Decoherence and Momentum Relaxation in Fermi-Polaron Rabi Dynamics: A Kinetic Equation Approach

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Despite the paradigmatic nature of the Fermi-polaron model, the theoretical description of its nonlinear dynamics poses challenges. Here, we apply a quantum kinetic theory of driven polarons to recent experiments with ultracold atoms, where Rabi oscillations between a Fermi-polaron state and a noninteracting level were reported. The resulting equations separate decoherence from momentum relaxation, with the corresponding rates showing a different dependence on microscopic scattering processes and quasiparticle properties. We describe both the polaron ground state and the excited repulsive-polaron state and we find a good quantitative agreement between our predictions and the available experimental data without any fitting parameter. Our approach not only takes into account collisional phenomena, but also it can be used to study the different roles played by decoherence and the collisional integral in the strongly interacting highly imbalanced mixture of Fermi gases.

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Introduction.—Polarons, i.e., impurities dressed by their environment, play a paradigmatic role in the understanding of many-body properties in a variety of physical systems. The Fermi polaron—a particle interacting with a reservoir of free fermions—became the object of intense study after its experimental realization with ultracold atomic gases [1–6], and more recently in monolayer semiconductors [7,8]. Studying this problem has offered valuable insights into the many-body physics of Fermi-Fermi or Bose-Fermi mixtures [4].

The high level of control allows the investigation of the dynamics of Fermi polarons [9]. In particular, experiments have studied the fate of coherent Rabi dynamics between a noninteracting impurity state and a polaron state [10–12], where collisional induced relaxation and decoherence are at play.

In thermal equilibrium, due to the simplicity of the Fermi-polaron model, fully solvable microscopic approaches, like simple variational Ansätze [13,14], have produced quantitatively valid predictions. However, the theoretical description of the problem becomes more challenging when considering the dynamics. For instance, the theoretical description of the Fermi-polaron Rabi dynamics based on a time-dependent variational approach revealed some nontrivial features [15], like the absence of decay from the repulsive to the attractive branch, which were required to achieve agreement with the experimental data [12].

Moreover, for such a problem—related to the open quantum system problem of a spin in a bath—it would be desirable to have a density matrix description and to identify the role of the various relaxation mechanisms.

In this Letter, we apply a quantum kinetic approach [16] to study the Fermi-polaron Rabi dynamics, which allows us to describe the system in terms of the evolution of the density matrix of a two-level system, whose properties are dressed by many-body effects. Importantly, we can distinguish between decoherence and momentum relaxation, whose rates show a different dependence on scattering processes and quasiparticle properties. Within a simple approximation for the scattering between the impurity and the bath, the solution of our quantum kinetic equations is shown to be in good agreement with the available experimental data without any fitting parameters. This indicates that the experiments realize the situation where

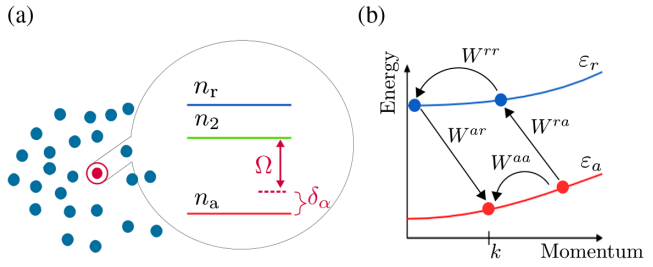

FIG. 1. (a) A highly imbalanced mixture of atoms in state $|1\rangle$ (majority, blue) and $|2\rangle$ (minority, red dot in a circle) is held at temperature T with zero interspecies interaction. The Rabi coupling Ω drives the transitions between states $|2\rangle$ and $|3\rangle$. In the latter, minority atoms occupy repulsive and attractive Fermi-polaron branches, formed due to the interaction with the majority component, while δ_α is the detuning between the noninteracting state and the polaron levels (here, only $\alpha = a$ is shown). (b) Polaron dispersion relations for atoms in state $|3\rangle$. The interaction with the majority atoms induces inter- and intraband transitions, with rates $W^{\alpha\beta}$, where $\alpha, \beta \in \{a, r\}$.

polarons (and not the bare particles) perform Rabi oscillations.

Model and kinetic equations.—We consider a homogeneous system composed of a bath of atoms in state $|1\rangle$, and impurity atoms in state $|3\rangle$, which are Rabi-coupled to a noninteracting state $|2\rangle$ [Fig. 1(a)]. All the states are different internal levels of the same isotope, and, therefore, they have the same mass m. The Hamiltonian of the system reads

$$\hat{H} = \sum_{i=1,2,3} \hat{H}_i + \hat{H}_{\rm int} + \hat{H}_\Omega, \tag{1}$$

where $\hat{H}_i = \sum_{\mathbf{k}} (\mathbf{k}^2/2m) \hat{c}^\dagger_{\mathbf{k},i} \hat{c}_{\mathbf{k},i}$, with $\hat{c}^{(\dagger)}_{\mathbf{k},i}$ the annihilation (creation) operators of a Fermi atom in the state $|i\rangle$ with momentum $\mathbf{k}$ (hereafter $\hbar = 1$) and $\hat{H}_{\rm int} = (U_0/V) \sum_{\mathbf{k},\mathbf{k}',\mathbf{q}} \hat{c}^\dagger_{\mathbf{k}'-\mathbf{q},1} \hat{c}^\dagger_{\mathbf{k}+\mathbf{q},3} \hat{c}_{\mathbf{k},3} \hat{c}_{\mathbf{k}',1}$ is the interaction potential. In absence of the Rabi coupling, when the impurity is in the state $|3\rangle$, the interaction with the bath leads to the emergence of quasiparticles known as Fermi polarons. In addition to a negative energy attractive branch, a, there exists a metastable repulsive, r, polaron branch at positive energy. The polaron dispersion relations at small momenta are given by $\varepsilon_\alpha(\mathbf{k}) \approx E_\alpha + \mathbf{k}^2/2m^*_\alpha$ ($\alpha = a$ or r), where E_α—with $E_a < 0$ and $E_r > 0$—and $m^*_\alpha > m$ are referred to as the polaron energy and polaron effective mass, respectively. The branches are also characterized by the quasiparticle weight $0 < Z_\alpha(\mathbf{k}) \leqslant 1$. The last term in Eq. (1) induces the Rabi oscillations between the states $|2\rangle$ and $|3\rangle$, and, in the rotating wave approximation, can be written as

$$\hat{H}_\Omega = \sum_{\mathbf{k}} \left[\frac{\Omega}{2} (\hat{c}^\dagger_{\mathbf{k},3} \hat{c}_{\mathbf{k},2} + \hat{c}^\dagger_{\mathbf{k},2} \hat{c}_{\mathbf{k},3}) + \Delta \hat{c}^\dagger_{\mathbf{k},2} \hat{c}_{\mathbf{k},2} \right], \tag{2}$$

where Ω is the bare Rabi frequency and Δ is the bare detuning from the transition $|2\rangle \leftrightarrow |3\rangle$.

In the experiments by Scazza *et al.* [10] for equal masses and by Kohstall *et al.* [9] for large mass imbalance, as well as by Oppong *et al.* [11] in quasi-2D geometry, it has been shown that it is possible to drive long-lived coherent Rabi oscillations between the noninteracting state and both the repulsive or the attractive polaron states. The resonant energies and the renormalized Rabi frequencies of the oscillations were found to be in reasonable agreement with an analysis based on the assumption that Ω does not affect the polaron properties. On the other hand, the decay rate of the oscillations, especially for the supposedly long-lived attractive polaron, has not found a proper explanation yet. Recently, a variational approach has been able to capture the dynamics for the repulsive branch [12].

Our aim is to provide an equation of motion for the impurity, which is able to take into account the quasiparticle nature of the polarons, and to explain how the static polaron properties modify the Rabi oscillations. The main result of our study is the set of equations for the following single-particle density matrix:

$$\dot{n}_\alpha - i \frac{Z_\alpha \Omega}{2} (f_{2\alpha} - f^*_{2\alpha}) = I_\alpha, \tag{3a}$$

$$\dot{n}_2 + i \frac{\Omega}{2} (f_{2\alpha} - f^*_{2\alpha}) = 0, \tag{3b}$$

$$\dot{f}_{2\alpha} - i \tilde{Z}_\alpha \delta_\alpha f_{2\alpha} + i \frac{\tilde{Z}_\alpha \Omega}{2} (n_2 - n_\alpha) = -\frac{\Gamma^{\rm dec}_\alpha}{2} f_{2\alpha}, \tag{3c}$$

where we dropped the time t and $\mathbf{k}$ for brevity, and $\tilde{Z}_\alpha = 2/(1 + 1/Z_\alpha)$. The effective detuning $\delta_\alpha(\mathbf{k}) = \varepsilon_\alpha(\mathbf{k}) - \varepsilon_2(\mathbf{k}) - \Delta \simeq 0$ determines whether the attractive $\alpha = a$ or the repulsive $\alpha = r$ polaron is involved in the dynamics. In this notation, $n_\alpha(\mathbf{k}, t)$ is the occupation of the polaron branch α at momentum $\mathbf{k}$, $n_2(\mathbf{k}, t)$ is the occupation of the state $|2\rangle$, which is coherently coupled to the branch α, and $f_{2\alpha}(\mathbf{k}, t)$ is the coherence between atoms in the state $|2\rangle$ and α polarons.

The quantum kinetic equations (3) can be obtained following the general approach established by Kadanoff and Baym [17] for time-dependent Green's functions, thereby extending previous works (see, e.g., [18,19]) on the derivation of the kinetic equations for spin-1/2 Fermi quantum fluids in magnetic fields [20] to three-level systems in the highly imbalanced case, i.e., in the impurity limit (see the Supplemental Material (SM) for more details [21]). Here, we briefly discuss the main approximations. Within the usual Kramers-Moyal [19] expansion, necessary to derive time-local equations, we drop the backflow term [36] (see also below) and assume that the polaron spectral properties, in particular the parameters Z_α, E_α, and m_α, are time-independent. Moreover, the equations are derived by projecting on the energy shell of the impurity interacting

with the equilibrium bath (see, e.g., [14]), i.e., the polaron branch unmodified by the driving laser.

The left-hand side of the kinetic equations in Eq. (3) predicts coherent oscillations with a renormalized Rabi frequency $\sqrt{Z_\alpha\Omega^2+\tilde{Z}_\alpha^2\delta_\alpha^2}$. For $\delta_\alpha=0$, we obtain $\sqrt{Z_\alpha}\Omega$, which is in agreement with the expression obtained by using a stationary variational Ansatz, which already received experimental verification [9,10]. There is instead no measurement of the role of the detuning in determining the polaron oscillation, which, according to our result, is nontrivial. Interestingly a very recent experiment [37] has investigated such dependence for large Ω paving the way to further analysis on the role of the dressing in the polaron dynamics.

The right-hand side of Eq. (3), due to collisions between minority and majority atoms, contains (i) the redistribution of the polaron population n_α, described by the collision integral I_α, and (ii) the loss of coherence between the α polaron and the noninteracting state, described by the decoherence rate $\Gamma_\alpha^{\mathrm{dec}}$. The population of the α polaron branch can be changed as a result of both inter- and intrabranch collisions. In the impurity limit $n_\alpha(\mathbf{k})\approx 0$, the decoherence rate reads

$$\Gamma_\alpha^{\mathrm{dec}}(\mathbf{k})=\frac{2}{1+Z_\alpha(\mathbf{k})}\frac{1}{V}\sum_{\mathbf{k}',\beta}W_{\mathbf{k}'\mathbf{k}}^{\beta\alpha}, \tag{4}$$

and the collision integral [38] takes the intuitive form $I_\alpha=\sum_\beta I_{\alpha\beta}$, where

$$I_{\alpha\beta}(\mathbf{k})=\frac{1}{V}\sum_{\mathbf{k}'}\left[W_{\mathbf{k}\mathbf{k}'}^{\alpha\beta}n_\beta(\mathbf{k}')-W_{\mathbf{k}'\mathbf{k}}^{\beta\alpha}n_\alpha(\mathbf{k})\right], \tag{5}$$

where $W_{\mathbf{k}'\mathbf{k}}^{\beta\alpha}$ is the transition rates from the α branch with momentum $\mathbf{k}$ to the β branch with momentum $\mathbf{k}'$.

The above expressions show that the redistribution of population and the decoherence, originating from collisions, are different in nature. The redistribution results from the imbalance between *in* and *out* scattering processes [see Eq. (5)], the density independent decoherence rate, Eq. (4), is due to possible scattering processes between a polaron α at momentum $\mathbf{k}$ and a polaron β at momentum $\mathbf{k}'$. The population redistribution drives the minority atoms toward thermal equilibrium with the majority ones, being indeed $I_{\alpha\beta}=0$ considering the Boltzmann equilibrium distribution $n_\alpha^{\mathrm{eq}}(\mathbf{k})\propto e^{-\varepsilon_\alpha(\mathbf{k})/T}$, with T the temperature of the majority component. As expected, notice that for decreasing quasiparticle weight the prefactor $2/(1+Z_\alpha)$ in Eq. (4) increases the role of the decoherence term compared to the redistribution rate.

The transition rates are due to the scattering of a polaron with an atom of the majority component. They are given by the Fermi golden rule

$$\begin{aligned}W_{\mathbf{k},\mathbf{k}'}^{\alpha\beta}=&\frac{2\pi}{V}\sum_{\mathbf{Q}}|T_{\mathrm{sc}}[\mathbf{Q},\varepsilon_\beta(\mathbf{k}')+\varepsilon_1(\mathbf{Q}-\mathbf{k}')]|^2Z_\alpha(\mathbf{k})Z_\beta(\mathbf{k}')\\&\times\delta[\varepsilon_\alpha(\mathbf{k})+\varepsilon_1(\mathbf{Q}-\mathbf{k})-\varepsilon_1(\mathbf{Q}-\mathbf{k}')-\varepsilon_\beta(\mathbf{k}')]\\&\times n_1^{\mathrm{eq}}(\mathbf{Q}-\mathbf{k}')[1-n_1^{\mathrm{eq}}(\mathbf{Q}-\mathbf{k})],\end{aligned} \tag{6}$$

where $T_{\mathrm{sc}}(\mathbf{Q},\varepsilon)$ is the scattering matrix with $\mathbf{Q}$ and ε the total momentum and the total energy of the particles entering the collision, $n_1^{\mathrm{eq}}(\mathbf{k})=1/[e^{(\varepsilon_1(\mathbf{k})-\mu)/T}+1]$ is the thermal distribution of the majority atoms with chemical potential μ, where $\varepsilon_1(\mathbf{k})=\mathbf{k}^2/2m$. Notice how the quasiparticle weights of the initial and final states renormalize the transition rates.

For the attractive polaron branch, the decoherence rate $\Gamma_a^{\mathrm{dec}}(\mathbf{k})$ is particularly simple at small momenta, since the scattering of the majority atoms takes place in the vicinity of the Fermi surface. We may write it as

$$\begin{aligned}\Gamma_a^{\mathrm{dec}}(k)\approx&\frac{\tilde{Z}_a(0)Z_a(0)m^2|T_{\mathrm{sc}}[k_F,\varepsilon_F+\varepsilon_a(k)]|^2}{4\pi^3}\\&\times\int_0^\infty k'^2dk'\frac{\varepsilon_a(k')-\varepsilon_a(k)+T\ln\left[\frac{1+e^{\frac{\varepsilon_F}{T}}}{e^{\frac{\varepsilon_F}{T}}+e^{\frac{\varepsilon_a(k')-\varepsilon_a(k)}{T}}}\right]}{\max(k',k)[e^{\frac{\varepsilon_a(k')-\varepsilon_a(k)}{T}}-1]},\end{aligned} \tag{7}$$

and therefore the timescale of the decoherence rate in the Rabi oscillations is set by $|T_{\mathrm{sc}}[k_F,\varepsilon_F+\varepsilon_a(0)]|^2$ and a k-dependent function originating from the phase space, where $\varepsilon_F=k_F^2/(2m)$ is the Fermi energy and k_F is the Fermi wave vector.

Setting and parameters.—In the following, we compare the solution of our kinetic equations to the experimental results from Ref. [10]. The impurity atoms are initially prepared in the noninteracting state $|2\rangle$ and in thermal equilibrium with the bath at temperature T. The initial condition for solving the kinetic equations are thus $n_\alpha=0$, $f_{2\alpha}=0$ and $n_2^{\mathrm{eq}}(\mathbf{k})=[e^{(\varepsilon_2(\mathbf{k})-\mu_2)/T)}+1]^{-1}$, where the chemical potential μ_2 fixes the imbalance $x\equiv\rho_2/\rho_1$, where $\rho_i=(1/V)\sum_{\mathbf{k}}n_i^{\mathrm{eq}}(\mathbf{k})$. We compute the polaron parameters $Z_\alpha(\mathbf{k})$ [39], E_α, m_α, and $T_{\mathrm{sc}}(\mathbf{k},\omega)$ within the non-self-consistent T-matrix approach [14]. In particular, the scattering matrix reads

$$T_{\mathrm{sc}}^{-1}(\mathbf{Q},\omega)=\frac{m}{4\pi a}-\frac{1}{V}\sum_{\mathbf{k}}\left(\frac{1-n_1^{\mathrm{eq}}(\mathbf{k}+\mathbf{Q}/2)}{\omega-\frac{\mathbf{Q}^2+4\mathbf{k}^2}{4m}+i0^+}+\frac{m}{\mathbf{k}^2}\right), \tag{8}$$

where a is the s-wave scattering length between the atoms in states $|1\rangle$ and $|3\rangle$, whose relation to the two-body contact potential reads $U_0^{-1}=m/(4\pi a)-(1/V)\sum_{\mathbf{k}}m/k^2$. This approach, even close to unitarity, compares reasonably well with experimental outcomes and quantum Monte Carlo calculations for the polaron's energy, mass and residue [10,40].

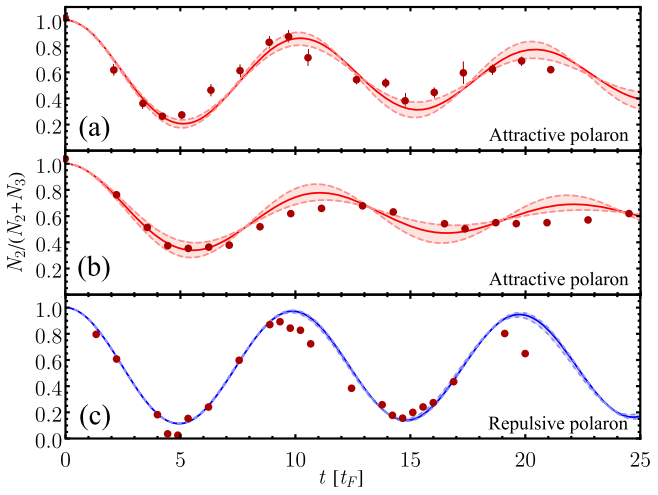

FIG. 2. Rabi oscillations of the attractive polarons for $1/(k_F a) = 0$ (a) and 0.25 (b) as a function of time t (in units of $t_F = 1/\varepsilon_F$) and of the repulsive polaron for $1/(k_F a) = 1.27$ (c). The red dots are from the experiment [10]. The solid red line is from Eq. (6) and the non-self-consistent T-matrix approximation in Eq. (6). The red shaded region shows the confidence interval for 20% relative uncertainty in the temperature. The parameters: $\Omega = 0.7\varepsilon_F$, $T = 0.135\varepsilon_F$.

To be consistent with the experiment [10], we set $T = 0.135\varepsilon_F$, $\Omega = 0.7\varepsilon_F$, and the imbalance $x = 0.15$. We take k_F equal to the effective κ_F from Ref. [10]. We numerically determine the observable $N_2/(N_2 + N_3)$, where N_i is the total atom number in the state $i = 2, 3$, and $N_2(t) + N_3(t) = N_2(t = 0)$.

Rabi oscillations.—To describe Rabi oscillations, we set $\delta_\alpha|_{\mathbf{k}=0} = 0$. In Figs. 2(a) and 2(b), we show the dynamics for $1/(k_F a) = 0$, 0.25, for which the polaron parameters $[E_a/\varepsilon_F, m_*/m, Z_a(0)]$ entering in the simulation are $(-0.625, 1.16, 0.775)$ and $(-0.858, 1.29, 0.673)$, respectively. The repulsive polarons are considered to not be populated, due to small Z_r and significant detuning. The shaded region takes into account the experimental uncertainty in the determination of the temperature, i.e., $T \pm \Delta T = 0.135\varepsilon_F(1 \pm 20\%)$. The agreement with the experimental data is quite remarkable. While the renormalized oscillation frequency $\sqrt{Z_a}\Omega$ can be explained just by static calculations, the decay rate within more standard variational approaches is usually found to be too small [15] or *ad hoc* assumptions have to be done [4]. Not only is our approach built to properly take into account collisional phenomena, but we disentangle the role played by decoherence and by the collisional integral. In particular, we find that for the Rabi oscillation dynamics, neglecting the collisional integral is still a good approximation. For completeness in Fig. 2(c), we show the comparison with the repulsive polaron experimental data [10] at $1/(k_F a) = 1.27$, where however the collisions are playing a minor role (see SM for further discussion).

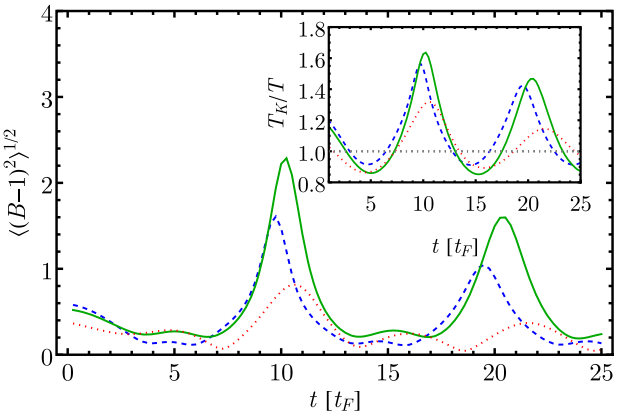

FIG. 3. Thermalization and violation of the detailed balance for attractive polarons. The main panel shows the degree of the violation of the detailed balance as quantified by $\sqrt{\langle [B(\mathbf{k}, \mathbf{k}') - 1]^2 \rangle}$. Inset shows the inferred temperature T_K (in the units of the bath temperature T) from the Kullback-Leibler divergence. Here, the parameters are $1/k_F a = -0.25$ (dashed blue), 0 (solid green), 0.25 (dotted red); $\Omega = 0.7\varepsilon_F$, $T = 0.135\varepsilon_F$, $\delta_a(0) = 0$.

Thermalization.—While we have found that the effect of the collisional integral I_α is negligible in the Rabi oscillation dynamics, it drives the system to a thermal state by redistributing the population. We focus again on the attractive polaron case, and consider the question how far from equilibrium the distribution $n_a(\mathbf{k}, t)$ is. To probe the violation of detailed balance, we define

$$B \equiv \frac{1 - n_a(\mathbf{k})}{n_a(\mathbf{k})} \frac{1 - n_1(\mathbf{k}_1)}{n_1(\mathbf{k}_1)} \frac{n_a(\mathbf{k}')}{1 - n_a(\mathbf{k}')} \frac{n_1(\mathbf{k}_1')}{1 - n_1(\mathbf{k}_1')}, \quad (9)$$

where we dropped the dependence on t. The two pairs of collision momenta $(\mathbf{k}', \mathbf{k}_1') \rightleftarrows (\mathbf{k}, \mathbf{k}_1)$ satisfy energy and momentum conservation laws. The values $B > 1$ ($B < 1$) indicate that, in the kinetic equation for $n_a(\mathbf{k})$, the rate of *in* processes is larger (smaller) than the *out* processes in the collision integral in the chosen momentum sector. In equilibrium, the *in* and *out* processes are in detailed balance and $B \equiv 1$ for all momenta. In our case, the bath is held at equilibrium and B becomes a function of $\mathbf{k}$ and $\mathbf{k}'$ only.

As an estimator of the lack of detailed balance we use $\sqrt{\langle [B(\mathbf{k}, \mathbf{k}') - 1]^2 \rangle}$, where the average is taken over the distribution $\propto n_a(\mathbf{k}, t) n_a(\mathbf{k}', t)$, whose dynamics is reported in Fig. 3. Our estimator shows a decaying oscillatory behavior in time with pronounced peaks—when the density of polarons is small—where the detailed balance is significantly violated, and the system is far from equilibrium (as can be seen also by looking at the momentum distribution itself; see SM).

In order to investigate the dynamics of thermalization, we introduce the concept of the inferred temperature T_K, which we define based on the Kullback-Leibler divergence $D(P|Q) \equiv -\sum_{\mathbf{k}} P_{\mathbf{k}} \ln(Q_{\mathbf{k}}/P_{\mathbf{k}})$, where P and Q are two

probability distributions. This non-negative object, widely used in the theory of nonequilibrium processes [41,42], which is nullified only if $P = Q$, quantifies the loss of information when the normalized distribution $Q_{\mathbf{k}}$ is used for the approximation of the true distribution $P_{\mathbf{k}}$; an optimization of the information loss is called as a moment projection [43]. To search for the best approximation of the polarons in terms of equilibrium states, we take $P_{\mathbf{k}} \propto n_a(t, \mathbf{k})$ and $Q_{\mathbf{k}} \propto n_a^{\mathrm{eq}}(T^*, \mathbf{k})$, which is a thermal state at a guessed temperature T^*. The optimal estimate of the temperature is obtained by the minimization of the Kullback-Leibler divergence over T^*, assuming that the densities of the distribution are the same. This optimal temperature, which we denote with T_K, is the information projection to a set of thermal states, and reads

$$T_K \equiv \underset{T^*}{\operatorname{argmin}} D(\tilde{n}_a(t)|\tilde{n}_a^{\mathrm{eq}}(T^*)), \tag{10}$$

where the tilde denotes normalization with respect to momentum $\mathbf{k}$. The chemical potential of equilibrium polarons $n_a^{\mathrm{eq}}(T^*, \mathbf{k})$ is adjusted so that its density coincides with the one given by $n_a(t, \mathbf{k})$.

In the inset of Fig. 3, we present the inferred T_K relative to the temperature T of the bath. We find that T_K is a meaningful characteristic of the distribution, as it yields a continuous function with the correct order of magnitude. The temperature T_K shows oscillations on the experimental timescale, with the peak value on the order of 1.3–$1.6T$ at times when the density of polarons is the lowest, i.e., at the maxima of the Rabi oscillations; cf. Fig. 2. In the long time limit, the temperature T_K approaches the bath temperature. In the examples, the fastest thermalization is observed for the case $1/(k_F a) = 0.25$.

Finally, close to thermal equilibrium, the violation of the detailed balance is small. We found that it is bounded by the Kullback-Leibler divergence, since for $n_a(t) \approx n_a^{\mathrm{eq}}$, we have

$$\langle [B(\mathbf{k}, \mathbf{k}') - 1]^2 \rangle \leqslant \zeta D[\tilde{n}_a(t)|\tilde{n}_a^{\mathrm{eq}}(T)], \tag{11}$$

where $\zeta > 0$ is a function of thermodynamic variables at T. Thus, in the approach to equilibrium, the Kullback-Leibler divergence D not only implies the inference of the correct temperature, i.e., $T_K \to T$, but also it sets a bound on the degree of the detailed balance violation. Our results show that on the timescale of the experiment, the system is far from the thermal state and a nonequilibrium description is required. More technical details on the thermalization process is discussed in SM.

Conclusions and perspectives.—In this Letter, we provided quantum kinetic equations for the single-particle density matrix of Fermi polarons coherently driven between two internal levels. In the impurity limit, and as long as the polarons are well defined quasiparticles, our kinetic equations take the intuitive form shown in Eq. (3)—a many-body version of a dressed two-level system. An important feature of our approach is that it manifestly separates decoherence from momentum relaxation mechanisms, the latter entering the equations as a collision integral that redistributes particles between different momenta. We compared our results with available experimental data for the ground state (attractive) polaron, finding good agreement without any fitting parameters. The theory is also applicable to the metastable (repulsive) polaron branch far from unitary limit. Our approach provides a general tool to study out-of-equilibrium problems related to the fundamental concept of quasiparticles in many-body quantum systems, such as impurity thermalization, generation of quasiparticles in presence of strong Rabi coupling, and repulsive-attractive polaron coherence, to mention a few.

Finally, very recently a new experiment on Rabi dynamics was reported [37] on which our theory could be tested. In the same experiment a sharp transition between weak and large Rabi drive strength has been observed. Our approach can be applied to this situation provided we use the dressed Rabi particles as initial ingredients. Such a route has been developed in [18] for Fermi liquid theory in large transverse magnetic fields.

The data presented in this article are available from [44].

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