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Orthogonality Catastrophe and Decoherence in a Trapped-Fermion Environment

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The Fermi-edge singularity and the Anderson orthogonality catastrophe describe the universal physics which occurs when a Fermi sea is locally quenched by the sudden switching of a scattering potential, leading to a brutal disturbance of its ground state. We demonstrate that the effect can be seen in the controllable domain of ultracold trapped gases by providing an analytic description of the out-of-equilibrium response to an atomic impurity, both at zero and at finite temperature. Furthermore, we link the transient behavior of the gas to the decoherence of the impurity, and to the degree of the non-Markovian nature of its dynamics.

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A Fermi gas may be shaken up by the switching of even a single, weakly interacting impurity, producing a complete rearrangement of the many-body wave function, which loses essentially any overlap with the initial, unperturbed one. This is the essence of Anderson's orthogonality catastrophe [1,2], witnessed by the singular (edgelike) behavior of the excitation energy distribution. Such a many-body effect comes into play in x-ray photoemission spectra from most simple metals, where the expected sharp symmetric peak at the binding energy of a core level is converted into a power law singularity, as predicted by Mahan-Nozières–De Dominicis (MND) theory [3,4]. Similar patterns have been observed in electron emission from carbon based nanomaterials [5] and for quantum dots [6]. Fermi-edge resonance and orthogonality catastrophe have also been revealed by nonequilibrium current fluctuations in nanoscale conductors [7] and enter prominently the physics of phenomena as diverse as the Kondo effect [2,8] and the scattering or sticking of low-energy atoms or ions on metal surfaces [9,10].

Recently, it has been proposed to observe this universal physics with ultracold atoms, probing the singular behavior either in the time domain, by Ramsey interference [11], or in the frequency domain, by radio-frequency spectroscopy [12]. However, an analytic framework for the case of a *trapped* Fermi gas is lacking. In this Letter, we provide such an analytic description and discuss the transient response of a harmonically trapped Fermi gas following the *sudden* switching of an embedded two-level atom excited by a fast pulse. The interaction with the impurity produces a local quench of the gas, giving rise to the Anderson catastrophe. We study the Fermi-edge physics at zero and finite temperature and both in the frequency domain, by looking at the excitation spectrum of the gas, and in the time domain, by analyzing the dynamics of the impurity. Thus, we link the Fermi-edge behavior of the excitation energy distribution to the decoherence of the impurity. In particular, we investigate the Loschmidt echo [13,14] and non-Markovian nature, using recently developed tools [15–19], employed so far to study open systems in different environments, ranging from spins [20] to Bose-Einstein condensates [21], and experimentally tested in optical setups [22,23]. We find that the non-Markovian nature of the decoherent dynamics of the impurity provides a novel interpretation of the essential physics of the shakeup process.

We consider a gas of noninteracting cold fermions confined by a one-dimensional trapping harmonic potential of frequency ω, described by the Hamiltonian $\hat{H}_0 = \sum_{n,\xi} \varepsilon_n \hat{c}^\dagger_{n\xi} \hat{c}_{n\xi}$, with $\hat{c}_{n\xi}$ being the annihilation operator for the nth single-particle state of energy $\varepsilon_n = \hbar\omega(n+1/2)$ and spin ξ. We add a two-level impurity (an atom of a different species from the trapped component), with internal states $|g\rangle$, $|e\rangle$ and Hamiltonian $\hat{H}_I = \sum_{i=e,g} \epsilon_i |i\rangle\langle i|$, trapped in an auxiliary potential and brought in contact with the Fermi gas. This can be achieved using a species selective dipole potential that has a frequency much greater than the trap which contains the gas, so that the impurity motion is essentially frozen. We assume that when the impurity is in the $|g\rangle$ state, it has a negligible scattering interaction with the gas; hence, the Hamiltonian of the composite system is given by $\hat{H} = \hat{H}_0 + \hat{H}_I + \hat{V} \otimes |e\rangle\langle e|$. With the fermions in their equilibrium configuration, set by $\hat{H}_0$, we suppose the impurity to be quickly excited; the gas then feels a sudden perturbation $\hat{V}(t) = \hat{V}\theta(t)$, assumed to have an s-wave-like character.

At sufficiently low temperatures, the pseudopotential approximation is invoked, which amounts to replacing the complicated atomic interaction potential with an effective short range potential of strength V_0, localized at the minimum of the harmonic well, which we scale with the trap length x_0 such that $V(x) = \pi V_0 x_0 \delta(x)$. Because of the parity of the single-particle wave functions, only the fermions lying in even-parity states ($n = 2r$, with $r = 0, 1, 2, \dots$) feel the impurity and are involved in the shakeup process. Explicitly, the fermion-impurity interaction is given by $\hat{V} = \sum_{r,r',\xi} V_{rr'} \hat{c}^\dagger_{2r\xi} \hat{c}_{2r'\xi}$, where $V_{rr'} = V_0(-1)^{r+r'} \gamma_r^{1/2} \gamma_{r'}^{1/2}$, and $\gamma_r = 2^{-2r}\pi^{1/2}(2r)!/r!^2$ [24]. We label the highest occupied level by $n_F = 2r_F$, with r_F a positive integer, so that the Fermi energy reads $\varepsilon_F = \hbar\omega(2r_F + 1/2)$.

A key quantity for the following is the *vacuum persistence amplitude*

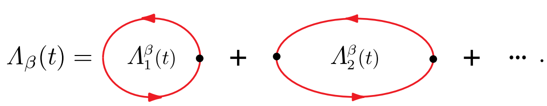
$$\nu_\beta(t > 0) = \langle e^{(i/\hbar)\hat{H}_0 t} e^{-(i/\hbar)(\hat{H}_0 + \hat{V})t} \rangle, \tag{1}$$

with $\langle \cdots \rangle$ denoting the grand canonical average over the unperturbed fermion state. $\nu_\beta(t)$ is the probability amplitude that the gas will retrieve its equilibrium state at time t, after the switching on of the perturbation, and its modulus gives the decoherence factor for the impurity (see below).

The Fourier transform $\tilde{\nu}_\beta(E)$ gives the excitation spectrum of the gas. In the interaction picture, we get

$$\begin{aligned} \nu_\beta(t) &= \langle T e^{(1/i\hbar) \int_0^t dt' \tilde{V}(t')} \rangle, \\ \tilde{V}(t) &= e^{(i/\hbar)\hat{H}_0 t} \hat{V} e^{-(i/\hbar)\hat{H}_0 t}, \end{aligned} \tag{2}$$

which, by virtue of the linked cluster theorem, reduces to an exponential sum of connected Feynman diagrams $\nu_\beta(t) = e^{\Lambda_\beta(t)}$, with



The closed graphs in $\Lambda_\beta(t)$ contain products of vertices ($V_{rr'}$), and lines (G_r^β) representing the unperturbed propagators

$$i\hbar G_r^\beta(t) = e^{-i\varepsilon_{2r}t/\hbar}[\theta(t) f_r^- - \theta(-t) f_r^+], \tag{3}$$

where $f_r^\pm = [1 + e^{\pm\beta(\varepsilon_{2r} - \mu)}]^{-1}$ are the particle-hole distributions, and μ denotes the chemical potential [24].

We focus on the lowest-order loops, namely,

$$\hbar\Lambda_1^\beta(t) = -it\chi_s V_0 \lambda_+^\beta(0), \tag{4}$$

$$\hbar^2 \Lambda_2^\beta(t) = -\chi_s V_0^2 \int_0^t dt' \int_0^{t'} dt'' \lambda_+^\beta(t'') \lambda_-^\beta(t''), \tag{5}$$

with $\chi_s = (2s + 1)$ accounting for the spin degeneracy and $\lambda_\pm^\beta(t) = \sum_{r=0}^\infty \gamma_r e^{\pm 2ir\omega t} f_r^\pm$.

This approximation will prove to accurately describe the singular response of the gas (contained in the two-vertex term) and to give the dominant contribution to the shakeup process if the interaction strength is small in the energy scale of the problem. The latter is set by both the level separation $\hbar\omega$ and Fermi energy ε_F, and we introduce $\alpha = \chi_s V_0^2 / 2\hbar\omega\varepsilon_F$ as a sensible interaction strength parameter.

The contribution (4) may be written as $\hbar\Lambda_1^\beta(t) = -itE_1^\beta$. Here,

$$E_1^\beta = \sqrt{2\chi_s \hbar\omega\varepsilon_F \alpha} \sum_{r=0}^\infty \gamma_r f_r^+ \tag{6}$$

is the first-order shift to the gas energy, as provided by the Rayleigh-Schrödinger perturbation theory. The behavior of the unperturbed energy $E_0^\beta = \chi_s \sum_n \varepsilon_n f_{n/2}^+$, and of its first and second-order corrections vs ε_F is shown in Fig. 1 for various temperatures. We notice that E_0^β is 1 to 3 orders of magnitude larger than E_1^β for $\alpha \lesssim 1$, while temperature plays an appreciable role in both E_0^β and E_1^β for $\beta\hbar\omega \lesssim 0.05$.

While $\Lambda_1^\beta(t)$ only brings a phase factor to $\nu_\beta(t)$, which corresponds to shifting the spectrum $\tilde{\nu}_\beta(E)$ by E_1^β, the two-vertex connected graph gives the crucial contribution to the persistence amplitude. It can be split into three parts with well defined trends and physical meaning [24], i.e., $\Lambda_2^\beta(t) = \Lambda_{2S}^\beta(t) + \Lambda_{2G}^\beta(t) + \Lambda_{2P}^\beta(t)$. These represent a (further) energy shift, a Gaussian envelope due to finite temperature effects, and periodic terms originating from the equal spacing of the unperturbed single-particle states, respectively, separately analyzed in Figs. 1(c), 2(a), and 2(b).

The first one $\hbar\Lambda_{2S}^\beta(t) = -itE_2^\beta$ provides the second-order correction to the energy of the gas (the $n >$ two-vertex graphs would complete the perturbation series):

$$E_2^\beta = \alpha\varepsilon_F \sum_{r \neq r' = 0}^\infty \frac{f_r^+ \gamma_r \gamma_{r'} f_{r'}^-}{r - r'}. \tag{7}$$

Comparing Figs. 1(b) and 1(c), we notice that the chosen value of α lets E_2^β take absolute values smaller than E_1^β. However, E_2^β is more sensitive to temperature than E_1^β for $\beta\hbar\omega < 0.05$.

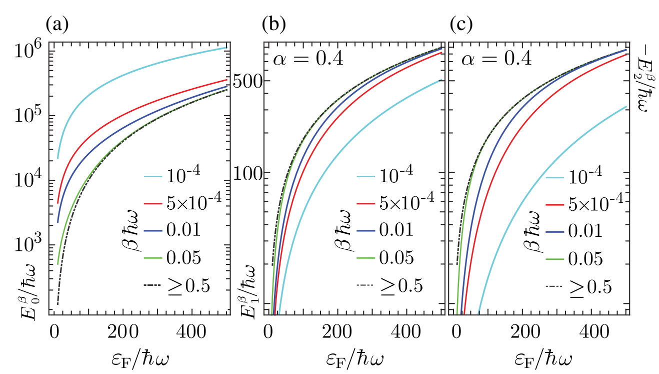


FIG. 1 (color online). Equilibrium energy E_0^β of a spin-1/2 gas into (a) a harmonic trap and perturbation corrections (b) E_1^β [Eq. (6)] and (c) E_2^β [Eq. (7)] due to the impurity potential $V(x)$. All energy curves are reported in units of $\hbar\omega$ vs $\varepsilon_F/\hbar\omega$ for different values of $\beta\hbar\omega$ and fixed coupling parameter $\alpha = 0.4$.

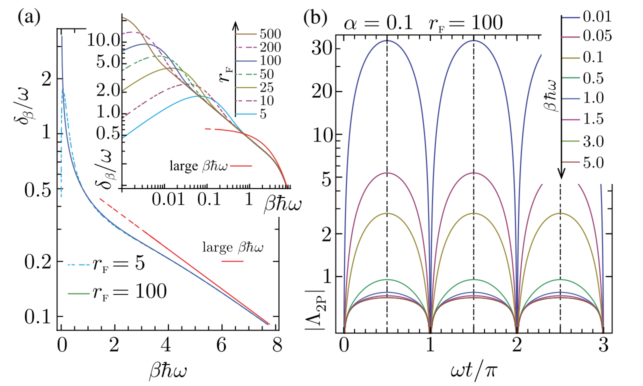


FIG. 2 (color online). (a) Standard deviation in the Gaussian power law (8), expressed in ω units, vs $\beta\hbar\omega$ for $r_F = 5$–100 and $\alpha = 0.1$. The low thermal energy approximation introduced in the text is also reported. (b) Periodic component $\Lambda^{\beta}_{2P}(t)$ vs $\omega t/\pi$, for $r_F = 100$, $\beta\hbar\omega = 0.01$–5, and $\alpha = 0.1$.

The second contribution $\Lambda^{\beta}_{2G}(t) = -\delta^2_{\beta}\omega^2 t^2/2$ produces a Gaussian damping in $\nu_{\beta}(t)$ and, therefore, a Gaussian broadening in $\tilde{\nu}_{\beta}(E)$ with standard deviation

$$\delta_{\beta} = \sqrt{2\alpha g_{\beta}}, \qquad g_{\beta} = \frac{\varepsilon_F}{\hbar\omega}\sum_{r=0}^{\infty}\gamma_r^2 f_r^+ f_r^-. \tag{8}$$

The coefficient g_{β} is weakly influenced by the Fermi energy but strongly affected by temperature, changing by various orders of magnitude for $\beta\hbar\omega \lesssim 0.5$. No damping or broadening effects are present at the absolute zero, since $\delta_{\beta} \to 0$ for $\beta\hbar\omega \to \infty$ [Fig. 2(a)].

The most important content of the second diagram, giving a nontrivial structure to $\nu_{\beta}(t)$, arises from the third contribution [24]:

$$\Lambda^{\beta}_{2P}(t) = -\frac{\alpha\varepsilon_F}{2\hbar\omega}\sum_{r\neq r'}^{\infty}\gamma_r f_r^+ \frac{1 - e^{2i(r-r')t\omega}}{(r-r')^2}\gamma_{r'} f_{r'}^-. \tag{9}$$

Because of the harmonic form of the trapping potential, this is periodic in time with frequency 2ω; see Fig. 2(b). The zeroes of this subgraph (at $\omega t = m\pi$ with $m = 0, \pm 1, \pm 2$), when combined with the Gaussian damping (8), yield modulations in the vacuum persistence amplitude which, as discussed below, are a signature of non-Markovian dynamics of the impurity.

Leaving aside the shifts, the persistence amplitude is then

$$\nu'_{\beta}(t) = e^{-\delta_{\beta}\omega^2 t^2/2} e^{\Lambda^{\beta}_{2P}(t)}. \tag{10}$$

Of particular interest for the discussion below is the behavior of $|\nu_{\beta}(t)|$ exhibiting spikes at $\omega t \sim \pi, 2\pi, \ldots$, which becomes more and more pronounced with increasing $\beta\hbar\omega$; see the left panels in Fig. 3. The periodicity in the time domain is reflected in the excitation spectrum $\tilde{\nu}_{\beta}(E)$ that offers an asymmetric, broadened signature of the singular

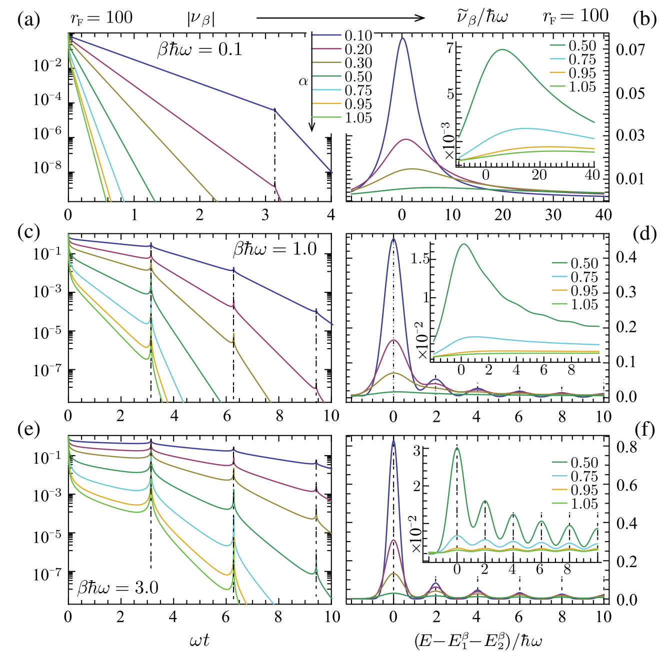


FIG. 3 (color online). Absolute value of (a),(c),(e) the decoherence factor $|\nu_{\beta}(t)|$ and (b),(d),(f) excitation spectrum $\nu'_{\beta}(E)$, calculated from Eq. (10) by numerically computing the Gaussian damping (8) and the periodic subdiagram $\Lambda^{\beta}_{2P}(t)$, for $\beta\hbar\omega = 0.1$–3.0, $r_F = 100$, and $\alpha = 0.1$–1.05.

behavior of the Fermi gas. The monotonic structure turns into a sequence of subpeaks, separated by $2\hbar\omega$ and related to even-level transitions in the gas as $\beta\hbar\omega$ gets above ~ 0.5 [see Fig. 3(b)]. These features are observed for any r_F in the range of 5 to 100 [24].

The coefficient (8) of the Gaussian power law and the periodic contribution $\Lambda^{\beta}_{2P}(t)$ can be approximated as [24]

$$g_{\beta} \approx 2\sum_{m=1}^{\infty}(-1)^m m\frac{e^{\beta\hbar\omega m/2}}{e^{\beta\hbar\omega m} - 1} \tag{11}$$

and

$$\Lambda^{\beta}_{2P}(t) \approx \alpha\sum_{m=-\infty}^{\infty}\ln\frac{e^{2\tau_m\omega} - 1}{e^{2(it+\tau_m)\omega} - 1}. \tag{12}$$

At low temperatures, the leading behavior of the Gaussian standard deviation is $\delta_{\beta} \approx 2\alpha^{1/2}e^{-\beta\hbar\omega/4}$ for thermal energies $\beta\hbar\omega \gtrsim 6$ [see Fig. 2(a)]. On the other hand, Eq. (12) contains a singularity at the absolute zero that we regularized by introducing a cutoff parameter τ_0. This regularization is needed to remove a zero temperature indefiniteness of the analytic approximation, whereas the numerical evaluation of the vacuum persistence amplitude does not suffer from divergence problems. As shown below and as detailed in Ref. [24], a similar parameter enters the original MND theory, and we can interpret it as the time scale over which transitions occur in the gas. On the other

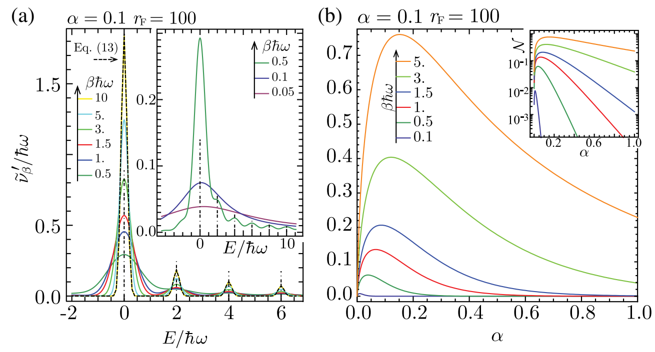


FIG. 4 (color online). (a) Absorption spectrum $\nu'_\beta(E)$, calculated numerically from Eq. (10) with $\beta\hbar\omega = 0.1$–∞, $r_F = 100$, and $\alpha = 0.1$, and analytical approximation $\nu'_\beta(E)$ obtained from Eqs. (11) and (12) with $\hbar\omega\beta = 10$ and $\omega\tau_0 = 0.001$. (b) Measure of the non-Markovian nature as a function of the critical parameter α for various temperatures.

hand, thermal fluctuations introduce other characteristic times $\tau_m = m\beta\hbar$.

Taking g_β and $\Lambda^\beta_{2P}(t)$ as in Eqs. (11) and (12) and using them in Eq. (10) gives an accurate approximation to the numerical results for $\beta\hbar\omega \gtrsim 0.1$, a number of particles larger than 10, and for suitable values of the cutoff parameter, say, $\omega\tau_0 < 0.02$ [see Fig. 4(a) and Ref. [24]]. In particular, at $T = 0$, the vacuum persistence amplitude takes the form

$$\nu'_{\beta\to\infty}(t) \approx \left[\frac{e^{2\tau_0\omega} - 1}{e^{2\omega(\tau_0 + it)} - 1}\right]^\alpha. \tag{13}$$

To compare our findings to the one-dimensional free-fermion theory, one needs to fix α and let the harmonic frequency go to zero by keeping the number of particles in the gas ($2r_F \approx \varepsilon_F/\hbar\omega$) finite. No Gaussian damping occurs in this case, and the two-vertex graph tends to

$$\Lambda_{\mathrm{MND}}(t) = -\alpha \ln(it/\tau_0 + 1), \tag{14}$$

yielding the Nozières–De Dominicis propagator $\nu_{\mathrm{MND}}(t) = (it/\tau_0 + 1)^{-\alpha}$, originally calculated for a suddenly switched-on core hole in a free electron gas [4]. Equation (14) was obtained by a long-time limit solution of the generalized Dyson equation for the electron Green's function in a constant window potential of width $\hbar/\tau_0$. For this reason, the MND spectrum lacks formal justification away from the threshold. In the present derivation, we have taken into account the full perturbation at an arbitrary time $t > 0$, retaining only the first nonadiabatic contribution in the linked cluster expansion [25]. We expect the effect of higher-order diagrams to be mainly concerned with the adiabatic correction to the equilibrium energy and some additional broadening of the excitation peaks. The latter should provide a renormalization to the critical parameter. Nevertheless, in the investigated ranges of temperatures and particle numbers, our definition of α produces a markedly singular response with the same range of criticality as the MND edge response parameter ($\alpha = 0$–1).

From this comparison with the free-gas case, we learn that the trapping frequency ω enters crucially the physics of the shakeup process. Indeed, it modifies the long-time response of the gas, as all single-particle excitations involve energy exchanges which are now even multiples of $\hbar\omega$. This gives rise to the periodic part of the fermion response and to the corresponding spectral peaks broadened at finite temperatures due to the Gaussian envelope, the latter being a typical effect of suddenly switched perturbations [2]. Up to now, we treated the response of the Fermi gas without any reference to the dynamics of the impurity that has just been assumed in the excited states for $t \geq 0$. If, instead, the two-level atom is subject to (say) a fast $\pi/2$ pulse and quickly prepared in the superposition $(|g\rangle + |e\rangle)/\sqrt{2}$, it experiences a purely dephasing dynamics due to the coupling with the gas, and its state at later times is $\rho_{\mathrm{IMP}}(t) = (|g\rangle\langle g| + |e\rangle\langle e| + \nu_\beta(t)|g\rangle \times \langle e| + \mathrm{H.c.})/2$. The decoherence factor entering the off-diagonal elements is just the persistence amplitude obtained before, going to zero at long times due to the orthogonality catastrophe. In the theory of open systems, one typically uses the so-called Loschmidt echo $L(t) = |\nu_\beta(t)|^2$, which gives a measure of the environmental response to the perturbation induced by the system [13,14,26] and which is linked with the non-Markovian nature of the open system dynamics [19]. The degree of the non-Markovian nature of a dynamical map can be evaluated in different manners [15–18], which are essentially equivalent for a purely dephasing quantum channel [19,27]. By adopting the information flow approach of Ref. [15], one finds

$$\mathcal{N} = \sum_n \{|\nu_\beta(t_{\max,n})| - |\nu_\beta(t_{\min,n})|\}, \tag{15}$$

where the summation is performed over all maxima and minima of $|\nu_\beta(t)|$, occurring at $t_{\max,n}$ and $t_{\min,n}$, respectively. Using our previous results for the amplitude, we obtain the non-Markovian nature of the dynamics of a two-level system in a trapped Fermi environment. The results are shown in Fig. 4(b), where we see that $\mathcal{N}$ depends on the temperature and on the critical parameter α. In particular, it has a maximum at small α, increasing with low temperatures, and goes to zero both for large temperatures (as thermal fluctuations suppress oscillations in the persistence amplitude) and for $\alpha > 1$. In the latter case, excitations are generated at every energy scale in the gas, as witnessed by the fact that the spectrum becomes structureless. This implies that the gas becomes more and more stiff (in the sense that it is not able to react on the impurity any more) and explains why $\mathcal{N}$ is zero: the open system does not receive information back, its Loschmidt echo decays monotonically and the dynamics is Markovian. A non-Markovian dynamics, then, can be characterized in our case by the appearance of specific spectral features in the excitation energy distribution [28].

We conclude with two remarks. First, the spectral distribution of energy excitations obtained here coincides with the so-called work distribution function, which is a central quantity in nonequilibrium processes [29,30]. In the setup described above, it is simple enough to conceive a "reverse" protocol, with the Fermi gas brought to thermal equilibrium in the presence of the impurity (i.e., with the two-level atom in the excited state) which is then switched off. The comparison of the work distribution functions in the direct and reverse protocols would lead to a direct experimental test of the Crooks relation in the quantum regime [31]. The second remark is on the experimental realization of the model that we have described. Many experiments have recently dealt with impurities in trapped Fermi gas [32], and state-dependent scattering lengths have been discussed [33]. This would lead to a direct test of our theory. Another viable candidate could be a gas of hard-core bosons in one dimension, where the Loschmidt echo is equivalent to that of the corresponding Fermi gas [34] and in which impurities have recently been experimentally generated [35].

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