

Phenomenological model of decaying Bose polarons

R. Alhyder^{1,2,*}, G. M. Bruun², T. Pohl^{2,3}, M. Leshchko¹, and A. G. Volosniev^{2,1}¹*Institute of Science and Technology Austria (ISTA), Am Campus 1, 3400 Klosterneuburg, Austria*²*Department of Physics and Astronomy, Aarhus University, Ny Munkegade 120, DK-8000 Aarhus C, Denmark*³*Institute for Theoretical Physics, Vienna University of Technology, Wiedner Hauptstraße 8–10, 1040 Vienna, Austria*

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Cold atom experiments show that a mobile impurity particle immersed in a weakly interacting Bose-Einstein condensate forms a well-defined quasiparticle (Bose polaron) for weak to moderate impurity-boson interaction strengths, whereas a significant line broadening is consistently observed for strong interactions. Motivated by this, we introduce a phenomenological theory based on the assumption that the most relevant states are characterized by the impurity correlated with at most one boson, since they have the largest overlap with the uncorrelated states to which the most common experimental probes couple. These experimentally relevant states can, however, decay to lower energy states characterized by correlations involving multiple bosons, and we model this using a minimal variational wave function combined with a complex impurity-boson interaction strength. We first motivate this approach by comparing to a more elaborate theory that includes correlations with up to two bosons. Our phenomenological model is shown to recover the main results of two recent experiments probing both the spectral and the nonequilibrium properties of the Bose polaron. Our work offers an intuitive framework for analyzing experimental data and highlights the importance of understanding the complicated problem of the Bose polaron decay in a many-body setting.

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Introduction. Polaron quasiparticles are key to our description of mobile particles interacting with a surrounding many-body medium. Quantum simulation experiments with cold atoms have provided a wealth of detailed information advancing our understanding of such polarons across a wide range of parameter regimes and interaction strengths, regarding both their spectral [1–14] and nonequilibrium properties [15–21]. This in turn, reinforced by a wealth of theoretical advancements on two-dimensional polaron physics [22–29], has been used for interpretation of experiments that create polarons in two-dimensional semiconductors, demonstrating the universal importance of this problem [30].

In spite of these advances, experimental data indicate that our understanding of impurity particles in Bose-Einstein condensates (BECs) is incomplete. In particular, the observed spectral “polaronic” signal is always broad for strong interactions, even for recent experiments performed in box potentials where the density is homogeneous apart from at the edges, which should reduce any density broadening [20]. The consistent observation of a broad spectral peak rather than a sharp quasiparticle peak naturally raises questions about the spectral weight and even existence of a long-lived Bose polaron for strong interactions. Second, the nonequilibrium impurity

dynamics for repulsive interactions were recently observed to lose coherence much more rapidly than expected [19,20].

A major reason for our incomplete understanding of the Bose polaron is that contrary to the Fermi polaron, there is no Pauli exclusion principle to suppress correlations between the impurity and $N \geq 2$ bosons, including the possible formation of bound states (trimers, tetramers, etc.) [31].

On the one hand, such few-body correlations in a many-body setting make the theoretical description of Bose polarons challenging [30]. On the other hand, experiment shows that relatively simple theories such as the ladder approximation [32] including only two-body correlations are able to recover the peak position (but not the width) of the spectral signal fairly well, as well as the nonequilibrium formation dynamics except the damped beating signal between the repulsive and attractive polarons [19].

Motivated by this, we explore here the conjecture that the experimentally important polaron states are characterized by the impurity being dressed by at most one boson. These many-body states do not necessarily correspond to the ground state but rather to broad resonances embedded in a many-body continuum. In other words, these states have a finite lifetime due to decay to low-lying states with a small spectral weight.

We first frame this conjecture by analyzing the spectral properties of a mobile impurity in a BEC using a well-known variational wave function that includes up to two Bogoliubov modes. This analysis suggests that the polaronlike state where the impurity is dressed by one boson appears as a broad resonance with a large spectral weight. This state lies above the long-lived state with a lower energy but also a small spectral weight. Instead of microscopically calculating the

*Contact author: ragheed.alhyder@ist.ac.at

decay rate of the polaron state, which is very challenging as it requires the inclusion of all relevant decay channels in a many-body system, we develop a phenomenological extension of the ladder approximation that describes the dominant spectral response. We then show that this phenomenological theory captures both the spectral and nonequilibrium features observed in two recent experiments [19,20].

Framework. We consider an impurity particle of mass m immersed in a homogeneous weakly interacting BEC with density n_0 . Since our focus is on the recent ^{39}K experiments [19,20], we assume that the mass of the bosons is also m , but the theory can be easily extended to mass-imbalanced systems. We use the two-channel Hamiltonian $\hat{H} = \hat{H}_0 + \hat{H}_{\text{int}}$ to describe the system with

$$\begin{aligned} \hat{H}_0 &= \sum_{\mathbf{k}} \epsilon_{\mathbf{k}} \hat{c}_{\mathbf{k}}^\dagger \hat{c}_{\mathbf{k}} + \sum_{\mathbf{k}} \xi_{\mathbf{k}} \hat{\beta}_{\mathbf{k}}^\dagger \hat{\beta}_{\mathbf{k}} + \sum_{\mathbf{k}} (k^2/4 + E_0) \hat{d}_{\mathbf{k}}^\dagger \hat{d}_{\mathbf{k}}, \\ \hat{H}_{\text{int}} &= g \sum_{\mathbf{k}} \left(\sqrt{n_0} \hat{d}_{\mathbf{k}}^\dagger \hat{c}_{\mathbf{k}} + \sum_{\mathbf{q}} \hat{d}_{\mathbf{q}}^\dagger \hat{c}_{\mathbf{q}-\mathbf{k}} \hat{b}_{\mathbf{k}} \right) + \text{H.c.}, \end{aligned} \quad (1)$$

where the operators $\hat{c}_{\mathbf{k}}^\dagger/\hat{\beta}_{\mathbf{k}}^\dagger/\hat{b}_{\mathbf{k}}^\dagger$ create an impurity/Bogoliubov mode (phonon)/boson with momentum $\mathbf{k}$, and $\hat{d}_{\mathbf{k}}^\dagger$ creates a closed channel (bare Feshbach) molecule, a two-body bound state of the impurity with one bath boson. We have as usual $\hat{b}_{\mathbf{k}} = u_{\mathbf{k}} \hat{\beta}_{\mathbf{k}} - v_{\mathbf{k}} \hat{\beta}_{-\mathbf{k}}^\dagger$ with $u_{\mathbf{k}} = \sqrt{(\epsilon_{\mathbf{k}} + n_0 g_B)/(2\xi_{\mathbf{k}}) + 1/2}$, $v_{\mathbf{k}} = \sqrt{(\epsilon_{\mathbf{k}} + n_0 g_B)/(2\xi_{\mathbf{k}}) - 1/2}$, and $g_B = 4\pi a_B$, with $a_B > 0$ being the boson-boson scattering length. The energies are $\epsilon_{\mathbf{k}} = k^2/2$ and $\xi_{\mathbf{k}} = \sqrt{\epsilon_{\mathbf{k}}(\epsilon_{\mathbf{k}} + 2g_B n_0)}$. Here and in what follows, we take the system volume, $\hbar$, and m to be unity. The bare molecule energy E_0 , open-closed channel coupling g , and ultraviolet momentum cutoff Λ can be eliminated in favor of the boson-impurity scattering length a and effective range R_e [33,34], which also gives rise to the three-body parameter $a_- = 2467 R_e$ determining when the first Efimov trimer becomes stable [35,36].

To find the eigenstates of the Hamiltonian in Eq. (1), we use the variational ansatz [37,38]

$$\begin{aligned} |\Psi\rangle &= \left(A \hat{c}_0^\dagger + \sum_{\mathbf{k}} B_{\mathbf{k}} \hat{c}_{-\mathbf{k}}^\dagger \hat{\beta}_{\mathbf{k}}^\dagger + \frac{1}{2} \sum_{\mathbf{k}\mathbf{p}} C_{\mathbf{k},\mathbf{p}} \hat{c}_{-\mathbf{k}-\mathbf{p}}^\dagger \hat{\beta}_{\mathbf{k}}^\dagger \hat{\beta}_{\mathbf{p}}^\dagger \right. \\ &\quad \left. + D \hat{d}_0^\dagger + \sum_{\mathbf{k}} E_{\mathbf{k}} \hat{d}_{-\mathbf{k}}^\dagger \hat{\beta}_{\mathbf{k}}^\dagger \right) |\text{BEC}\rangle. \end{aligned} \quad (2)$$

The first three terms describe the dressing of the impurity by $N = 0, 1, 2$ Bogoliubov modes, whereas the fourth and fifth terms give the closed channel molecule components of the wave function including the dressing by up to $N = 1$ Bogoliubov mode. Such a wave function has been used to successfully recover the experimental results regarding the spectral maximum of an impurity in a BEC [21,30]. Here, we will also use the ansatz in Eq. (2) with time-dependent coefficients $A(t)$, $B_{\mathbf{k}}(t)$, $C_{\mathbf{k},\mathbf{p}}(t)$, $D(t)$, and $E_{\mathbf{k}}(t)$ to study quench dynamics [36].

For the ground state, the variational parameters in Eq. (2) can be found by minimizing the expectation value $\langle \Psi | \hat{H} | \Psi \rangle$ of the Hamiltonian. To overcome the challenge of accessing the excited states, we use the Krylov subspace method [39–42], which iteratively constructs a reduced basis that captures the

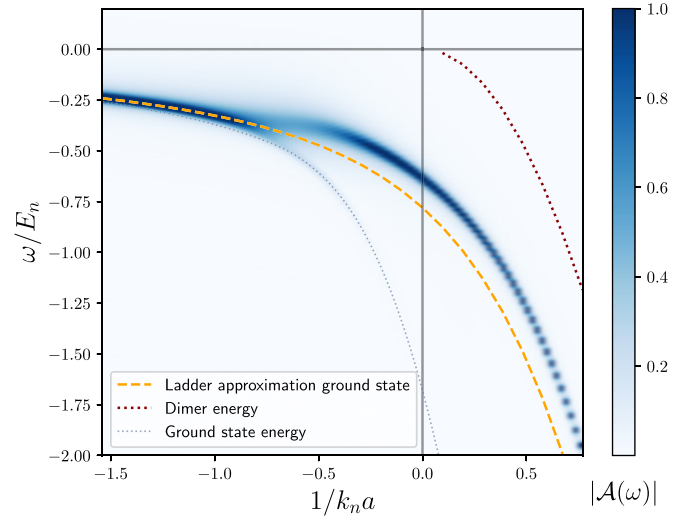


FIG. 1. The zero-momentum impurity spectral function calculated using the ansatz in Eq. (2) for $n_0^{1/3} a_- = -1$ and $n_0^{1/3} a_B = 0.02$. The ground-state energy is shown as a faint dotted curve, the dashed yellow curve shows the ground-state energy within the ladder approximation, and the dotted curve shows the energy of the dimer, $-1/a^2$. We have defined $E_n = (6\pi^2 n_0)^{2/3}/2$ and $1/k_n = 1/\sqrt{2E_n}$. Note that the spectral function is computed numerically for a discrete set of interaction strengths, which accounts for the apparently discontinuous behavior of the high-energy branch in the repulsive regime.

action of the Hamiltonian on the initial state, thereby forming a Krylov subspace [36]. It offers efficient access to the low-energy spectrum, essential for the present study. To construct the spectral function from the spectrum, we apply the standard finite-broadening prescription; that is, the delta peaks associated with the discrete eigenspectrum are replaced by narrow Lorentzian functions.

Spectral function. The zero-momentum impurity spectral function obtained from the Krylov subspace method is shown in Fig. 1 for the parameters $n_0^{1/3} a_- = -1$ and $n_0^{1/3} a_B = 0.02$ allowing for a direct comparison with Ref. [38]. Near the unitary limit $1/k_n a = 0$, the spectral function features two branches: a sharp low-energy branch, which is barely visible, reflecting its vanishing spectral weight, in agreement with the properties ground state of the problem studied in Ref. [38], and a broad high-energy branch present at strong interactions. We furthermore see that the maximum in the spectral density of this branch is in the vicinity of the results of the ladder approximation, which is equivalent to including terms with maximally one Bogoliubov mode in Eq. (2). This indicates that the upper branch corresponds to states dressed by one boson, which decay to the ground state dressed by two bosons within this approximation. Importantly, the spectral weight of the upper branch vastly exceeds that of the ground-state branch suggesting that it is the primary contributor to most experimental observables. We find that these qualitative features are robust for $n_0^{1/3} |a_-| \gtrsim 1$.

Phenomenological model. The ansatz in Eq. (2) has been extended to include up to three Bogoliubov modes for $a_B = 0$ and $1/a$ obtaining no clear convergence for the ground-state

energy [35], which may reflect an underlying bosonic orthogonality catastrophe for an impurity in an ideal BEC for a single-channel interaction [43]. In general, calculating the full spectral response from both ground and excited states encounters scalability challenges with the number of included Bogoliubov modes. Furthermore, the Hamiltonian in Eq. (1) neglects boson-boson interactions beyond mean field, which makes the accuracy of this approach unclear for states involving the dressing by multiple bosons [30,44–46]. We therefore adapt in the following a simpler phenomenological approach.

Instead of the multichannel model given by Eq. (1), we now use the interaction

$$\hat{\mathcal{H}}_{\text{int}} = g_I \sum_{\mathbf{k}, \mathbf{q}} c_{\mathbf{k}+\mathbf{q}}^\dagger c_{\mathbf{k}} b_{\mathbf{k}'-\mathbf{q}}^\dagger b_{\mathbf{k}'} \quad (3)$$

For real values of the interaction strength g_I , Eq. (3) is simply the standard single-channel point interaction, but we will instead employ a *complex* g_I where the imaginary part introduces an imaginary part to the energy [47–51]. This allows us to describe the decay and decoherence of the polaron state due to any omitted part of the Hilbert space in a relatively straightforward way. An imaginary part of g_I can in fact be shown to naturally emerge due to $N \geq 3$ -body losses [36], but we will in the following not provide a microscopic analysis of the origin of this loss, regarding it instead as a free parameter.

Moreover, since we know that the experimentally observed maximum of the spectral response is fairly accurately recovered by the ladder approximation [30], which is equivalent to a variational wave function truncated at a single Bogoliubov mode level, we fix $C = D = E = 0$ in Eq. (2). As Fig. 1 shows, the physical reason behind this approximation is that the states dressed with at most one boson have a large overlap (in comparison to low-lying states) with the noninteracting plane wave state to which both rf and Ramsey experiments couple. Consequently, the dominant experimental signal does not, in general, track the ground state but an excited, one-boson-dressed polaron that decays into low-lying states.

The calculated spectral function depends on a_B rather weakly in the ladder approximation [30], and to gain physical intuition we first consider an ideal BEC with $a_B = 0$, which reproduces key spectral features at $T = 0$ [52]. In this case, the pole of the impurity propagator is found by solving

$$\omega \left(\frac{1}{g_I} + \frac{\Lambda}{2\pi^2} - i \frac{\sqrt{\omega}}{4\pi} \right) = n_0. \quad (4)$$

To fit the experimental data, we use

$$\frac{1}{g_I} \simeq \frac{1}{4\pi a} - \frac{\Lambda}{2\pi^2} + i \frac{\Gamma + \gamma n_0^{1/3} |a|}{4\pi |a|}. \quad (5)$$

The first two terms are the standard prescription giving an undamped attractive polaron with energies $\omega \rightarrow 4\pi n a$ and $\omega \rightarrow -1/a^2$ for $a \rightarrow 0^-$ and $a \rightarrow 0^+$ respectively, and a repulsive polaron with energy $\omega \rightarrow 4\pi n a$ for $a \rightarrow 0^+$. As discussed above, these energies recover reasonably well the maximum of the observed spectroscopic signal.

The last term in Eq. (5) moves the pole into the complex plane, thereby introducing damping. In the spirit of our phenomenological approach, this term is constructed to

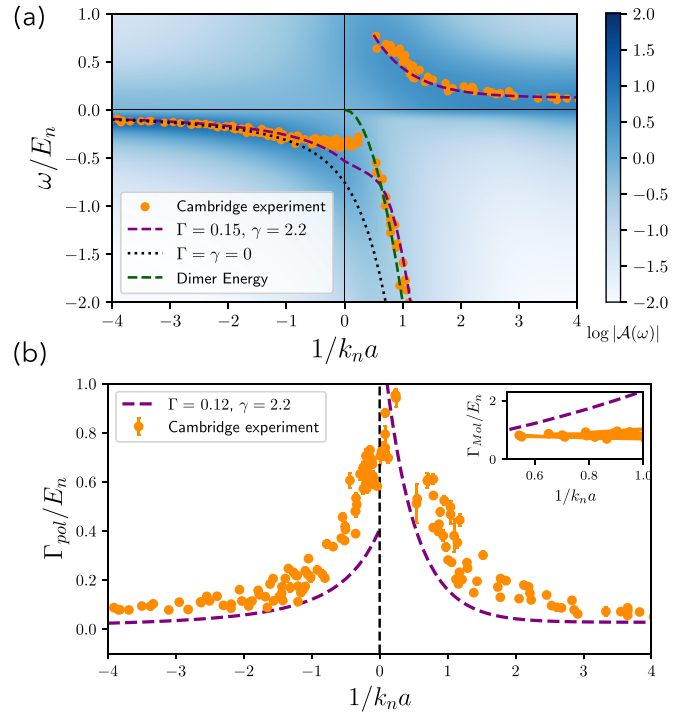


FIG. 2. (a) The zero-omentum impurity spectral function from our phenomenological model. The black dotted curve shows the ladder approximation with real g_I , the green dashed curve is the molecular state, and the orange dots are the experimental data of Ref. [20]. (b) Extracted widths Γ_{pol} of the attractive polaron for $a < 0$ and repulsive polaron for $a > 0$. To match the experiment [20], we use $n^{1/3} a_B = 0.005$, and for fitting the signal we use $\Gamma = 0.12$ and $\gamma = 2.2$. Inset: Extracted width Γ_{Mol} of the attractive polaron for $a > 0$.

interpolate between the observed linewidth in the weakly and strongly interacting regimes rather than being based on microscopic calculations. For weak interactions $a \rightarrow 0$, we find from Eqs. (4) and (5) that the complex energy is $\omega \simeq 4\pi n a - i 4\pi n |a| \Gamma$ for $\Gamma \ll 1$. This recovers the mean-field energy with a damping that scales linearly with the scattering length as recently observed for the Bose polaron in a box potential [20]. Likewise, a straightforward but more cumbersome analysis yields $\omega \simeq -(4\pi n)^{2/3} - 2i\gamma (4\pi)^{1/3} n^{2/3}/3$ for $1/a \rightarrow 0$, which also agrees with the observed $n^{2/3}$ scaling of the damping at unitarity [20].

Spectral properties. We now explore the spectral properties of a zero-momentum impurity as predicted by our phenomenological model, focusing on comparing with the recent experiment [20]. Figure 2(a) plots the spectral function of the impurity obtained numerically using $\Gamma = 0.12$ and $\gamma = 2.2$ with the experimentally realistic value $n^{1/3} a_B = 0.005$. We clearly see a broad spectral response with a large linewidth in the vicinity of the unitary limit. Following Ref. [20], we extract the position E_p and width Γ_{pol} of the signal, using a Lorentzian fit for the attractive polaronic branch ($E_p < 0$) for $a < 0$ and the repulsive branch ($E_p > 0$) for $a > 0$, and Gaussian fit for the attractive branch ($E_p < 0$) for $a > 0$ [53]. The obtained widths are plotted in Fig. 2(b). We observe reasonable agreement with the experimental results of Ref. [20]

plotted as yellow dots, regarding both the peak positions of the branches and their widths. A better agreement may be achieved by performing a global fit to determine the parameters γ and Γ .

Both the experimental and theoretical data exhibit a shift of the attractive-branch maximum toward the molecular energy for $a > 0$ as well as the asymmetric behavior of the widths of the attractive and repulsive branches, which is a manifestation of different physics at play. Indeed, the repulsive polaron is inherently metastable even for $\gamma = \Gamma = 0$ due to the presence of a low-lying molecular state [54]; therefore, its width is weakly influenced by the complex interaction. It is worth noting that the agreement between the experimental and theoretical results presented in Fig. 2 can be further improved by incorporating broadening mechanisms specific to this experiment, like magnetic field fluctuations, three-body losses, and inhomogeneous density into our theoretical model. Including these mechanisms would increase the estimated polaron width and result in larger error bars for the peak position. Consistent with these trends, a recent solvable-model analysis reproduces central features of the observed spectra and provides a complementary microscopic perspective on the damping systematics [55].

Dynamical properties. The dynamical properties of impurities and polaron formation can be probed using Ramsey interferometry, which measures the overlap $A(t) = \langle \text{BEC} | \hat{c}_{\mathbf{k}=0}(t) \hat{c}_{\mathbf{k}=0}^\dagger(0) | \text{BEC} \rangle$ between the state created by injecting a zero-momentum impurity at time $t = 0$ and taking it out again at time t [56–58]. This technique was pioneered to explore Fermi polarons [15,16] and subsequently modified to explore Bose polarons [17].

Such experiments correspond to measuring $A(t)$ in Eq. (2) with an initial “bare” wave function for which $A(t = 0) = 1$ and $B_{\mathbf{k}}(t = 0) = 0$. The impurity is subsequently dressed by excitations in the BEC resulting in $|A(t)| < 1$ and $B(\mathbf{k}, 0) \neq 0$ [59]. We compare the results with a recent experiment measuring $A(t)$ for $a > 0$ [19]. The results for $a < 0$ are presented in the Supplemental Material [36].

In Fig. 3, we present our calculations of the contrast $|A(t)|$ using different methods and compare them with the experimental data [19]. First, we see that the ladder approximation with a real interaction strength captures the initial decay of $|A(t)|$ for the present case just as well as for the attractive case $a < 0$ [17]. It also predicts oscillations due to the interference between the repulsive and attractive polaron branches [19] (not present for $a < 0$), which, however, are much more pronounced than what is observed experimentally even after trap averaging and magnetic field fluctuations are included [19]. The results based upon the ansatz from Eq. (2) are closer to the measured values featuring a quicker decay of $|A(t)|$. This demonstrates that a more accurate description of the attractive high-energy branch and its decay is important for describing the observed dynamics [60]. Still, this improved description does not fully agree with the experimental results where the decay of oscillations is even more pronounced.

Figure 3 shows that our phenomenological model with $\Gamma = 0.05$, $\gamma = 0.5$ captures the experimental dynamics accurately within experimental uncertainties [61], including the overall decrease of $|A(t)|$ as well as the superimposed damped oscillations (data for other values of $k_n a$ are shown in Ref.

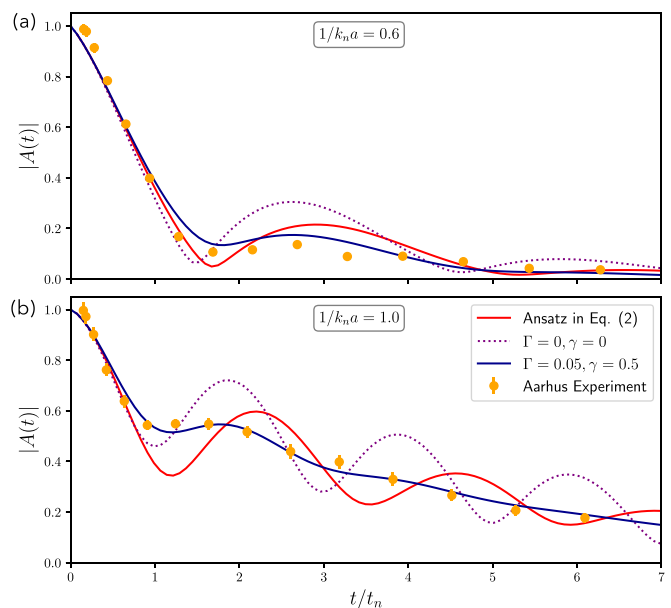


FIG. 3. Time dynamics of the contrast $|A(t)|$ on the repulsive side of the Feshbach resonance for $(k_n a)^{-1} = 0.6$ in panel (a) and $(k_n a)^{-1} = 1.0$ in panel (b). Dotted purple curves correspond to the standard ladder approximation ($\gamma = \Gamma = 0$); solid blue curves are calculated using our phenomenological model with $\Gamma = 0.05$, $\gamma = 0.5$; red solid curves are from the ansatz in Eq. (2) in the two-channel Hamiltonian (1) with $n^{1/3} a_- = -100$, $n^{1/3} a_B = 0.005$ similar to the experiment [19]. The orange dots represent experimental data [19]. We account for magnetic dephasing and three-body loss when computing the blue and red solid curves.

[36]). Furthermore, we find that parameter values in the ranges $\Gamma \simeq 0.05\text{--}0.1$ and $\gamma \simeq 0.5\text{--}2$ actually capture both the Cambridge [20] and the Aarhus experiments [19]. The fact that our phenomenological model can describe two independent experiments with consistent values of parameters indicates that it captures the most important physics. The differences in the optimal values of γ and Γ for the two experiments can be partially explained by the differences in the trapping potentials since the value of Γ_{pol} may be influenced by the inhomogeneity of the trap [62]. Further, we account for magnetic decoherence and three-body losses only when analyzing the Aarhus experiment following the procedure described in Ref. [19]. Consequently, larger values of Γ and γ are required for the Cambridge experiment to accommodate these additional sources of broadening.

To enable the extraction of the experiment independent values of Γ and γ , one would need to exclude the broadening mechanisms specific to a given experimental setup. Such a systematic analysis is beyond the scope of the present work and is left for future studies.

Summary. We investigated a phenomenological model of Bose polaron states, focusing on scenarios where the impurity is primarily correlated with a single boson, as these configurations couple most strongly to the two most commonly used experimental probes. To describe the possible decay of these states to lower-lying states that involve correlations with more bosons, we introduced an imaginary part to the impurity-boson interaction. Using a simple variational wave

function, we then showed that this model can describe two recent measurements of the spectral and dynamical properties of the Bose polaron, including both the energies and the decay rates.

Our findings pave the way for deeper investigations into the limits of the quasiparticle description in the Bose polaron problem. First, our phenomenological approach offers a practical framework for parametrizing experimental data, thereby enabling meaningful comparisons across different experiments. Second, our findings underscore the need for microscopic theories of Bose polarons that include the decay of the experimentally relevant polaron states. This entails analyzing the complex problem of few-body correlations and bound states within a many-body environment, which we expect to be most pronounced for light impurities. It would also be interesting to apply our model to systems with different masses and measurement protocols such as ejection spectroscopy [10,12]. Finally, the effects of a nonzero temperature

[57,58,63–67] as well as of a lattice [68–71] on the decay of the Bose polaron should be explored.

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Data availability. The data that support the findings of this article are openly available [72].

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- $$L(\omega) = A \frac{\Gamma^2}{(\omega - \omega_0)^2 + \Gamma^2} + C$$
- and define the peak position and width as $E_p \equiv \omega_0$, $\Gamma_{\text{pol}} \equiv \Gamma$ (HWHM). For the attractive branch at $a > 0$, consistent with Ref. [20], we use a Gaussian profile
- $$G(\omega) = A e^{-(\omega - \omega_0)^2 / (2\sigma^2)} + C$$
- and report $E_p \equiv \omega_0$, $\Gamma_{\text{pol}} \equiv \sigma$. Fits are performed with non-linear least squares, and initial guesses are set by the local maximum of $A(\omega)$.
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- [59] It is easy to understand the time dynamics of $A(t)$ for a noninteracting Bose gas ($a_B = 0$) using the following ansatz for $A(t)$, which fits the numerical data well [36]:
- $$A(t) = \sqrt{Z}e^{-i\omega_1 t} + (1 - \sqrt{Z})e^{-i\omega_2 t}, \quad (6)$$
- where $|Z| = 8\pi na(12\pi na - \omega)^{-1}$ is the residue and ω_1 (ω_2) is the solution of Eq. (4) with $\text{Im}(\omega) < 0$; in the limit $\Gamma = 0$, ω_1 (ω_2) is the standard attractive (repulsive) polaron [12]. The coefficients B read
- $$B_{\mathbf{k}}(t) = \frac{\omega_1 \sqrt{Z} f_{\mathbf{k}}(\omega_1)}{(2\pi)^{3/2}(\omega_1 - 2\epsilon_{\mathbf{k}})} + \frac{\omega_2 (1 - \sqrt{Z}) f_{\mathbf{k}}(\omega_2)}{(2\pi)^{3/2}(\omega_2 - 2\epsilon_{\mathbf{k}})}, \quad (7)$$
- where $f_{\mathbf{k}}(\omega) = e^{-i\omega t} - e^{-2i\epsilon_{\mathbf{k}} t}$. The first term in $f_{\mathbf{k}}$ is due to the pole, and the second one is due to the bosonic continuum. The latter contribution describes the particles that move away from the impurity and do not interact with it at later times. For an interacting Bose gas, this would correspond to sound waves emitted by the impurity (see, e.g., Ref. [73] for a further visualization of this process).
- [60] Note that our phenomenological model agrees with the variational ansatz in Eq. (2) regarding the dynamics when $\Gamma = 0.01$ and $\gamma = 0.1$.
- [61] For example, for the data shown in Fig. 3(a), a weighted goodness-of-fit analysis yields a reduced $\chi^2 = 3.15$ and a Pearson correlation coefficient of 0.996, confirming that the theoretical model reproduces the experimental trend accurately. The somewhat elevated χ^2 value primarily reflects the very small experimental uncertainties ($\approx 1\%$), which do not include systematic effects such as magnetic field jittering and amplitude calibration. When a realistic systematic uncertainty of a few percent is taken into account, the reduced χ^2 approaches unity, indicating statistical consistency.
- [62] Note that we performed a trap averaging for Ref. [19] (see Ref. [36]) and concluded that it has a weak effect on the time dynamics as the decoherence dynamics mainly happens at shorter timescales. It is known, however, that the presence of the trap broadens the polaron peak in the spectroscopic experiments [11].
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