

Rabi oscillations of strongly driven Bose polarons

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Understanding the dynamical behavior of quasiparticles is essential for uncovering quantum many-body phenomena. Among these phenomena, the polaron in ultracold atomic gases has attracted considerable interest due to its precise controllability. By engineering the underlying Hamiltonian, polarons serve as a versatile platform for studying both equilibrium properties and nonequilibrium dynamics. In this work, we investigate the quantum dynamics of strongly driven Bose polarons, where minority atoms are described as mobile impurities with spin-1/2, and majority atoms are bosons. The spin- $\uparrow$ impurity interacts with majority atoms through a tunable scattering length a , while the spin- $\downarrow$ impurity remains noninteracting. After turning on the Rabi coupling, we calculate the evolution of the total magnetization using a trial wavefunction. We identify the exhibition of anomalous Rabi oscillation and steady-state magnetization for $a > 0$, due to the interplay between attractive and repulsive polarons. Our results provide a concrete example that illustrates how Rabi oscillations are dressed by system-environment coupling.

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Introduction. The emergence of quasiparticles is a cornerstone of many-body physics, providing an effective description of complex interactions through renormalized excitations. From Landau's Fermi liquid theory to fractionalized excitations in topological phases, quasiparticles serve as a powerful conceptual framework for describing collective behavior. Among these, the polaron, an impurity dressed by excitations of a surrounding medium, plays an important role in understanding the properties of various physical systems, ranging from solid-state materials to ultracold atomic gases [1–4]. In ultracold atomic gases, the setup consists of minority atoms immersed in a cloud of majority atoms at low temperatures. The interaction between the minority atoms and the majority atoms can be tuned using the Feshbach resonance [5]. Depending on the statistics of the majority atoms, the polaron is classified as either a Bose polaron [6–20, 22–40, 40–61] or a Fermi polaron [62–88]. A Bose polaron forms when minority atoms are dressed by phonon excitations, whereas a Fermi polaron arises from the excitation of particle-hole pairs. Studies also highlight the signature of Efimov physics in the Bose polaron with large mass imbalance [59, 89–96].

Most studies on Bose/Fermi polarons focus on spectral functions near thermal equilibrium, typically measured using ejection and injection spectroscopy. Experimentally, this requires weak coupling of the internal state of the minority atoms to an auxiliary state, where the interaction between the

minority and majority atoms is negligible. After a short time, the population of the minority atoms is measured, following Fermi's golden rule [28, 29, 69, 71, 72, 97]. Recently, there has been growing interest in exploring the opposite regime, where a strong Rabi coupling is applied for a sufficiently long time, allowing the system to reach a steady state [87]. On the one hand, this regime investigates the dressing of Rabi oscillations in the presence of system-environment coupling, with a close relation to quantum thermalization. On the other hand, it provides a concrete example of the fate of quasiparticles under a strong driving field. For Fermi polarons, theoretical calculations near unitarity using the Green's function approach have shown good agreement with experimental results [85]. However, an analysis of strongly driven Bose polarons is still lacking.

In this work, we take the first step toward understanding the Rabi oscillations of strongly driven Bose polarons by formulating a trial wavefunction approach [62] to investigate their nonequilibrium dynamics. When the majority atoms are noninteracting, our approach allows for an analytical treatment, providing an accurate description of the dynamical evolution for arbitrarily long times. The results reveal the emergence of anomalous Rabi oscillations and steady-state magnetization on the Bose-Einstein condensate (BEC) side, where the attractive and repulsive branches coexist. We explain these phenomena by relating the Fourier transform of the impurity polarization density over time to the steady-state spectral functions, which can be approximated using perturbation theory. Furthermore, we provide numerical evidence that introducing a weak interaction between majority atoms does not lead to a qualitative change in our conclusions. We expect that our theoretical prediction can be readily tested on state-of-the-art experimental platforms.

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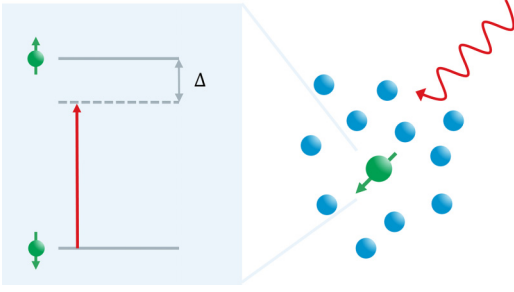


FIG. 1. We present a schematic of the setup of our system. We consider an impurity atom (green) immersed in a cloud of bosonic atoms (blue). The impurity has two internal states, labeled by $\uparrow$ and $\downarrow$. Initially, the impurity is prepared in the state $|\downarrow\rangle$, while the interaction between the impurity and the environment is only present for $|\uparrow\rangle$. We study the dressed Rabi oscillation of the system after a strong radio-frequency field is turned on, using the trial wavefunction approach.

Model. Our model consists of a single spin-1/2 impurity immersed in a cloud of bosonic atoms, as illustrated in Fig. 1. The bosons either are noninteracting or have weak repulsive interactions. Both cases can be captured by the Bogoliubov Hamiltonian

$$H_b = \sum_{k \neq 0} \gamma_k \beta_k^\dagger \beta_k. \quad (1)$$

Here, $\gamma_k = \sqrt{\epsilon_{b,k}(\epsilon_{b,k} + 2g_b n_0)}$ is the dispersion of Bogoliubov quasiparticles. We define $\epsilon_{b,k} = k^2/2m_b$ as the dispersion for noninteracting bosons, and g_b is related to the scattering length a_b between bosons by the relation $g_b = 4\pi a_b/m_b$. The annihilation operator of the Bogoliubov quasiparticle β_k is related to the annihilation and creation operators of atoms $b_k/b_{-k}^\dagger$ via the transformation $\beta_k = u_k b_k + v_k b_{-k}^\dagger$, where $u_k^2 = \frac{1}{2}(\frac{\epsilon_{b,k} + g_b n_0}{\gamma_k} + 1)$ and $u_k^2 - v_k^2 = 1$. The impurity has two internal states, $|\uparrow\rangle$ and $|\downarrow\rangle$, which are coupled through a radio-frequency field with Rabi coupling Ω_0 and detuning Δ for $t > 0$. The Hamiltonian of the impurity reads

$$H_i(t) = \sum_k \left(\epsilon_{i,k} f_{k,\uparrow}^\dagger f_{k,\uparrow} + (\epsilon_{i,k} + \Theta(t)\Delta) f_{k,\downarrow}^\dagger f_{k,\downarrow} + \frac{\Omega_0}{2} \Theta(t) (f_{k,\downarrow}^\dagger f_{k,\uparrow} + f_{k,\uparrow}^\dagger f_{k,\downarrow}) \right). \quad (2)$$

Here, $f_{k,\sigma}$ annihilates an impurity atom with spin $\sigma \in \{\uparrow, \downarrow\}$ and $\epsilon_{i,k} = k^2/2m_i$. $\Theta(t)$ is the Heaviside step function. Since we focus on the single-impurity sector, the statistics of the impurity are irrelevant to our analysis. Nevertheless, we assume the impurity is fermionic for concreteness. Finally, the impurity interacts with the bosons only if it is in the state $|\uparrow\rangle$. This leads to

$$H_{bi} = \frac{g}{V} \sum_{k,k',q} b_{k'}^\dagger f_{q-k',\uparrow}^\dagger f_{q-k,\uparrow} b_k. \quad (3)$$

The interaction term has an ultraviolet divergence, which is regularized by the renormalization relation $\frac{1}{g} = \frac{m}{2\pi a} - \frac{1}{V} \sum_k \frac{2m}{k^2}$, where $m^{-1} = m_b^{-1} + m_i^{-1}$ is the inverse reduced

mass. The Hamiltonian of the total system is therefore given by $H(t) = H_b + H_i(t) + H_{bi}$.

Trial wavefunction. For $t < 0$, the impurity is prepared in the spin- $\downarrow$ state with zero momentum, and the majority atoms are in the ground state with no quasiparticle excitation. This corresponds to a product state

$$|\psi(0)\rangle = f_{0,\downarrow}^\dagger |0\rangle_i |\text{BEC}\rangle_b \equiv f_{0,\downarrow}^\dagger |G\rangle, \quad (4)$$

where $\beta_k |\text{BEC}\rangle_b = 0$. At $t = 0$, the Rabi coupling is turned on and the system begins to evolve. In particular, after the transition into the spin- $\uparrow$ state, the impurity-boson interaction can excite Bogoliubov quasiparticles. Truncating the Hilbert space to at most one excitation, we write out the trial wavefunction

$$|\psi(t)\rangle = \sum_\sigma \left(\psi_\sigma^{(0)}(t) f_{0,\sigma}^\dagger + \frac{1}{\sqrt{V}} \sum_k \psi_{k,\sigma}^{(1)}(t) f_{-k,\sigma}^\dagger \beta_k^\dagger \right) |G\rangle. \quad (5)$$

Here, V is the size of the total system. The summation over k implicitly excludes $k = 0$. The factor of $1/\sqrt{V}$ is introduced such that wavefunctions satisfy the normalization condition

$$\sum_\sigma \left[|\psi_\sigma^{(0)}(t)|^2 + \int \frac{d\mathbf{k}}{(2\pi)^3} |\psi_{k,\sigma}^{(1)}(t)|^2 \right] = 1. \quad (6)$$

Similar trial wavefunctions have been employed to study both the equilibrium properties and quench dynamics of Bose polarons in the absence of an external drive [18,32,33,44]. It serves as a good approximation across the Feshbach resonance, capturing the crossover between the polaron state and the molecule state for the attractive branch, as well as the presence of a repulsive branch for $a > 0$. It is also straightforward to include contributions from two-excitation sectors, which contain three-body correlations. This is particularly important for Bose polarons with a significant mass imbalance, as the Efimov effect gives rise to qualitative signatures in the spectral function [95]. In this work, we instead focus on the simplest scenario with $m_b = m_i$, which is realized by using different hyperfine states of the same atomic species for both the impurity and the bosonic environment.

Leaving the details to the Supplemental Material [98], for a generic scattering length a_b , the evolution of the wavefunction reads

$$\begin{aligned} i\partial_t \psi_\uparrow^{(0)} &= g n_0 \psi_\uparrow^{(0)} + \frac{\Omega_0}{2} \psi_\downarrow^{(0)} + \frac{g}{V} \sqrt{n_0} \sum_k \mathcal{R}_k \psi_{k,\uparrow}^{(1)}, \\ i\partial_t \psi_{k,\uparrow}^{(1)} &= (\epsilon_k + g n_0) \psi_{k,\uparrow}^{(1)} + \frac{\Omega_0}{2} \psi_{k,\downarrow}^{(1)} \\ &\quad + g \sqrt{n_0} \mathcal{R}_k \psi_\uparrow^{(0)} + \frac{g}{V} \sum_{k'} \mathcal{D}_{k,k'} \psi_{k',\uparrow}^{(1)}, \\ i\partial_t \psi_\downarrow^{(0)} &= \Delta \psi_\downarrow^{(0)} + \frac{\Omega_0}{2} \psi_\uparrow^{(0)}, \\ i\partial_t \psi_{k,\downarrow}^{(1)} &= (\epsilon_k + \Delta) \psi_{k,\downarrow}^{(1)} + \frac{\Omega_0}{2} \psi_{k,\uparrow}^{(1)}. \end{aligned} \quad (7)$$

Here, we have introduced $\epsilon_k = \epsilon_{i,k} + \gamma_k$, $\mathcal{R}_k = u_k - v_k$, and $\mathcal{D}_{k,k'} = u_k u_{k'} + v_k v_{k'}$ for conciseness. n_0 is the density of the condensate. Since we focus on a weakly interacting or non-interacting Bose gas, we do not distinguish n_0 from the total

density of the majority atoms. The initial condition is $\psi_{\downarrow}^{(0)} = 1$, while all other components of the wavefunction vanish. In the following sections, we analyze this set of equations to determine the dynamics of strongly driven Bose polarons. In particular, we focus on the magnetization of the impurity, defined as

$$\mathcal{M}(t) = |\psi_{\uparrow}^{(0)}(t)|^2 - |\psi_{\downarrow}^{(0)}(t)|^2 + \int \frac{d\mathbf{k}}{(2\pi)^3} (|\psi_{k,\uparrow}^{(1)}(t)|^2 - |\psi_{k,\downarrow}^{(1)}(t)|^2). \quad (8)$$

Results. We first focus on the limit of $a_b = 0$, where bosons are noninteracting. Equation (7) is greatly simplified since $u_k = 1$ and $v_k = 0$. After performing the Laplace transform of the wavefunction $f(s) \equiv \int_0^\infty e^{-st} f(t) dt$, we can obtain analytical expressions for all components of the wavefunction.

The central observation is that we have $\frac{g}{V} \sum_k \mathcal{R}_k \psi_{k,\uparrow}^{(1)} = \frac{g}{V} \sum_k \mathcal{D}_{k',k} \psi_{k,\uparrow}^{(1)} = \frac{g}{V} \sum_k \psi_{k,\uparrow}^{(1)}$. Therefore, Eq. (7) allows all components of the wavefunction in terms of $\frac{g}{V} \sum_k \psi_{k,\uparrow}^{(1)}$. For example, this leads to the relation

$$\begin{aligned} \psi_{\uparrow}^{(0)}(s) &= \frac{\frac{i\Omega_0}{2} + (is - \Delta)\sqrt{n_0}\chi(s)}{D_0}, \\ \psi_{k,\uparrow}^{(1)}(s) &= \frac{(is - \epsilon_k - \Delta)\chi(s)}{D_k}, \end{aligned} \quad (9)$$

where, for conciseness, we define $D_k = (is - \epsilon_k)(is - \epsilon_k - \Delta) - (\frac{\Omega_0}{2})^2$. We also introduce $\chi(s) = g\sqrt{n_0}\psi_{\uparrow}^{(0)}(s) + \frac{g}{V} \sum_k \psi_{k,\uparrow}^{(1)}(s)$. Combining this definition with Eq. (9), we derive the result for $\chi(s)$ as

$$\left(\Pi(is) - \frac{(is - \Delta)n}{D_0} \right) \chi(s) = \frac{i\Omega_0\sqrt{n}}{2D_0}, \quad \Pi(is) = \frac{m}{2\pi a} - \int \frac{d\mathbf{k}}{(2\pi)^3} \left[\frac{is - \epsilon(k) - \Delta}{D_k} + \frac{2m}{k^2} \right], \quad (10)$$

where we have used the renormalization relation to absorb the divergence in $\Pi(is)$. In the limit of $\Omega_0 = 0$, $\Pi(E)|_{\Omega_0=0} \equiv \Sigma(E)$ matches (the inverse of) the two-body T matrix. The integration over $\mathbf{k}$ can be performed analytically for arbitrary Ω_0 , and its explicit form is provided in the Supplemental Material [98]. This leads to a closed-form expression for both $\psi_{\sigma}^{(0)}(s)$ and $\psi_{k,\sigma}^{(1)}(s)$. However, we are not able to perform the inverse Laplace transform analytically. Therefore, we numerically determine $\psi_{\sigma}^{(0)}(t)$ and $\psi_{k,\sigma}^{(1)}(t)$ (for a discretized set of k values) and then compute the magnetization using Eq. (8).

Results: Short-time oscillation. Numerical results in both the BEC ($k_n a = 1$) and Bardeen-Cooper-Schrieffer (BCS) ($k_n a = -1$) regimes, for different values of detuning Δ , are presented in Fig. 2, where the density is parametrized as $n_0 = k_n^3/6\pi^2$. We summarize several interesting features of

the magnetization dynamics. At short times, we observe many-body Rabi oscillations with a damped amplitude due to coupling with the environment. For either detuning $\Delta \gtrsim 0$ or $k_n a = -1$, the oscillation is regular and approximately single frequency. This is a close analog of Rabi oscillations of noninteracting spin-1/2 particles. In contrast, the oscillation becomes highly irregular for $k_n a = 1$ and $\Delta/k_n^2 \in \{-2, -1.34\}$, indicating a large deviation from the noninteracting case. We refer to this phenomenon as anomalous Rabi oscillation.

To uncover the underlying physical origin, we study the impurity spectral function in the presence of the Rabi coupling. This is done by analyzing the eigenenergy of Eq. (7), replacing $i\partial_t$ with E . Straightforward calculations yield $E = \Delta + \frac{(\Omega_0/2)^2}{E - n\Pi^{-1}(E)}$. This equation should be interpreted as looking for the pole of the impurity Green's function

$$G_{\downarrow}^R(\Omega) = \frac{1}{\Omega - \Delta - \frac{(\Omega_0/2)^2}{\Omega - n\Pi^{-1}(\Omega + i0^+)}}. \quad (11)$$

Here, we have introduced the retarded Green's function $G_{\downarrow}^R(t) \equiv -i\theta(t)\langle\{f_{0,\downarrow}(t), f_{0,\downarrow}^\dagger(0)\}\rangle$. Then, we can compute the spectral function using $A_{\downarrow}(\Omega) = -\frac{1}{\pi}\text{Im}(G_{\downarrow}^R(\Omega))$. The numerical results for $A_{\downarrow}(\Omega)$ are presented in Figs. 3(a) and 3(b) for $k_n a = -1$ and $k_n a = 1$, respectively. For $k_n a < 0$, there are two branches of excitations, formed due to the avoid-crossing between the attractive polaron formed by the spin- $\uparrow$ state and the noninteracting spin- $\downarrow$ state. This is well captured by a perturbative analysis of Ω_0 , which makes the single-pole approximation $\frac{1}{\Omega - n\Pi^{-1}(\Omega)} \approx \frac{Z_a}{\Omega - E_a}$. Here, the polaron energy E_a and the quasiparticle residue Z_a are determined by solving $E_a = \Sigma(E_a)$ and $Z_a = (1 - \partial_\Omega \Sigma(E_a))^{-1}$. Under this approximation, the excitation energies Ω^* are given by

$$\Omega^* = \frac{E_a + \Delta \pm \sqrt{(E_a - \Delta)^2 + Z_a \Omega_0^2}}{2} \quad \text{for } a < 0, \quad (12)$$

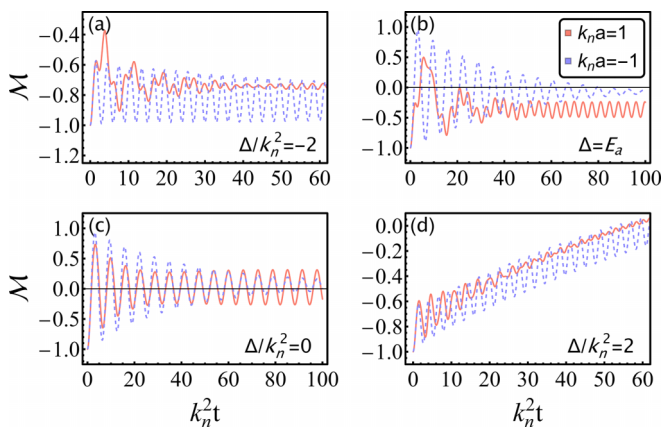


FIG. 2. We present numerical results for the magnetization dynamics during Rabi oscillations with $k_n a_b = 0$ and $\Omega_0/k_n^2 = 1$ for different impurity-boson interactions with $k_n a = \pm 1$. The analysis includes four different detunings: $\Delta = -2k_n^2, E_a, 0, 2k_n^2$, where E_a represents the energy of the attractive polaron (without the Rabi coupling), which satisfies $E_a = \Sigma(E_a)$. For $k_n a = 1$, this leads to $\Delta \approx -1.34k_n^2$.

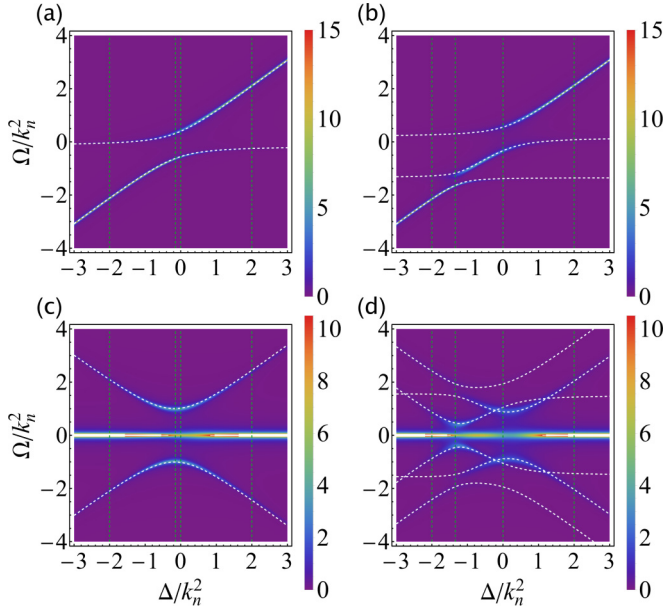


FIG. 3. We present the density plot for the impurity spectral function $A_I(\Omega)$ [panels (a) and (b)] and the Rabi spectrum [panels (c) and (d)]. Here, we choose $k_n a = -1$ for panels (a) and (c), and $k_n a = 1$ for panels (b) and (d). In all plots, we set $k_n a_b = 0$ and $\Omega_0/k_n^2 = 1$. Vertical dashed lines show the choices of Δ for Fig. 2. White dashed lines show the theoretical predictions from the perturbative calculation, which match the numerical results with good accuracy.

which are plotted in Fig. 3(a) using white dashed lines. A similar relation has been discovered for driven Fermi polarons [85]. For comparison, there are three branches of excitations for $k_n a > 0$ due to the coexistence of attractive and repulsive polarons [83]. Therefore, we approximate $\frac{1}{\Omega - n\pi^{-1}(\Omega)} \approx \frac{Z_a}{\Omega - E_a} + \frac{1 - Z_a}{\Omega - E_r}$, where E_r is the energy of the repulsive polaron, satisfying $E_r = \text{Re}[\Sigma(E_r)]$. This leads to the equation for excitation energies

$$(\Omega^* - \Delta) = \frac{(\Omega_0/2)^2 Z_a}{\Omega^* - E_a} + \frac{(\Omega_0/2)^2 (1 - Z_a)}{\Omega^* - E_r} \quad \text{for } a > 0. \quad (13)$$

Finding explicit solutions is possible, though tedious. We plot the results in Fig. 3(b) using white dashed lines, which also match the numerical results to good accuracy.

To proceed, we consider the magnetization dynamics. We perform a Fourier transform of $N_{\downarrow}^{(0)} \equiv |\psi_{\downarrow}^{(0)}(t)|^2$, which is referred to as the Rabi spectrum. By definition, $N_{\downarrow}^{(0)}$ is quadratic in the wavefunction. Therefore, it is natural to expect that the spectrum is determined by the energy differences between excitation energies Ω^* . This prediction is demonstrated in Figs. 3(c) and 3(d). For $a < 0$, the two excitation energies in Eq. (12) lead to a single peak at $\Omega > 0$. As a result, the magnetization is nearly single frequency. For $a > 0$, the coexistence of attractive and repulsive polarons gives rise to a complex Rabi spectrum. When Δ is negative but has a small magnitude, we find a finite spectral weight in all branches. Moreover, the hybridization between the attractive polaron and the spin- $\downarrow$ state becomes significant due to the renormalization of the Rabi coupling. These effects lead to the observed anomalous

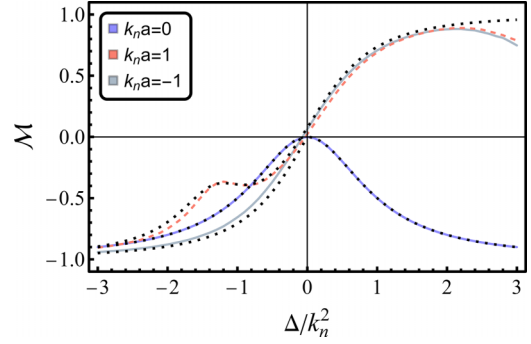


FIG. 4. We present the steady-state magnetization as a function of detuning Δ for $k_n a = 0, 1, -1$, with fixed parameters $k_n a_b = 0$ and $\Omega_0/k_n^2 = 1$. A prominent peak near $\Delta \approx -1.34 k_n^2 \approx E_a$ serves as a signature of the attractive polaron (see discussions in the main text). The decrease for $k_n a = \pm 1$ at large Δ arises from the insufficient evolution time to fully reach the steady state.

Rabi oscillations in the magnetization, exhibiting a sharp contrast with noninteracting spins.

Results: Steady-state magnetization. We now examine the long-time behavior of many-body Rabi oscillations, where the system approaches its steady state. The steady-state magnetization is defined as $\overline{M} \equiv \frac{1}{T} \int_{t_0}^{t_0+T} dt \mathcal{M}(t)|_{T \rightarrow \infty}$. For noninteracting spins, standard Rabi oscillations yield $\overline{M} = -\frac{\Delta^2}{\Delta^2 + \Omega_0^2}$, which decreases monotonically for $\Delta > 0$. This behavior arises because large detuning suppresses the transitions between the $|\uparrow\rangle$ and $|\downarrow\rangle$ states. When interactions between the impurity and bosons are introduced, the $|\downarrow\rangle$ state can couple to the $|\uparrow\rangle$ state by exciting an additional phonon. This interaction mechanism leads to a qualitative difference in the steady-state magnetization for large Δ , as shown in Fig. 4. The decrease for $k_n a = \pm 1$ at large Δ arises from the insufficient evolution time to fully reach the steady state. For $k_n a = -1$, the steady-state magnetization is well approximated by the ground state of noninteracting quantum spins, resulting in [87,98]

$$\overline{M} = \frac{\Delta}{\sqrt{\Delta^2 + \Omega_0^2}} \quad \text{for } a < 0. \quad (14)$$

In contrast, for $k_n a = 1$, we observe the emergence of an additional peak near $E = E_a$. This feature is attributed to the contribution from attractive polarons, which significantly alters the behavior of the system. We propose a simple formula

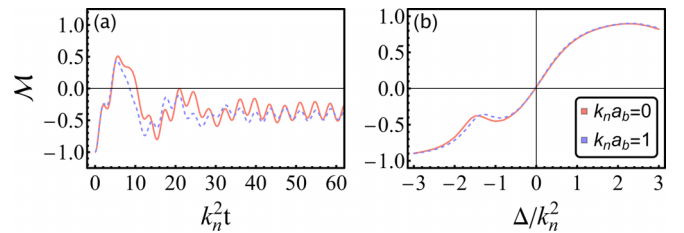


FIG. 5. We present numerical results for the magnetization dynamics and the steady-state magnetization during Rabi oscillations for $k_n a_b = 0, 1$, with fixed parameters $k_n a = 1$ and $\Omega_0/k_n^2 = 1$. In panel (a), we choose $\Delta = E_a \approx -1.34 k_n^2$. Both show a weak dependence on $k_n a_b$.

to capture both contributions, which reads

$$\overline{\mathcal{M}} = \frac{\Delta}{\sqrt{\Delta^2 + \Omega_0^2}} + \frac{Z^2 \Omega_0^2}{(\Delta - E_a)^2 + Z\Omega_0^2} \quad \text{for } a > 0. \quad (15)$$

Z is the quasiparticle residue of the second solution of Eq. (13) at $\Delta = E_a$. These predictions are validated in Fig. 4 using black dotted lines.

Results: Interacting BEC. Finally, we show that the above signatures of Rabi oscillations in driven Bose polarons remain valid even after introducing weak repulsive interactions, $k_n a_b > 0$. In this case, the solution of Eq. (7) becomes more involved and is deferred to the Supplemental Material [98]. The numerical results for the magnetization evolution $\mathcal{M}(t)$ are presented in Fig. 5, exhibiting the expected anomalous Rabi oscillation behavior for $k_n a > 0$ and $\Delta < 0$. The magnetization of the steady state also shows a weak dependence on $k_n a_b$. This demonstrates the stability of these intriguing dynamical behaviors.

Summary. In this work, we have explored the quantum dynamics of strongly driven Bose polarons in ultracold atomic gases, revealing rich nonequilibrium behavior

through our trial wavefunction approach. Our analysis has uncovered anomalous Rabi oscillations and steady-state magnetization patterns that emerge due to the interplay between attractive and repulsive polaron branches, particularly in the BEC regime. These findings demonstrate how system-environment coupling fundamentally alters the dynamics of driven quasiparticles, bridging the gap between equilibrium polaron physics and nonequilibrium quantum dynamics and providing insights for understanding and controlling quasiparticle behavior in strongly correlated systems.

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Data availability. The data that support the findings of this article are not publicly available. The data are available from the authors upon reasonable request.

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