Optical Signals to Monitor the Dynamics of Phonon-Modified Rabi Oscillations in a Quantum Dot

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1. Introduction

Semiconductor quantum dots (QDs) with their discrete energy structure are perfectly suited to study light–matter interaction effects known from quantum optics with atoms in a solid-state platform. Several phenomena known from cavity quantum electrodynamics have already been observed in QD structures, among them the Vacuum Rabi splitting, the Mollow triplet as well as the generation of single and entangled photons. Due to the outstanding photon properties, QDs are promising candidates to be used as photon sources in quantum information technology and photonic circuits.

Unlike atoms, QDs are solid-state objects and therefore subject to the electron–phonon interaction. For a long time phonons were just seen as an unwanted problem which hinders effects known from atomic physics like damping of Rabi oscillations or reduction of the photon indistinguishability. As such, it has been desirable to reach the decoupling regime of the electron–phonon interaction, where a reappearance of the signals occurs.

Recently, instead of avoiding them, phonons are explicitly exploited, because they allow for new phenomena and possibilities. A prominent example is the phonon-assisted state preparation, which has also been shown to result in a high quality photon emission. Key to the phonon-assisted state preparation is using an off-resonant laser pulse to excite the QD above the exciton transition and exploiting the phonon relaxation to the lower dressed state as well as the adiabatic undressing.

In this paper, we study the phonon influence of the time-resolved optical signals from a constantly driven QD exhibiting Rabi oscillations for detuned excitation. We will show that the evolution of the probe signals reflects well the phonon-mediated relaxation into the lower dressed state, which results in different occupations of the bare QD states depending on the sign of detuning.

Another well established optical signal of a constantly driven two-level system is the resonance fluorescence signal with the well-known Mollow triplet consisting of a strong center peak and two weaker side peaks at the Rabi frequency. The Mollow triplet has also been measured for QDs. Driving with two laser pulses, the peaks split up even further resulting in a complex optical spectrum. The resonance fluorescence signal is typically evaluated using the quantum regression theorem. Here we will consider an optical pump-probe set-up, which will allow us to also monitor the dynamics of the system during the Rabi oscillations. A sketch of the system is shown in Figure 1. In the time-resolved dynamics we find an oscillation between absorptive and dispersive behavior, a typical signal for the occurrence of superposition states. We expect that our theoretical predictions can be measured in heterodyne detection set-ups used to measure four-wave-mixing signals of single QDs or in differential transmission experiments available for single QDs. Being able to provide analytical equations, we give an illustrative and insightful view on the corresponding dynamics and optical signals.

2. System and Theoretical Background

2.1. Phonon Free Case

We take the standard model of a QD as two-level system coupled to the light field in dipole and rotating wave approximation such
Rabi frequency respectively. The coupling to the electric field is given by the occupation of the excited state \( \rho \). Because we consider a constant driving pulse the diagonalization yields the dressed states \( \Omega \). The Hamiltonian consists of the free phonon part with the dispersion \( \omega_q \) and the interaction Hamiltonian with the coupling matrix element \( g_q \). Here we will consider only the interaction with longitudinal acoustic phonons, which have been identified as major source of decoherence in QDs. As typical for the electron-phonon interaction, it is of pure dephasing type in the bare state basis, that is, unlike the electric field it cannot induce transitions between the two states. In the dressed state basis, the Hamiltonian is transformed to

\[
H_{ph}^{DS} = \sum_q g_q (b_q^i + b_q^\dagger) \frac{1}{2\Omega} \times \left[ (\Omega - \Delta_q)|u\rangle\langle u| + (\Omega + \Delta_q)|l\rangle\langle l| - \Omega g_q |l\rangle\langle u| + |u\rangle\langle l| \right]
\]

This reveals that the phonon interaction results in two effects: On the one hand there is a renormalization effect (term \( \sim |u\rangle\langle u| \)) and \( \sim 0\langle l|l\rangle \) and on the other hand transitions between the two states can be induced (term \( \sim |u\rangle\langle l| + |l\rangle\langle u| \)). While in the bare state basis, the phonon effects result in a pure dephasing mechanism, in the dressed state basis, we have both a pure dephasing part and a transition part. In calculations using the polaron transformation, therefore, only the renormalization terms are referred to as pure dephasing terms. Note that pure dephasing Hamiltonians typically cannot be solved analytically, because they cause non-Markovian effects. However, here we will focus on the transitions between the dressed states, which in good approximation can be treated within a rate equation approach.

To use the rate approach, we need to estimate the transition rate between the dressed states, which is given by

\[
\gamma = \frac{1}{\tau_{ph}} = \frac{\pi}{2} \left( \frac{\Omega g_q}{\Omega} \right)^2 J(\Omega)
\]

with the phonon spectral density \( J(\omega) \). The phonon spectral density depends on the geometry of the QD and for the electronic properties it is sufficient to assume a spherical symmetric QD. A good approximation of the phonon spectral density is given by a simple model using the cut-off frequency \( \omega_c \) and

\[
J(\omega) = A \omega^3 \exp \left( -\frac{\omega^2}{\omega_c^2} \right)
\]
with \( A \) being a constant. We stress that the relaxation rate \( \gamma \) depends strongly on the driving of the system due to the phonon spectral density, but also on the mixing of the dressed states. For our calculations we use for illustrative reasons a phonon coupling using the parameters \( \alpha_t = 2.5 \text{ ps}^{-1} \) and \( A = 1/(2 \pi \cdot 14.5) \text{ ps}^2 \), which are close to the parameters of a strongly confined GaAs QD. For the calculations we set \( \Omega_r = 3 \text{ ps}^{-1} \) and the temperature to zero.

### 3. Dynamics

The dynamics of an optically driven QD without phonons is well known.\([16]\) Considering the case of a laser field that is switched on instantaneously and the initial Bloch vector \((0, 0, -1)\), this corresponds to the two-level dynamics of the Bloch vector

\[
\begin{bmatrix}
\rho_{11}^\text{DS} \\
\rho_{12}^\text{DS} \\
\rho_{22}^\text{DS}
\end{bmatrix} = \begin{bmatrix}
-2 \frac{\Omega}{\Omega_t} \sin^2 \left( \frac{\Omega}{2} t \right) \\
\frac{\Omega}{\Omega_t} \sin(\Omega t) \\
\frac{\Omega^2}{\Omega_t^2} \sin^2 \left( \frac{\Omega}{2} t \right) - 1
\end{bmatrix}
\]  

(10)

The superindex \( s \) here indicates the phonon-free case. The solution are Rabi oscillations with the period \( T = 2\pi/\Omega \). Only for resonant excitation, that is, for \( \Delta_n = 0 \), the amplitude of the oscillation is one, while it decays rapidly for higher detunings. In terms of state preparation in QDs, Rabi oscillations are thus only useful when used in resonance. However, in QDs the electron–phonon interaction is present, which leads to a damping of the Rabi oscillations for detuned pulses.\([16]\) This phenomenon has been used experimentally.\([28–30]\) As a similar sign dependence has been found for detuning an oscillatory behavior emerges, which in the resonant case and without phonons has an amplitude of 1. We remind that without phonons the amplitude of the oscillations reduces drastically when a finite detuning is introduced. When including phonons a quite different behavior emerges and a clear asymmetry due to the phonon relaxation depending on the sign of the detuning is visible. For the resonant case, the Rabi oscillations are damped and the stationary inversion reaches 0, that is, the exciton occupation is one half. For negative detuning, the stationary value of the Rabi oscillations clearly is below 0, that is, the exciton occupation is rather low. However, for positive detuning, the system can reach stationary values of well above 0 and actually arbitrarily close to 1.\([55]\) This phenomenon has been used in phonon-assisted state preparation, which has been also been seen experimentally.\([28–30]\) A similar sign dependence has been found for excitation with chirped laser pulses, where a phonon dephasing takes place for one sign of the chirp and is unaffected for the other sign of chirp.\([24,64–66]\) When the detuning becomes too strong (in our case \( \Delta_n > 5 \text{ ps}^{-1} \)) the stationary inversion \( r_s \) goes to \(-1\), because here the Bloch oscillations are strongly off-resonant and the phonons are not active, because the splitting between the dressed states lies beyond the phonon spectral density.\([21]\)

Let us briefly comment on the validity of these analytical expressions in comparison to more sophisticated methods to calculate the influence of phonons on the QD dynamics. These rate equations describe the electronic dynamics of the QD system well, which was also confirmed by comparison with a correlation expansion approach.\([39]\) However, due to the renormalization terms in Equation (7) a slight change of the Rabi frequency occurs, which is not captured by the rate equations, but can be accounted for by an adjustment of the Rabi frequency compared to the case without phonons. As such effects are of non-Markovian nature,\([47]\) a more evolved theory is needed like a polaron transformation,\([46,55]\) a correlation expansion\([50]\) or a path integral method\([56,57]\) to accurately the energy renormalizations. Also in the case of strong electron–phonon coupling as it occurs in other materials like GaN, the phonon effects can become so strong, that approximate approaches like the correlation expansion or the rate equations, might break down.\([34]\)

Because we neglect the non-Markovian effects, we also do not describe the acoustic phonon sideband, which in GaAs dots is about three orders of magnitudes below the Mollow triplet peaks.\([58,59]\) Also for more evolved photon properties\([60,61]\) or entangled photons,\([62,63]\) the rate equations might fail depending on the parameter regime. Nonetheless, to describe the quantum dynamics of the electronic states in a QD, rate equations capture the essential features given that a non-monotonous phonon spectral density is taken into account and by their simplicity offer beautiful insight in the underlying physics.

In Figure 2a we plot the inversion \( r_s \) calculated by Equation (12c) as function of time and detuning. For each detuning an oscillatory behavior emerges, which in the resonant case and without phonons has an amplitude of 1. We remind that without phonons the amplitude of the oscillations reduces drastically when a finite detuning is introduced. When including phonons a quite different behavior emerges and a clear asymmetry due to the phonon relaxation depending on the sign of the detuning is visible. For the resonant case, the Rabi oscillations are damped and the stationary inversion reaches 0, that is, the exciton occupation is one half. For negative detuning, the stationary value of the Rabi oscillations clearly is below 0, that is, the exciton occupation is rather low. However, for positive detuning, the system can reach stationary values of well above 0 and actually arbitrarily close to 1.\([55]\) This phenomenon has been used in phonon-assisted state preparation, which has been also been seen experimentally.\([28–30]\) A similar sign dependence has been found for excitation with chirped laser pulses, where a phonon dephasing takes place for one sign of the chirp and is unaffected for the other sign of chirp.\([24,64–66]\) When the detuning becomes too strong (in our case \( \Delta_n > 5 \text{ ps}^{-1} \)) the stationary inversion \( r_s \) goes to \(-1\), because here the Bloch oscillations are strongly off-resonant and the phonons are not active, because the splitting between the dressed states lies beyond the phonon spectral density.\([21]\)
traced back to the oscillation of the Bloch vector being either in the left or right part of the Bloch sphere.

4. Optical Signals

4.1. Theoretical Description

Having seen that there is a strong difference in the occupation dynamics between positive and negative detuning due to the electron–phonon interaction, the questions arises, whether this is also visible in optical spectra. To calculate the optical spectra, we assume a pump-probe set-up, where the pump pulse is a continuous wave (cw) beam, which is switched on instantaneously at \( t = 0 \) with the Rabi frequency \( \Omega_R \) and the detuning \( \Delta \omega \). A second, ultrafast pulse (the probe pulse) is then used to monitor the dynamics of the system. The second pulse has its pulse maximum at time \( t = \tau \) being the time delay to the set-in of the cw light. We model the ultrafast probe pulse by a \( \delta \)-pulse, such that the total coupling to the electric field is given by

\[
\tilde{\Omega}(t) = \Omega_b \delta(t) + \Omega_{pr} e^{i\varphi} \delta(t - \tau)
\]  

with the pump pulse \( \Omega_b \delta(t) \) and the probe pulse having the pulse area \( \Omega_{pr} \). For the probe pulse a phase \( e^{i\varphi} \) is added. To calculate the probe signal we use a matrix multiplication method discussed in ref. [39, 51, 67]. The method is based on the assumption that the influence of the phonons can be neglected during a ultrafast pulse. To calculate the action of the \( \delta \)-pulse, the equation of motion is rewritten as function of the pulse area \( \Omega_{pr} \). The solution of the corresponding equation can be gained by transforming the density matrix before the probe pulse \( \rho_{\text{before}} \) with the matrix \( S \) to calculate the density matrix \( \rho_{\text{after}} \) after action of the probe pulse

\[
\rho_{\text{after}} = S^{-1} \rho_{\text{before}} S \quad \text{with}
\]

\[
S = \begin{pmatrix}
\cos(\frac{1}{2} \Omega_{pr}) & i e^{i\varphi} \sin(\frac{1}{2} \Omega_{pr}) \\
 ie^{-i\varphi} \sin(\frac{1}{2} \Omega_{pr}) & \cos(\frac{1}{2} \Omega_{pr})
\end{pmatrix}
\]  

(14)

The density matrix \( \rho_{\text{before}} \) is the given by Equation (12) at time \( t = \tau \). We then apply Equation (14) to gain \( \rho_{\text{after}} \), which is then the new starting point for the evolution of the density matrix according to Equation (11). The solution in the dressed state basis can be obtained by simple integration. The resulting density matrix in the bare states and in particular the polarization \( \rho_{\text{pr}}(t; \tau, \varphi) \) is a function of \( \tau \) and depends parametrically on \( \tau \). It further contains terms going with different phases \( e^{i\varphi n} \), \( n \in \{0, \pm1, \pm2\} \). These different phases correspond to different optical signals, for example, the term for \( n = 1 \) correspond to the probe signal and terms for \( n = 2 \) describe four wave mixing signals. The phase filtering is directly related to heterodyne detection techniques, where different frequencies for the pulses are used and then a phase filtering is performed. In experiments using extended structures, the phases correspond to different directions of the pulse propagation. We gain the probe signal by extracting all terms proportional to \( e^{i\varphi} \) of the polarization, that is, \( \rho_{\text{pr}}(t; \tau, \varphi) = \rho_{\text{pr}}(t; \tau, \varphi)_{\varphi = \varphi_0} \). To calculate the spectrum we take the imaginary part of the Fourier transform of the probe polarization

\[
a_{\text{pr}}(\omega; \tau) \sim \text{Im} \left\{ \mathcal{F} \mathcal{T} \left[ \rho_{\text{pr}}(t; \tau) \right] \right\}
\]  

(15)
to obtain the probe spectrum as

$$a(\omega; \tau) \sim -\frac{\Omega_g}{4\Omega_r^2} \left[ -\Delta_\omega r_\omega(\tau) + \Omega_r r_\omega(\tau) \right] \frac{\gamma - \Delta_\omega r_\omega(\tau) \omega}{\gamma^2 + \omega^2} \right.$$

$$\left (+ \frac{(\Omega - \Delta_\omega)}{8\Omega_r^2} \frac{[\Omega_r r_\omega(\tau) - (\Omega - \Delta_\omega ) r_\omega(\tau)]}{(\gamma/2)^2 + (\omega - \Omega)^2} \right) \right.$$

$$\left (\frac{[\Omega_r r_\omega(\tau) + (\Omega + \Delta_\omega ) r_\omega(\tau)]}{(\gamma/2)^2 + (\omega + \Omega)^2} \right) \right)$$

This equation shows the emergence of three features at $\omega = 0$ and $\omega = \pm \Omega$. Each feature is composed of two functions, a Lorentzian function (Cauchy distribution) of the form $g/(g^2 + \omega^2)$ and a function of the form $\alpha/\left( g^2 + \omega^2 \right)$, which we will call dispersive line shape. We will now analyze the details of this equation while showing the corresponding figures below.

### 4.2. Time-Dependence of the Probe Signals

Figure 3 shows the probe spectra as a function of delay $\tau$ for positive detuning $\Delta_\omega = +3 \text{ps}^{-1}$ (left) and negative detuning $\Delta_\omega = -3 \text{ps}^{-1}$ (right). In the bottom panel we show the spectra for short delays up to $\tau = 6 \text{ ps}$, that is, shortly after the cw-pulse has been switched on, for better visualization. Overall we find for both detunings an oscillatory behavior as function of time delay, which reflects the Rabi oscillations occurring in the system. On the short time scales the phonon relaxation is not effective and the two spectra are mirrored by each other. For positive detuning (Figure 3c) we see a strong peak at $\omega = -\Omega$, while a weak peak occurs at $\omega = \Omega$. Additionally there is a peak at $\omega = 0$. The position of the peak corresponds to the energy splitting between the dressed state energies, as known from the Mollow triplet.\(^{34}\)

It is instructive to analyze the time behavior for short time delays in more detail. For this we show spectra for four different detunings $\Delta_\omega = 0, 1, 2, 3 \text{ ps}^{-1}$ within the first oscillation period in Figure 4. Note that in the first oscillation period the spectra with opposite sign, that is, $\Delta_\omega = -1, -2, -3 \text{ ps}^{-1}$, can be obtained by just taking $\omega \rightarrow -\omega$. The spectra are shown for $\tau = 0, T/4, T/2, 3T/4$. During the first period the relaxation induced by phonons is negligible and we can treat the Bloch vector as in the case without phonons. However, for the probe signals, the phonons results in the decay of the probe polarization. Different to a simple decay, which would result in the typical 2:1 amplitude of the lines found in the Mollow triplet, the decay of the phonons results in different width of the middle and side peak with a ratio of 2:1.\(^{39}\)

Both Figures 3 and 4 show that the spectra oscillate between absorption/gain and dispersive behavior. The absorptive behavior occurs whenever the system is in a minimum or maximum of the inversion $r_\omega$, here at $\tau = 0$ or $T/2$, respectively. At $\tau = 0$ (Figure 4a) all optical signals independent of the detuning are positive, while at $\tau = T/2$ (Figure 4c) we find negative and positive peaks. All peaks have a purely Lorentzian shape. To understand our finding, we inspect Equation (16). For the resonant case $\Delta_\omega = 0$, it is clear that for $\tau = 0, T/2$ the coherence is zero, that is, $r_\omega = 0$, hence no dispersive line shapes appear. For detuned excitation the imaginary part of the coherence is also zero $r_\omega(0) = r_\omega(T/2) = 0$, but at the maximum a finite real part occurs with $r_\omega(T/2) = -2\Delta_\omega \Omega_\omega^2/\Omega_r^2$. Nonetheless, this results in only Lorentzian lineshapes.

The signs of the Lorentzian are connected to absorption (positive sign) and gain (negative sign). Hence it is expected, that whenever the system is in the ground state with $r_\omega = -1$, the spectrum shows only absorption, that is, $a(\omega) > 0$ for all $\omega$. This is also in agreement with Equation (16). The situation is more complicated when the system is in a maximum of the occupation at $\tau = T/2$. In the resonant case at $\tau = T/2$ all peaks are negative, because here the system is fully inverted. When we increase the detuning, the peaks at $\omega = \Omega$ and $\omega = 0$ stay...
For increasing $\Delta$, the side of the peak. With increasing detuning, the Lorentzian $\omega = \Omega$, looking closely, because the line does not cross zero at characteristic zero crossing. The Lorentzian part becomes visible the feature is dominated by the dispersive line shape with the Equation (16b)) the absorptive part remains rather small, hence the real part of the coherence weighted by the corresponding prefactor. It is already remarkable that the coherence can play such a part in optical signals.

At all other time delays $\tau$ not being a multiple of $T/2$, we also have a non-vanishing imaginary part of the coherence $r_y$. According to Equation (16), this results in dispersive line shapes. The imaginary part has its maximal strength at $\omega = 0$. For resonant excitation the inversion is zero at $\tau = T/4$ and $3T/4$, hence no feature is seen in the spectrum. For increasing $\Delta_\omega$, it holds that $[\Delta_\omega, r_y(T/4) + \Omega_y r_y(T/4)] = 0$ (cf. Equation (16a)) and only the imaginary part of the coherence $r_y$ contributes to the spectra. Accordingly, we see always a purely dispersive feature at $\omega = 0$ for all detunings.

At $\tau = T/4$ the side feature at $\omega = \pm \Omega$ comprises a mixture of Lorentzian and dispersive line shape for a finite detuning, due to the remaining finite value of $r_y(T/4)$ and the non-zero inversion $r_s(T/4)$. The respective strength of Lorentzian and dispersive part depend on the sign of the detuning as seen in Equation (16). For positive detuning, in the feature at $\omega = \Omega$ (cf. Equation (16b)) the absorptive part remains rather small, hence the feature is dominated by the dispersive line shape with the characteristic zero crossing. The Lorentzian part becomes visible when looking closely, because the line does not cross zero at $\omega = \Omega$, but slightly beneath. For positive detuning the feature at $\omega = -\Omega$ (cf. Equation (16c)) is dominated by a Lorentzian. Here the dispersive behavior is only visible by a slight negativity at the side of the peak. With increasing detuning, the Lorentzian part becomes stronger, such that it dominates the feature. At $\tau = 3T/4$, the overall behavior is similar, only the sign of the dispersive feature has changed, while the Lorentzian features remain the same. This reflects well the oscillation of the coherence $r_y$ and $r_s$, where we find that $r_s(T/4) = r_s(3T/4)$, while for the imaginary part it holds that $r_y(T/4) = -r_y(3T/4)$.

We stress that dispersive signatures in probe signals are a sign of a quantum mechanical superposition state with complex phase or in other words a direct influence of the imaginary coherence between the states. Besides the optically driven QD, where a superposition between ground and excited state occurs, such dispersive line shapes have been predicted theoretically for the probe spectra of QDs doped with a single manganese atom, where a superposition of bright and dark exciton can be prepared. Experimentally, an oscillatory behavior between absorptive/dispersive line shapes has been seen in a superposition of trion states in excellent agreement with theory.

The influence of the phonons on the dynamics of the occupation and polarization of the two-level system (cf. Figure 2) becomes visible in the probe spectra for longer time delays $\tau$ as shown in Figure 3 (upper panels). Note that in both cases the system relaxes into the lower dressed state, which due to its composition (cf. Equation (5)) has different weights of the bare states resulting in the different dynamics of the inversion as shown in Figure 2. For both detunings we find that the amplitude of the middle feature at $\omega = 0$ decays in strength, while the side features show a different behavior depending on the detuning. For negative detuning shown in Figure 3b, we find that the strong peak at $\omega = \Omega$ remains in its strength, while only the oscillatory behavior vanishes. In contrast, for positive detuning the strong feature at $\omega = -\Omega$ first reduces in amplitude and then reverses its sign to show a strong gain feature. Let us briefly comment on the time scale of the relaxation, which for our parameters takes place within a few tens of ps, which is typical for a QD coupled to phonons. The radiative life time of such QDs is typically on the order of a ns. In this case the radiative lifetime can be safely neglected. However, for some QDs in particular in photonic structure, the radiative lifetime can be drastically shorter such that

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{figure4.png}
\caption{Probe signals for different detunings $\Delta_\omega = 0, 1, 2, 3 \text{ ps}^{-1}$ at time delays a) $\tau = 0$, b) $\tau = T/4$, c) $\tau = T/2$, and d) $\tau = 3T/4$.}
\end{figure}
a competing process between phonon relaxation and radiative decay can take place.

The behavior of the spectrum for long time delays as function of detuning $\Delta_\omega$ is summarized in Figure 5 together with the spectrum for zero time delay. The probe spectra in these cases can be evaluated to

$$\alpha(\omega, \tau = 0) \sim \frac{\Omega_\omega^2}{4\Omega^2 \gamma^2 + \omega^2} \frac{\gamma}{8\Omega^2} \frac{(\Omega - \Delta_\omega)^2}{(\gamma/2)^2 + (\omega - \Omega)^2}$$

$$+ \frac{(\Omega + \Delta_\omega)^2}{8\Omega^2} \frac{\gamma}{(\gamma/2)^2 + (\omega + \Omega)^2}$$

(17)

$$\alpha(\omega, \tau \to \infty) \sim \frac{\Omega_\omega^2}{4\Omega^2 \gamma^2 + \omega^2} \frac{\gamma}{8\Omega^2} \frac{(\Omega - \Delta_\omega)^2}{(\gamma/2)^2 + (\omega - \Omega)^2}$$

$$- \frac{(\Omega + \Delta_\omega)^2}{8\Omega^2} \frac{\gamma}{(\gamma/2)^2 + (\omega + \Omega)^2}$$

(18)

For zero time delay the influence of the phonons is only visible in the widths of the side peaks compared to the center peak as discussed in ref. [39]. Because the system is in the ground state at $\tau = 0$ all peaks are positive as seen above. In the limit $\tau \to \infty$ the equations confirm the vanishing of the center feature at $\omega = 0$, while the side peaks remain. If $\Delta_\omega \gg \Omega_\omega$ we can approximate $\Omega - \Delta_\omega \to 0$ and $\Omega + \Delta_\omega \to 2\Delta_\omega$, underlining that depending on the sign of $\Delta_\omega$ one peak becomes strong, while the other vanishes, as we also see in Figure 5 for increasing detuning. Figure 5 further summarizes nicely that for all detunings a relaxation into the lower dressed state takes place with the peak at negative frequencies having a negative amplitude and the peak at positive frequencies having a positive amplitude.

5. Conclusion

In conclusion, we have presented an analytical description of the probe spectra of a continuously driven QD under the influence of phonons. The probe spectra were characterized by two parts: i) Lorentzians with positive and negative sign corresponding to absorption and gain and ii) dispersive lines, which could be traced back to the imaginary part of the coherence between states as part of the superposition. We accounted for the relaxation of the phonons into the lower dressed state, which results in a strong dependence of the occupation dynamics on the sign of detuning. The dynamics is reflected by a changing peak amplitude occurring at the Rabi frequency. All features are well described using analytical equations, which gives an insightful understanding of the interplay of occupations and coherences in QD dynamics. With this, our results propose a way to measure quantum mechanical superposition states in solid state few-level systems like QDs.

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Conflict of Interest

The authors declare no conflict of interest.

Data Availability Statement

Research data are not shared.

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