum transport through correlated quantum dots has been the source of exciting many-body physics over the past decades [13,14]. The Kondo effect [15–17] and its different incarnation and generalization in and out of equilibrium, for example, provide a test bed for theories of strongly correlated electron systems [18]. Coupling quantum dots with quantum point contacts acting as charge detectors have been studied in the past both theoretically and experimentally [19–27]. More recently, controlled dissipative environments based on hybrid circuits including microwave resonators have been coupled to quantum dots [28–32], allowing to track down the fluctuations of the monitoring process [26,33,34]. On a different experimental front, transport via quantum impurities with ultracold atoms has also made tremendous progress in recent years [35–39]. These settings naturally include fast dissipative processes such as dephasing due to stimulated emission [40,41] or correlated particle losses [42–45].

Theoretical investigations on dissipative quantum impurity models, where localized levels coupled to a bath are exposed to fast Markovian dissipation, have been considered in recent years, both within a Lindblad description [46–58] or in the so-called no-click limit corresponding to a full postselection of quantum trajectories and an effective non-Hermitian evolution [59–61]. The fate of Kondo physics in the presence of local dissipation arising from continuous monitoring or dephasing has been discussed using perturbative approaches such as weak [49] or strong coupling methods [62]. In Ref. [62], it has been shown that the competition between electronic correlation and local dissipation can give rise to Kondo-Zeno crossover in the dynamics. A crucial question concerns the interplay between Kondo physics and dissipation in the steady-state transport, which is of direct experimental relevance. Transport through a noninteracting monitored dot has been recently shown to contain rich physics [51], and it is therefore natural to investigate the interacting case. At a more fundamental level, one could ask what is the effect of dephasing and monitoring on Kondo temperature and whether the former acts as an infrared cutoff like temperature or bias voltage for the conventional (unitary) Kondo problem [63].

In this article, we study the steady-state transport via an interacting monitored quantum dot, described by a dissipative Anderson impurity model (AIM) with charge or spin dephasing. Using the auxiliary master equation approach (AMEA) [64–66], we compute the steady-state dot Green's function and from these obtain linear and nonlinear conductance as a function of the monitoring rate and the interaction. We show that weak charge dephasing leaves visible signatures of Kondo physics, both in the spectral function and in the conductance, as opposed to acting on the spin sector that rapidly destroys many-body coherent effects. We interpret this result in terms of effective heating of low-energy degrees of freedom that is enhanced when the spin rather than the charge of the dot is monitored. Finally, we show that the nonlinear conductance shows a perfect scaling collapse with a dephasing-dependent

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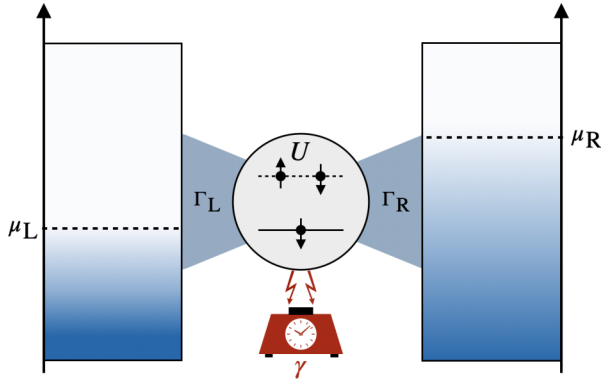


FIG. 1. Sketch of the setup for the two leads dissipative Anderson impurity model: a quantum dot subjected to Markovian dissipation due to continuous monitoring with strength γ and coupled to two large metallic leads via a hybridization $\Gamma_{L,R}$. The two leads are held at different chemical potentials, leading to a current flowing and nonequilibrium transport.

scale, suggesting that Kondo universality survives in the presence of dissipation.

II. DISSIPATIVE ANDERSON IMPURITY MODEL

We consider a model for an interacting spinful quantum dot coupled to two metallic reservoirs at different chemical potentials (see sketch in Fig. 1), as described by the Anderson impurity model [16,18] with Hamiltonian

$$\mathcal{H} = \mathcal{H}_{\text{leads}} + \mathcal{H}_{\text{hyb}} + \mathcal{H}_{\text{dot}}. \quad (1)$$

Here, the first term describes the Hamiltonian of the noninteracting leads $\mathcal{H}_{\text{leads}} = \sum_{\sigma,\mathbf{k}\alpha} \epsilon_{\mathbf{k}\alpha} c_{\mathbf{k},\sigma\alpha}^\dagger c_{\mathbf{k},\sigma\alpha}$, where $c_{\mathbf{k},\sigma\alpha}^\dagger$, $c_{\mathbf{k},\sigma\alpha}$ are the fermionic operators of the lead $\alpha = L, R$, with momentum $\mathbf{k}$ and spin σ . The leads are characterized by the single-particle energy $\epsilon_{\mathbf{k}\alpha} = \epsilon_{\mathbf{k}} - \mu_\alpha$, where the chemical potentials of the leads are set by the applied bias voltage $\mu_L - \mu_R = eV$. The second term describes the hybridization between dot and the leads $\mathcal{H}_{\text{hyb}} = \sum_{\mathbf{k},\sigma\alpha} (V_{\mathbf{k}\alpha} d_\sigma^\dagger c_{\mathbf{k},\sigma\alpha} + V_{\mathbf{k}\alpha}^* c_{\mathbf{k},\sigma\alpha}^\dagger d_\sigma)$, with coupling $V_{\mathbf{k}\alpha}$. Finally, the last term describes the dot Hamiltonian $\mathcal{H}_{\text{dot}} = \sum_\sigma (V_g - \frac{U}{2}) d_\sigma^\dagger d_\sigma + U n_\uparrow n_\downarrow$, where $d_\sigma^\dagger$ and d_σ are the dot creation/annihilation operators, V_g is the gate voltage, and U the Coulomb interaction, with $n_\sigma = d_\sigma^\dagger d_\sigma$. In the following, we consider the fermionic leads to be at finite temperature T and assume them to have a flat (constant) density of states $\Gamma_\alpha(\epsilon) = \pi \sum_{\mathbf{k}} |V_{\mathbf{k}\alpha}|^2 \delta(\epsilon - \epsilon_{\mathbf{k}\alpha}) = \Gamma_\alpha \Theta_{\text{sm}}(D - |\epsilon|)$, where D is the half-bandwidth of the lead [67]. Although simplistic, this density of states encodes the main properties of a metallic conduction lead, with a finite bandwidth and a finite weight at the Fermi level.

In addition to coherent processes on the dot described by $\mathcal{H}_{\text{dot}}$, we are interested in a situation where the impurity is exposed to local dissipative processes, which we assume to be Markovian and described by a Lindblad master equation [68]. These dissipative processes originate from some fast Markovian environment, whose microscopic degrees of freedom are not under our control and so can be traced out from the start. This has to be contrasted with the quantum bath described by the fermions $c_{\mathbf{k},\sigma\alpha}$, $c_{\mathbf{k},\sigma\alpha}^\dagger$ that play a key role in the many-body

physics of the quantum impurity. As a result of this local dissipation, the entire system (metallic leads plus quantum dot) is described by a density matrix ρ_t that evolves in time according to the Lindblad equation

$$\partial_t \rho_t = -i[H, \rho_t] + L_b \rho_t L_b^\dagger - \frac{1}{2} \{L_b^\dagger L_b, \rho_t\}, \quad (2)$$

where we have denoted by L_b , $L_b^\dagger$ the jump operators acting on the impurity system only, that is to say they are written only as functions of the operators d_σ , $d_\sigma^\dagger$, and $b = \{\text{charge, spin}\}$ denotes the operator that is being monitored. In this work, we consider the case of continuous monitoring of some Hermitian operator of the dot that leads to dissipation in the form of dephasing. In particular, we will consider separately dephasing of the total charge or spin along the z axis, i.e., jump operators of the form

$$L_{\text{charge}} = \sqrt{\gamma_{\text{charge}}} \sum_\sigma n_\sigma, \quad L_{\text{spin}} = \sqrt{\gamma_{\text{spin}}} \sum_\sigma \sigma. \quad (3)$$

We note that these jump operators commute with the total particle number; i.e., they describe energy rather than particle exchange with the Markovian bath and furthermore preserve particle-hole symmetry of the Lindblad master equation, which is therefore guaranteed for the neutrality point $V_g = U/2$. A comment on experimental feasibility is in order here: Any monitoring protocol that reduces to our Lindblad master equation upon averaging over the measurement outcomes would be enough to capture the physics discussed in this work. This includes both photocounting protocols, corresponding to quantum jump monitoring, and homodyne detection [26] as routinely done, for example, using a microwave resonator and coupling it to the charge or spin of the dot [31].

Our results are obtained using the Auxiliary Master Equation Approach [64,66,69,70] (Appendix B). In AMEA, the metallic leads are replaced by an auxiliary environment composed of a finite number N_B of bath sites and additional Markovian reservoirs. The resulting many-body Lindblad equation, which fully accounts for electronic correlations on the impurity, can be solved exactly using numerical methods. AMEA is a nonperturbative approach, and its accuracy is governed solely by the deviation between the hybridization functions of the original leads and those of the auxiliary baths. This deviation decreases exponentially with N_B , as demonstrated in Ref. [71]. In this work, we solve the auxiliary Lindblad problem using the configuration interaction scheme introduced in Ref. [66], which enables access to the Kondo scaling regime of the conductance for interaction strengths up to $U/\Gamma \sim 8$. A key advantage of AMEA is its equal applicability to both equilibrium and nonequilibrium impurity problems producing real frequency spectra. For the present case, AMEA is particularly well suited: The inclusion of the Lindblad dissipative term does not incur any additional computational cost. In Appendix B, we benchmark AMEA against the exact solution for a noninteracting dissipative impurity ($U = 0$) and find perfect agreement.

III. SPECTRAL AND DISTRIBUTION FUNCTION OF THE DOT

We start our discussion from the spectral function of the dot, which is a key quantity to characterize Kondo effect and

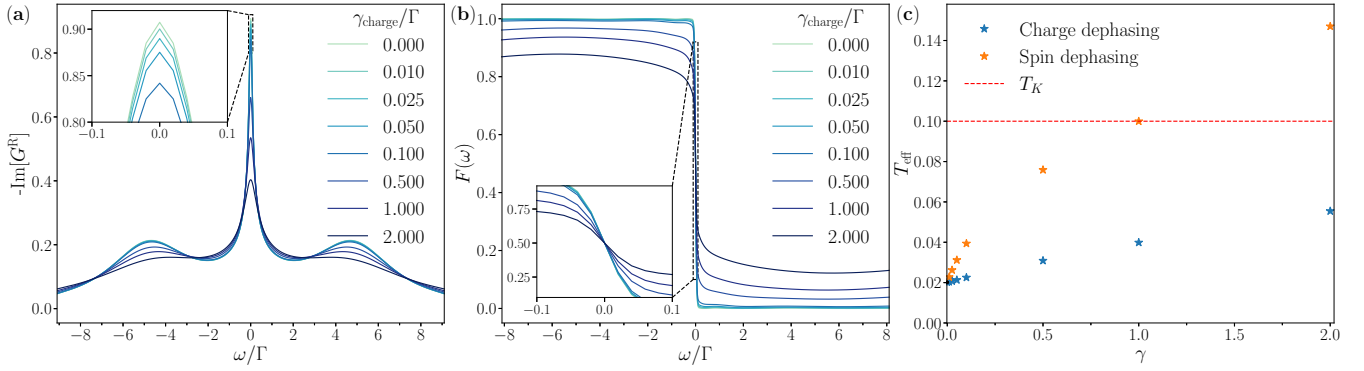


FIG. 2. Dissipative Anderson impurity model. (a) Impurity spectral function at half-filling for interaction strength $U = 8\Gamma$, $V_g = 0$, shown for increasing values of the charge dephasing rate γ_{charge} . (b) Corresponding distribution function $F(\omega)$ as γ_{charge} increases. (c) Effective temperature T_{eff} extracted from the distribution function, as a function of γ in the charge dephasing case, compared with the spin-dephasing case and the Kondo temperature of the nondissipative Anderson impurity model (taken from Ref. [66]). The leads are held at a finite temperature $T = 0.02\Gamma$ and are modeled with a flat (wideband) density of states.

it is defined as

$$A(\omega) = -\frac{1}{\pi} \text{Im}[G^R(\omega)], \quad (4)$$

where in the time domain the retarded Green's function is defined as

$$G^R(t, t') = -i\theta(t - t')\langle\{d_\sigma(t), d_\sigma^\dagger(t')\}\rangle. \quad (5)$$

We plot the spectral function in Fig. 2(a) for a fixed temperature $T = 0.02\Gamma$ and $U = 8\Gamma$, well within the Kondo regime [$T_K = 0.1\Gamma$ for these parameters; see dashed line in panel (c)] and increasing values of γ_{charge} . We see that charge dephasing reduces the height of the Kondo peak, which is expected since dissipative impurities lack the equivalent of Luttinger theorem [62], but does not broaden the resonance at least up to values of $\gamma_{\text{charge}} \sim \Gamma$, well above T_K . The incoherent excitations at high frequency, corresponding to the Hubbard bands, are suppressed already for $\gamma_{\text{charge}} < U/2$. This supports the indication that Kondo effect is robust to moderate charge dephasing. For comparison, we show in Appendix C that spin dephasing suppresses the Kondo resonance already for $\gamma_{\text{spin}} \ll T_K$, as expected for a relevant perturbation of the Kondo fixed point. To understand better the difference between the two dissipative processes, we look at the Keldysh Green's function defined as

$$G_\sigma^K(t, t') = -i\langle\{d_\sigma(t), d_\sigma^\dagger(t')\}\rangle. \quad (6)$$

By expressing the Keldysh Green's function in frequency space and relating it to the spectral function, one can write

$$G^K(\omega) = (1 - 2F(\omega))[G_\sigma^R(\omega) - G_\sigma^A(\omega)], \quad (7)$$

where $F(\omega)$ is the distribution function. In Fig. 2(b), we plot the distribution function for increasing values of γ_{charge} . We see that the main effect of charge dephasing is to suppress uniformly the occupation of finite frequency states, a signature of Markovianity, while at the same time reduce the slope of the distribution function at very low frequency (see inset). The slope allows us to extract an effective temperature T_{eff} to characterize the effect of dephasing. If the system was in true thermal equilibrium, the quantum fluctuation-dissipation theorem (FDT) would constrain the Keldysh and

retarded components to obey the relation

$$1 - 2F_{\text{eq}}(\omega) = \tanh\left(\frac{\beta\omega}{2}\right), \quad (8)$$

where $\beta = 1/T$ is the inverse temperature of the system and $\mu = 0$ in this particle-hole symmetric case. In the low-frequency limit, $\omega \ll T$, the equilibrium distribution function behaves as $1 - 2F_{\text{eq}}(\omega) \sim \omega/2T$, allowing us to extract the temperature from a linear fit.

In a nonequilibrium system on the contrary, there is no well-defined temperature and the FDT does not hold in general. Nonetheless, it is useful to use the left-hand side of the FDT relation in Eq. (8) to define an effective distribution function and extracting an effective temperature T_{eff} . We plot T_{eff} in Fig. 2(c), where we compare the effect of charge and spin dephasing. We see that T_{eff} increases with γ as expected, yet the rate is significantly smaller for the charge dephasing case, leading to $T_{\text{eff}} < T_K$ for all values of γ_{charge} considered in this work. On the other hand, the heating is parametrically stronger when dephasing acts on the spin degree of freedom, as seen from the effective temperature whose behavior at small γ_{spin} seems sublinear growing above T_K already for $\gamma_{\text{spin}} \sim \Gamma$ (see Appendix C for further data on the distribution function and the role of spin-flip noise). This can be understood since the effective temperature probes heating of low-energy degrees of freedom that are mainly spin fluctuations of the dot due to Kondo effect. The effective temperature depends also on the interaction U , particularly for the spin-dephasing case where the nonanalytic behavior at small dephasing is enhanced (see Appendix C). These results imply that monitoring of the charge degrees of freedom does not immediately destroy the Kondo regime. We now discuss the consequences of this finding for charge transport by analyzing the conductance through the dot.

IV. TRANSPORT

We now focus on the particle transport through the quantum dot both in the linear and nonlinear regimes and discuss the effect of monitoring on current and conductance. The current operator is defined as the overall flux of particles

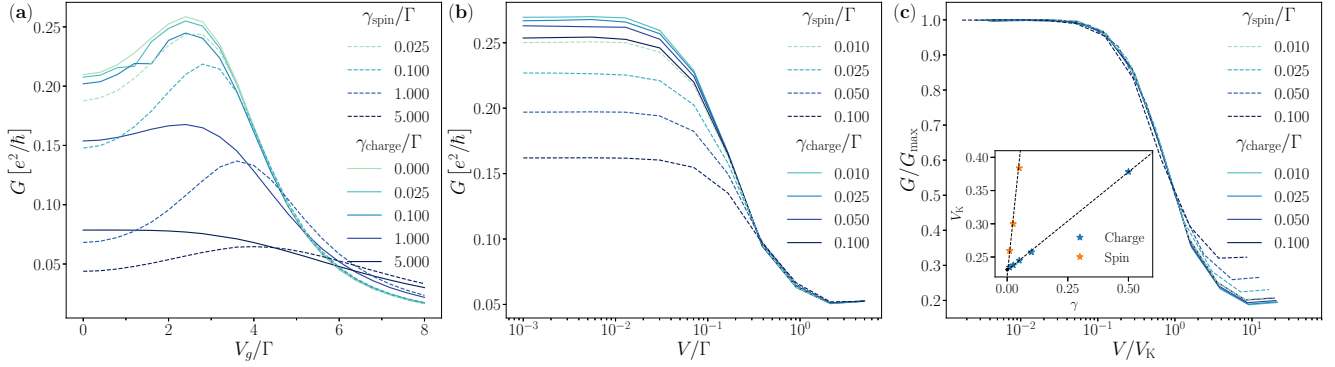


FIG. 3. Dissipative Anderson impurity model. (a) Differential conductance G as a function of the gate voltage V_g for increasing values of γ_{charge} and γ_{spin} . (b) Differential conductance G as a function of the bias voltage for increasing values of γ_{charge} and γ_{spin} . (c) The normalized conductance G/G_{max} displays a universal behavior as a function of the normalized bias V/V_K , highlighting the universality and scaling behavior across different charge and spin-dephasing rates. Here, G_{max} denotes the zero-bias conductance at a given γ , and V_K is defined as the bias voltage at which the conductance drops to half of the value it has with respect to $V = 0$ for a given γ .

flowing through the impurity at time t , namely, $\mathcal{I}(t) = \frac{1}{2}(\mathcal{I}_L(t) - \mathcal{I}_R(t))$, with $\mathcal{I}_\alpha = -e \frac{dN_\alpha}{dt}$ the instantaneous current flowing out from lead α and $N_\alpha = \sum_{\mathbf{k}\sigma} c_{\mathbf{k}\sigma,\alpha}^\dagger c_{\mathbf{k}\sigma,\alpha}$ the number of conducting electrons in lead α . The commutator can be evaluated explicitly due to the fact that the dissipation acts only at the level of the impurity, leading to $\langle \mathcal{I} \rangle_\alpha(t) = \frac{ie}{\hbar} \sum_{\mathbf{k}\sigma} V_{\mathbf{k}\alpha} \langle d_\sigma^\dagger(t) c_{\mathbf{k}\sigma,\alpha}(t) - c_{\mathbf{k}\sigma,\alpha}^\dagger(t) d_\sigma(t) \rangle$. For this reason, the expression for the current is the same as for a nondissipative dot and acquires the Meir-Wingreen form [51,72] (see Appendix A for a proof), which in the case of equal lead self-energies $\Gamma_R(\omega) = \Gamma_L(\omega) \equiv 2\Gamma(\omega)$ becomes

$$\langle \mathcal{I} \rangle = \frac{2e}{\hbar} \int \frac{d\omega}{2\pi} \Gamma(\omega) (F_R(\omega) - F_L(\omega)) \text{Im}[-G^R(\omega)]. \quad (9)$$

Here, $F_\alpha(\omega) = F_{\text{eq}}(\omega - \mu_\alpha)$ denotes the Fermi distribution of lead α with chemical potential μ_α .

We start discussing the linear conductance $G = \partial I / \partial V_{V=0}$ and its dependence on the gate voltage V_g and the dephasing rate γ . In Fig. 3(a), we plot $G(V_g)$ for different values of the charge γ_{charge} and spin γ_{spin} dephasing rates. In the standard (unitary), AIM conductance shows clear signatures of Kondo effect as well as Coulomb blockade physics. Here, we see that for $\gamma_{\text{charge}} = 0$ the conductance has a peak at $V_g \sim U/4$ and a dip at the symmetric point due to finite temperature effects. As we switch on the dephasing in either spin or charge sector, we see different behaviors: Charge dephasing does not affect transport until a value $\gamma_{\text{charge}} \sim \Gamma$ is reached, above which the suppression of coherent transport is clear. The spin dephasing, on the other hand, is very effective in reducing the conductance, particularly at $V_g \simeq 0$, leaving only a broad peak at $V_g \sim U/2$.

We then move to discuss the nonlinear transport at finite voltage. In Fig. 3(b), we plot the conductance as a function of voltage for different types of dephasing. We see that the qualitative behavior is similar; namely, dephasing reduces uniformly the low-bias conductance for $V \ll \Gamma$. Again, compared to the charge dephasing case the spin dephasing has a stronger effect on the suppression of transport. For voltages of the order $V \sim \Gamma$, all the γ dependence disappears. Quite remarkably, as we see in Fig. 3(c), we can rescale all the dephasing dependence and obtain an excellent scaling

collapse of the normalized conductance G/G_{max} . This is a remarkable result for a dissipative quantum impurity model, which suggests the existence of a universal scale V_K controlling the dephasing problem, analog to the Kondo temperature. Indeed, for the unitary AIM the conductance is known to display universal scaling with respect to T/T_K and V/T_K [66]. The dependence of the crossover scale V_K , shown in the inset of Fig. 3(c), implies that the scaling collapse survives over a broader range of parameters for charge dephasing, in agreement with our overall picture that this process does not directly affect Kondo physics.

V. CONCLUSIONS

In this work, we have studied transport through an interacting monitored quantum dot modeled as a dissipative Anderson impurity model. We have discussed in particular how dissipation due to monitoring interplays with the physics of the Kondo effect, considering the separate effects of charge and spin monitoring on both spectral and transport properties of the model. Our results highlight the robustness of Kondo physics to charge monitoring, while spin dephasing turns out to destroy rapidly any signatures of the correlated state. We have shown that this difference can be understood in terms of local heating that remains moderate when the charge of the dot is monitored, while rapidly brings the system away from the Kondo regime when acting on the spin. The robustness of Kondo physics to charge dephasing is strikingly demonstrated in the scaling collapse of the nonlinear conductance with respect to dephasing, suggesting that a form Kondo universality survives even in the presence of dissipation. An outstanding question that is motivated by our results would be to generalize Nozière's Fermi liquid picture [73] to the dissipative regime. Our work suggests a number of interesting extensions, notably to heat transport where one can expect measurements to play a role in cooling [6] and to nonreciprocal effects due to asymmetric couplings [51].

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DATA AVAILABILITY

The data that support the findings of this article are openly available in the TU Graz Repository [74].

APPENDIX A: DERIVATION OF THE MEIR-WINGREEN FORMULA IN THE PRESENCE OF DEPHASING

Due to the fact that the dephasing only acts on the impurity, the Meir-Wingreen formula for the current [72] holds in the present case as well. For completeness, we provide here a derivation closely following Ref. [72]. The current operator (from left to right) is defined as the overall flux of particles flowing through the impurity at time t , namely,

$$\mathcal{I}(t) = \frac{1}{2}[\mathcal{I}_L - \mathcal{I}_R], \quad (\text{A1})$$

where L/R refer to the left/right metallic lead and $\mathcal{I}_\alpha$ is the instantaneous current flowing out of lead α ,

$$\mathcal{I}_\alpha(t) = -e \frac{dN_\alpha}{dt}, \quad (\text{A2})$$

where $N_\alpha = \sum_{\mathbf{k}\sigma} c_{\mathbf{k},\sigma\alpha}^\dagger c_{\mathbf{k},\sigma\alpha}$ is the number of conducting electrons in lead α at time t . If one takes the trace over the density matrix evolving with the Lindblad master equation given in the main text, it follows that the average current from lead α reads

$$\langle \mathcal{I}_\alpha \rangle_t = \frac{ie}{\hbar} \sum_{\mathbf{k}\sigma} (V_{\mathbf{k}\alpha}^* \langle d_\sigma^\dagger c_{\mathbf{k},\sigma\alpha} \rangle_t - V_{\mathbf{k}\alpha} \langle c_{\mathbf{k},\sigma\alpha}^\dagger d_\sigma \rangle_t). \quad (\text{A3})$$

All operators in Eq. (A3) evolve in real time according to the Lindblad master equation. This result follows since the jump operators only act on the impurity.

In the stationary state, the expectation value of the current operator simplifies to (we now take $V_{\mathbf{k}\alpha}$ to be real for simplicity)

$$\begin{aligned} \langle \mathcal{I}_\alpha \rangle &= -i \frac{e}{\hbar} \sum_{\mathbf{k}\sigma} (V_{\mathbf{k}\alpha} \langle c_{\mathbf{k},\sigma\alpha}^\dagger d_\sigma \rangle - \text{c.c.}) \\ &= -\frac{e}{2\hbar} \sum_{\mathbf{k}\sigma} \int \frac{d\omega}{2\pi} [V_{\mathbf{k}\alpha} G_{\text{imp},\mathbf{k}\alpha}^K(\omega) + \text{c.c.}], \end{aligned} \quad (\text{A4})$$

where in the second line we have used the definition of the Green's function (which we take spin independent), $G_{\text{imp},\mathbf{k}\alpha}^K(t, t') = -i \langle [d_\sigma(t), c_{\mathbf{k},\sigma\alpha}^\dagger(t')] \rangle$ connecting the lead and impurity degrees of freedom. Since the leads are noninteracting, one can use the Dyson equation with Langreth rules to obtain

$$G_{\text{imp},\mathbf{k}\alpha}^K(\omega) = V_{\mathbf{k}\alpha} [G^R(\omega) g_{\mathbf{k}\alpha}^K + G^K(\omega) g_{\mathbf{k}\alpha}^A]. \quad (\text{A5})$$

Since the local $G^K(\omega)$ are purely imaginary, we can rewrite Eq. (A4) as

$$\begin{aligned} \langle \mathcal{I}_\alpha \rangle &= -\frac{e}{\hbar} \int \frac{d\omega}{2\pi} \sum_{\mathbf{k}\sigma} V_{\mathbf{k}\alpha}^2 (-\text{Im}[G^R(\omega)] \text{Im}[g_{\mathbf{k}\alpha}^K] + \text{Im}[G^K(\omega)] \text{Im}[g_{\mathbf{k}\alpha}^R]) \\ &= -\frac{e}{\hbar} \int \frac{d\omega}{2\pi} (-\Gamma_\alpha(\omega) \text{Im}[G^K(\omega)] + 2\text{Im}[G^R(\omega)] \Gamma_\alpha(\omega) (1 - 2F_\alpha(\omega))), \end{aligned} \quad (\text{A6})$$

where we have used

$$\sum_{\mathbf{k}\sigma} V_{\mathbf{k}\alpha}^2 \text{Im} g_{\mathbf{k}\alpha}^R = -2\Gamma_\alpha(\omega), \quad g_{\mathbf{k}\alpha}^K = 2i \text{Im}[g_{\mathbf{k}\alpha}^R] (1 - 2F_\alpha(\omega)). \quad (\text{A7})$$

Inserting Eq. (A6) in Eq. (A1) yields

$$\langle \mathcal{I} \rangle(t) = -\frac{2e}{\hbar} \int \frac{d\omega}{2\pi} \left(-\frac{1}{2} (\Gamma_L(\omega) - \Gamma_R(\omega)) (\text{Im}[G^K(\omega)] - 2\text{Im}[G^R(\omega)]) - 2\text{Im}[G^R(\omega)] (\Gamma_L(\omega) F_L(\omega) - \Gamma_R(\omega) F_R(\omega)) \right). \quad (\text{A8})$$

We can use

$$\text{Im}[G^K(\omega)] = 2\text{Im}[G^<(\omega)] + 2\text{Im}[G^R(\omega)] \quad (\text{A9})$$

to rewrite Eq. (A8) as

$$\langle \mathcal{I} \rangle(t) = -\frac{2e}{\hbar} \int \frac{d\omega}{2\pi} \{ (\Gamma_R(\omega) - \Gamma_L(\omega)) \text{Im}[G^<(\omega)] + 2(\Gamma_R(\omega) F_R(\omega) - \Gamma_L(\omega) F_L(\omega)) \text{Im}[G^R(\omega)] \}. \quad (\text{A10})$$

As pointed out in Ref. [72], in an equilibrium situation $F_R(\omega) = F_L(\omega)$ and $\text{Im}[G^<(\omega)] = -2\text{Im}[G^R(\omega)] F_{R/L}(\omega)$ so that the current vanishes. Due to the dephasing, however, the second condition is not fulfilled even in the case in which the first one is, i.e., leads in equilibrium. As pointed

out in Ref. [51], this can produce a nonreciprocal current flowing without bias under the absence of certain symmetry conditions. This is because the dephasing term can be seen also as a coupling to a bosonic bath at infinite temperature.

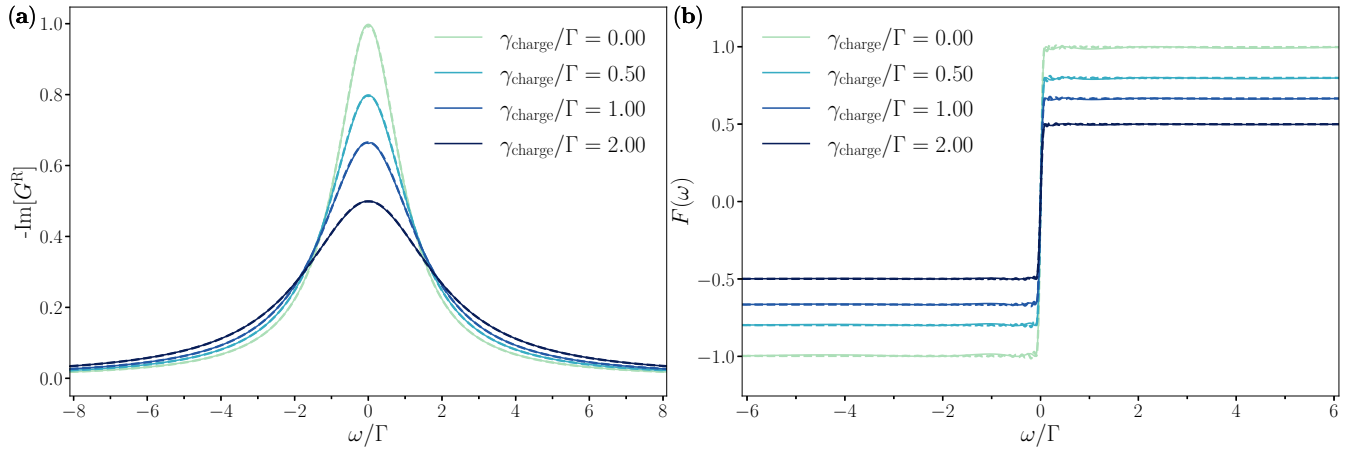


FIG. 4. Dissipative resonant level model with charge dephasing $L = \sqrt{\gamma_{\text{charge}}} \sum_{\sigma} n_{\sigma}$. (a) Impurity spectral function for the half filled dissipative resonant level model, corresponding to $U = 0$, and increasing values of γ_{charge} at fixed temperature $T = 0.02\Gamma$. (b) Corresponding distribution function $F(\omega)$. The solid lines represent results from the auxiliary master equation approach, while the dotted lines show the exact solution obtained via the Keldysh formalism.

APPENDIX B: AUXILIARY MASTER EQUATION APPROACH

The auxiliary master equation approach [64,69] is a nonequilibrium steady-state impurity solver and in the following we will give a brief overview. For details, the reader is referred to Refs. [64–66].

A correlated impurity problem is completely described by the hybridization function, which defines the uncorrelated fermionic reservoirs, the on-site energy, and the Hubbard interaction. Within AMEA, the “physical” fermionic reservoirs are mapped onto an “auxiliary” open system, which consists of a finite number of bath sites, N_B , that are in turn coupled to Markovian environments described by linear Lindblad terms. The Lindblad and hopping parameters of the auxiliary system are obtained by a fit to the physical hybridization function. The accuracy of the mapping is determined solely by the difference between the (retarded and Keldysh components) of the hybridization functions of the physical and the auxiliary reservoirs. This difference decreases exponentially with increasing N_B , as was shown Ref. [71], and is already quite small for $N_B \gtrsim 6$. This is due to the fact that the number of fit parameters is already quite large for this value of N_B . For finite, not too large N_B , the auxiliary many-body problem can be solved by exact methods such as Lanczos or Matrix product states, whereby the many-body density matrix has to be addressed. Concretely, one finds the steady-state density matrix and the Green’s functions directly in real frequencies. In this paper, we use a configuration interaction based approach [66], which uses Lanczos/Arnoldi for the computation of the steady state and the Green’s functions, which has been further improved by introducing a functional interpolation scheme [75]. Throughout this paper, we use $N_B = 8$.

APPENDIX C: ADDITIONAL RESULTS

In this Appendix, we present additional results on the dissipative Anderson impurity model: (1) benchmark of the

auxiliary master equation on the dissipative resonant level model; (2) spectral function and distribution function for the spin-dephasing case; and (3) the effect of the interaction U on the effective temperature.

1. Impurity spectral function and distribution function for the dissipative resonant level model

We now consider the dissipative but noninteracting case, with $U = 0$, corresponding to a dissipative Resonant Level Model. Due to the quadratic nature of the quantum jump, we stress that the system is still interacting, i.e., non-gaussian in the dissipative way. However, the specific nature of the charge and spin dephasing make it possible to close exactly the equations of motions for the retarded Green’s function in terms of a Dyson equation. [76] We obtain therefore in the real time domain

$$\Sigma(t, t') = \gamma_{\alpha} \delta(t, t') G(t, t'). \quad (\text{C1})$$

In Fig. 4, we plot the imaginary part of the impurity retarded Green’s function (left panel) and the corresponding distribution function obtained with the auxiliary master equation in the $U = 0$ case for increasing value of the charge dephasing rate. A comparison with the exact solution obtained from the Keldysh formalism is shown in dashed lines; we see an essentially perfect agreement between the auxiliary master equation approach and the exact result. We stress that this comparison is nontrivial and, for example, it fails for smaller number of bath sites in the AMEA approach.

As γ_{charge} increases, the resonance at the Fermi level is broadened and at the same time the value at $\omega = 0$ decreases with γ_{charge} , which is a signature of loss of coherence due to the Markovian environment. At the same time, the distribution function deviates from the equilibrium Fermi-Dirac form, reflecting the fact that the system is driven out of equilibrium and relaxes to a nonthermal state. Importantly, the dissipation acts uniformly across all frequencies, meaning that the dissipation does not discriminate between low- and high-energy

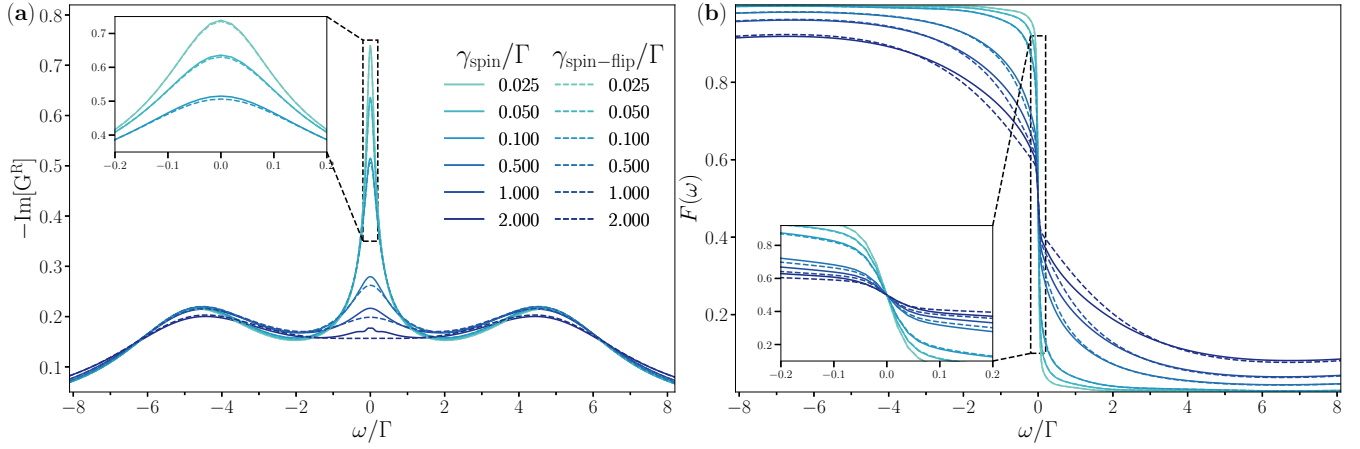


FIG. 5. Anderson impurity model with spin dephasing $L = \sqrt{\gamma_{\text{spin}}} \sum_{\sigma} \sigma n_{\sigma}$ and spin-flip noise [see Eq. (C3)]. (a) Impurity spectral function at half-filling for interaction strength $U = -2\epsilon_d = 8\Gamma$, shown for increasing values of the spin-dephasing rate γ_{spin} and spin-flip noise $\gamma_{\text{spin flip}}$. (b) Corresponding distribution function $F(\omega)$ as γ_{spin} and $\gamma_{\text{spin flip}}$ increase.

excitations; it affects all many-body states with equal strength, regardless of their energy.

In the noninteracting Anderson impurity model at half-filling, all states are dynamically equiprobable, and the dissipation does not preferentially target low- or high-energy excitations. As a result, the distribution function becomes progressively flatter with increasing dephasing, reflecting a uniform smearing of occupations across the spectrum. This energy-insensitive behavior is typical of dephasing processes, which tends to push the system toward a fully mixed state, similar to an infinite temperature.

2. Spectral properties of the Anderson impurity model subjected to spin dephasing and spin-flip noise

In addition to the charge dephasing case discussed in the main text, we present here further results concerning the spectral properties of the Anderson impurity model subjected to spin dephasing, which we recall is implemented via the

quantum jump operator

$$L_{\text{spin}} = \sqrt{\gamma_{\text{spin}}} \sum_{\sigma} \sigma n_{\sigma}, \quad (\text{C2})$$

and spin-flip fluctuations, which is modeled by the jump operators

$$\begin{aligned} L_{\text{spin flip}+} &= \sqrt{\gamma_{\text{spin flip}+}} S^{+} = \sqrt{\gamma_{\text{spin flip}+}} c_{\uparrow}^{\dagger} c_{\downarrow}, \\ L_{\text{spin flip}-} &= \sqrt{\gamma_{\text{spin flip}-}} S^{-} = \sqrt{\gamma_{\text{spin flip}-}} c_{\downarrow}^{\dagger} c_{\uparrow}, \end{aligned} \quad (\text{C3})$$

where we only consider the symmetric case of $\gamma_{\text{spin flip}} = \gamma_{\text{spin flip}+} = \gamma_{\text{spin flip}-}$. The latter is of importance, because it becomes relevant in any experimental setting.

In Fig. 5, we show the imaginary part of the Green's function and the corresponding distribution function for increasing values of the spin-dephasing rate γ_{spin} and the spin-flip noise rate $\gamma_{\text{spin flip}}$. For $\gamma_{\text{spin}} = \gamma_{\text{spin flip}} < T_K$, they are almost indistinguishable and even for higher values they do not differ much. Therefore, the following discussion applies to spin dephasing as well as spin-flip noise. The increase of γ_{spin}

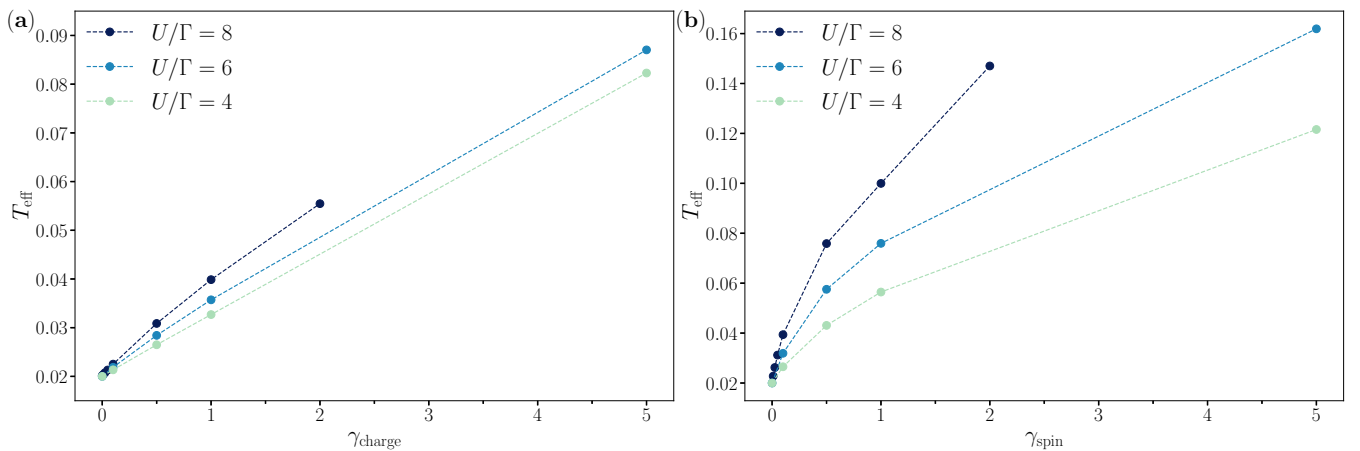


FIG. 6. Dissipative Anderson impurity model. Effective temperature T_{eff} of the dissipative Anderson impurity model as a function of the dissipation strength, for various values of the interaction U . The left panel corresponds to charge dephasing, while the right panel shows the case of spin dephasing. The leads are held at a finite temperature $T = 0.02\Gamma$ and are modeled with a flat (wideband) density of states.

leads to a pronounced suppression of the Kondo resonance in the spectral function. This behavior signals the breakdown of coherent spin fluctuations that are essential for the formation of the Kondo singlet. Notably, while the sharp Kondo peak is rapidly suppressed, the high-energy Hubbard bands remain largely unaffected by spin dephasing, indicating that charge excitations are more robust against this type of dephasing. The contrasting behavior between the Kondo peak and the Hubbard bands highlights that spin dephasing acts primarily on the low-energy spin degrees of freedom without significantly perturbing the charge sector [62].

This interpretation is further supported by the behavior of the distribution function. In contrast to charge dephasing, whose effect tends to be more frequency independent, spin dephasing predominantly alters the distribution function in the low-frequency region, corresponding to the energy window where spin fluctuations dominate. This selective suppression around $\omega \approx 0$ is consistent with the loss of Kondo coherence. On the other hand, the distribution remains almost unchanged at high frequencies $|\omega| \gg 1$, where charge fluctuations prevail, further confirming that spin dephasing acts mainly on the spin degrees of freedom and effectively suppresses Kondo

correlations without significantly affecting the charge fluctuation regime.

3. Effect of the interaction U on the effective temperature

In this Appendix, we provide additional results on the effective temperature T_{eff} extracted from the distribution function $F(\omega)$. Specifically, we discuss the influence of the Coulomb interaction U on T_{eff} for both the charge and spin dephasing. In Fig. 6, we plot the effective temperature T_{eff} as a function of dephasing rate γ_α (with $\alpha = \text{spin, charge}$) for different values of interaction U . The left panel corresponds to charge dephasing ($a = \text{charge}$) and the right panel to spin dephasing ($b = \text{spin}$). We see that in both cases interaction leads to an increase of the effective temperature, an effect that is particularly evident for spin dephasing where increasing U makes T_{eff} almost nonanalytic at small γ , as if heating in the spin-dephasing case was controlled by a nonperturbative effect. We can understand the strong dependence of heating from interaction in the spin-dephasing case by noticing that T_{eff} refers to the low-energy degrees of freedom, which are particularly sensitive to spin fluctuations in the strong interacting regime.

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