Measurements with a noninvasive detector and dephasing mechanism

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We study the dynamics of the measurement process in quantum-dot systems, where a particular state out of coherent superposition is observed. The ballistic point contact placed near one of the dots is taken as a noninvasive detector. We demonstrate that the measurement process is fully described by the Bloch-type equations applied to the whole system. These equations clearly reproduce the collapse of the density matrix into the statistical mixture in the course of the measurement process. The corresponding dephasing width is uniquely defined. We show that the continuous observation of one of the states in a coherent superposition may *accelerate* decay from this state—in contradiction with rapidly repeated observations, which slow down the transitions between quantum states (the quantum Zeno effect). [S0163-1829(97)04748-6]

I. INTRODUCTION

In recent years there have been many measurements in mesoscopic systems sensitive to the phase of the electronic wave function. We mention experiments with double-split systems,[1,2] quantum dot embedded in an Aharonov-Bohm ring,[3,4] and coupled quantum dots.[5] It is known that the phase of the wave function, or, more precisely, the off-diagonal density-matrix elements, can be destroyed by interaction with the environment, or with the measurement device. As a result, the density matrix becomes the statistical mixture. The latter does not display any coherence effects. Now the rapid progress in microfabrication technology allows us to investigate experimentally the dephasing process in mesoscopic systems, for instance, by observation of a particular state out of coherent superposition.[6]

Although dephasing (decoherence) plays an important role in different processes, its mechanism has not been elaborated upon enough. For instance, in many studies of the quantum measurement problems the dephasing is usually accounted for by introducing some phenomenological dissipating terms, associated with a detector (or an environment). Yet such a procedure cannot not illuminate the origin of the dephasing and its role in the measurement problem. The most appropriate way to approach the problem, however, is to start with a microscopic description of the measured system and the detector together with use of the Schrödinger equation, $i\dot{\sigma}=[\mathcal{H},\sigma]$, where $\sigma(\mathcal{S},\mathcal{S}';\mathcal{D},\mathcal{D}',t)$ is the total density matrix and $\mathcal{H}$ is the Hamiltonian for the entire system. Here $\mathcal{S}(\mathcal{S}')$ and $\mathcal{D}(\mathcal{D}')$ are the variables of the measured system and the detector, respectively. In this case the influence of the detector on the measured system can be determined by ''tracing out'' the detector variables in the total density matrix,

$$\sum_D \sigma(\mathcal{S},\mathcal{S}',\mathcal{D},\mathcal{D},t)\to\overline{\sigma(\mathcal{S},\mathcal{S}',t)}. \tag{1.1}$$

The decoherence would correspond to an exponential damping of the off-diagonal matrix elements in the reduced density matrix: $\overline{\sigma(\mathcal{S},\mathcal{S}',t)}\sim\exp(-\Gamma_d t)\to 0$ for $\mathcal{S}=\mathcal{S}'$, with Γ_d the decoherence rate.

In this paper we apply the above approach to study the decoherence, generated by measurement of a quantum-dot occupancy in multidot systems. As the measurement device (detector), we take the ballistic point contact in close proximity to the measured quantum dot.[7] Since the quantum-mechanical description of this detector is rather simple, it allows us to investigate the essential physics of the measurement process in great detail. In addition, the ballistic point contact is a noninvasive detector.[7] Indeed, the time which an electron spends inside it is very short. Thus the point contact would not distort the measured dot. (The first measurement of decoherence in a quantum dot generated by the point-contact was recently performed by Buks *et al.*[6])

The plan of this paper is the following: In Sec. II we describe the measurement of a quantum-dot occupation, when the current flows through this dot. We use the quantum rate equations,[8–12] which allow us to describe both the measured quantum dot and the point-contact detector in the most simple way. A detailed microscopic derivation of the rate equations for the point contact is presented in Appendix A. In Sec. III we investigate the decoherence of an electron in a double-well potential caused by the point-contact detector by measuring the occupation of one of the wells. Special attention is paid to a comparison with the result of rapidly repeated measurements. For a description of this system we use Bloch-type rate equations,[8,9,13,14] which are derived in Appendix B. Similar decoherence effects, but in dc current flowing through a coupled-dot system, are discussed in Sec. IV. Section V is a summary.

II. BALLISTIC POINT-CONTACT DETECTOR

Consider the measurement of the electron occupation of a semiconductor quantum dot by means of a separate measuring circuit in close proximity.[6,7] A ballistic one-dimensional point contact is used as a ''detector'' of whether resistance is very sensitive to the electrostatic field generated by an electron occupying the measured quantum dot. Such a setup is shown schematically in Fig. 1, where the detector is represented by a barrier, connected with two reservoirs at the chemical potentials μ_L and μ_R, respectively. The transmission probability of the barrier varies from T to T', depending

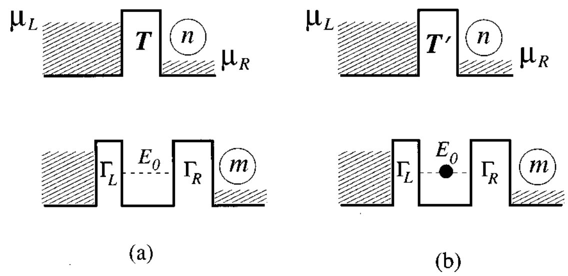

FIG. 1. Ballistic point contact near the quantum dot. $\Gamma_{L,R}$ are the corresponding tunneling rates. The penetration coefficient of the point contact is T for the empty dot (a), and T' for the occupied dot (b). The indices m and n denote the number of electrons penetrating to the right reservoirs at time t.

on whether or not the quantum dot is occupied by an electron, Figs. 1(a) and 1(b).

Initially all the levels in the reservoirs are filled up to the corresponding Fermi energies, and the quantum dot is empty. (For simplicity we consider the reservoirs at zero temperature.) Such a state is not stable, since electrons move from the left to the right reservoir. The time evolution of the entire system can be described by the master (rate) equations[8–12] (the microscopic derivation from the many-body Schrödinger equation is given in Appendix A and in Refs. 8 and 9).

In order to write down these equations we introduce the probabilities $\sigma_{aa}^{m,n}(t)$ and $\sigma_{bb}^{m,n}(t)$ of finding the entire system in the states $|a\rangle$ and $|b\rangle$ corresponding to empty or occupied dots [Figs. 1(a) and 1(b)]. Here m and n are the number of electrons penetrating to the right reservoirs of the measured system and the detector, respectively. The corresponding rate equations for these probabilities have the following forms:

$$\dot{\sigma}_{aa}^{m,n}=-(\Gamma_L+D)\sigma_{aa}^{m,n}+\Gamma_R\sigma_{bb}^{m-1,n}+D\sigma_{aa}^{m,n-1}, \tag{2.1a}$$

$$\dot{\sigma}_{bb}^{m,n}=-(\Gamma_R+D')\sigma_{bb}^{m,n}+\Gamma_L\sigma_{aa}^{m,n}+D'\sigma_{bb}^{m,n-1}, \tag{2.1b}$$

where $\Gamma_{L,R}$ are the transition rates for an electron tunneling from the left reservoir to the dot and from the dot to the right reservoir, respectively, and $D=T(\mu_L-\mu_R)/2\pi$ is the rate of electron hopping from the right to the left reservoir through the point contact (the Landauer formula).

The accumulated charge in the right reservoirs of the detector (d) and of the measured system (s) is given by

$$Q_d(t)=\sum_{m,n} n[\sigma_{aa}^{m,n}(t)+\sigma_{bb}^{m,n}(t)], \tag{2.2a}$$

$$Q_s(t)=\sum_{m,n} m[\sigma_{aa}^{m,n}(t)+\sigma_{bb}^{m,n}(t)]. \tag{2.2b}$$

(We choose the units where the electron charge $e=1$ and $\hbar=1$.) The currents flowing in the detector and in the measured system are $I_d(t)=\dot{Q}_d(t)$ and $I_s(t)=\dot{Q}_s(t)$. Using Eqs. (2.1) and (2.2), we obtain

$$I_d(t)=\sum_{m,n} n[\dot{\sigma}_{aa}^{m,n}(t)+\dot{\sigma}_{bb}^{m,n}(t)]=D\sigma_{aa}(t)+D'\sigma_{bb}(t), \tag{2.3a}$$

$$I_s(t)=\sum_{m,n} m[\dot{\sigma}_{aa}^{m,n}(t)+\dot{\sigma}_{bb}^{m,n}(t)]=\Gamma_R\sigma_{bb}(t), \tag{2.3b}$$

where $\sigma_{aa}\equiv\Sigma_{m,n}\sigma_{aa}^{m,n}$ and $\sigma_{bb}\equiv\Sigma_{m,n}\sigma_{bb}^{m,n}$ are the total probabilities of finding the dot empty or occupied. Obviously $\sigma_{aa}(t)=1-\overline{\sigma(t)}$, where $\overline{\sigma(t)}\equiv\sigma_{bb}(t)$. Performing the summation over m,n in Eqs. (2.1), we obtain the following rate equation for the quantum-dot occupation probability $\overline{\sigma}$:

$$\dot{\overline{\sigma(t)}}=\Gamma_L-(\Gamma_L+\Gamma_R)\overline{\sigma(t)}. \tag{2.4}$$

If the point contact and the quantum dot are decoupled, the detector current is $I_d^{(0)}=D$. Hence the occupation of the quantum dot can be measured through a variation of the detector current $\Delta I_d=I_d^{(0)}-I_d$. One readily obtains from Eq. (2.3a) that

$$\Delta I_d(t)=\frac{\Delta T V_d}{2\pi}\overline{\sigma(t)}, \tag{2.5}$$

where $V_d=\mu_L-\mu_R$ is the voltage bias, and $\Delta T=T-T'$. Thus the point contact is indeed the measurement device. In fact, Eq. (2.5) is a self-evident one. Indeed, the variation of the point-contact current is $\Delta T V_d/2\pi$, and $\overline{\sigma}$ is the probability for such a variation.

It follows from Eqs. (2.1) and (2.3) that the same current $I_s(t)=\Gamma_R\overline{\sigma(t)}$ would flow through the quantum dot in the absence of the detector $(D=D'=0)$. This means that the point-contact detector is a noninvasive detector. This is not surprising, since only an electron inside the point contact (under the barrier) can affect an electron in the quantum dot. The relevant (tunneling) time is very short. Actually, it is zero in the tunneling Hamiltonian approximation, Eqs. (A1) and (B1), used for the derivation of the rate equations.

III. DETECTION OF ELECTRON OSCILLATIONS IN COUPLED DOTS

The well-known manifestation of quantum coherence is the oscillations of a particle in a double-well (double-dot) potential. The origin of these oscillations is the interference between the probability amplitudes of finding a particle in different wells. Hence one can expect that the disclosure of a particle (electron) in one of wells would generate the ''dephasing'' that eventually destroys these oscillations, even without distorting the energy levels of the system.

Let us investigate the mechanism of this process by taking for detector a noninvasive point contact. A possible set up is shown in Fig. 2. We assume that the transmission probability of the point contact is T when an electron occupies the right well, and it is T' when an electron occupies the left well. Here $T'<T$, since the right well is away from the point contact.

Now we apply the quantum-rate equations[8,9] to the whole system. However, in the distinction with the previous case, the electron transitions in the measured system take place

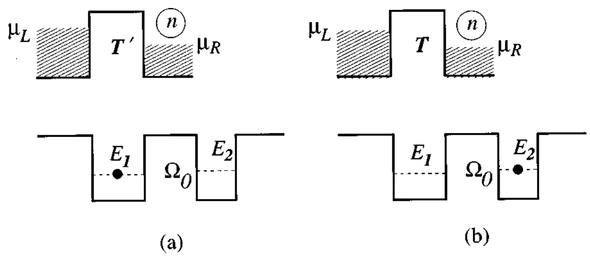

FIG. 2. Electron oscillations in the double well. The penetration coefficient of the point contact varies from T' to T when an electron occupies the left well (a) or right well (b), respectively. The index n denotes the number of electron accumulated in the collector at time t.

between the *isolated* states inside the dots. As a result, the diagonal density-matrix elements are coupled with the off-diagonal elements, so that the corresponding rate equations are the Bloch-type equations.[8,9,13,14]

We first start with the case of the double-well detached from the point-contact detector. The Bloch equations describing the time evolution of the electron-density matrix σ_{ij} have the following forms:

$$\dot{\sigma}_{aa}=i\Omega_0(\sigma_{ab}-\sigma_{ba}), \tag{3.1a}$$

$$\dot{\sigma}_{bb}=i\Omega_0(\sigma_{ba}-\sigma_{ab}), \tag{3.1b}$$

$$\dot{\sigma}_{ab}=i\epsilon\sigma_{ab}+i\Omega_0(\sigma_{aa}-\sigma_{bb}), \tag{3.1c}$$

where $\epsilon=E_2-E_1$, and Ω_0 is the coupling between the left and right wells. Here $\sigma_{aa}(t)$ and $\sigma_{bb}(t)$ are the probabilities of finding the electron in the left and right wells, respectively, and $\sigma_{ab}(t)=\sigma^*_{ba}(t)$ are the off-diagonal density-matrix elements (''coherences'').[14]

Solving these equations for the initial conditions and $\sigma_{aa}(0)=1$ and $\sigma_{bb}(0)=\sigma_{ab}(0)=0$, we obtain

$$\sigma_{aa}(t)=\frac{\Omega_0^2\cos^2(\omega t)+\epsilon^2/4}{\Omega_0^2+\epsilon^2/4}, \tag{3.2}$$

where $\omega=(\Omega_0^2+\epsilon^2/4)^{1/2}$. As expected, the electron initially localized in the first well oscillates between the wells with the frequency ω. Note that the amplitude of these oscillations is $\Omega_0^2/(\Omega_0^2+\epsilon^2/4)$. Thus the electron remains localized in the first well if the level displacement is large, $\epsilon\gg\Omega_0$.

Now we consider the electron oscillations in the presence of the point contact detector (Fig. 2). The corresponding Bloch equations for the entire system have the following forms (Appendix B):

$$\dot{\sigma}^n_{aa}=-D'\sigma^n_{aa}+D'\sigma^{n-1}_{aa}+i\Omega_0(\sigma^n_{ab}-\sigma^n_{ba}), \tag{3.3a}$$

$$\dot{\sigma}^n_{bb}=-D\sigma^n_{bb}+D\sigma^{n-1}_{bb}-i\Omega_0(\sigma^n_{ab}-\sigma^n_{ba}), \tag{3.3b}$$

$$\dot{\sigma}^n_{ab}=i\epsilon\sigma^n_{ab}+i\Omega_0(\sigma^n_{aa}-\sigma^n_{bb})-\frac{1}{2}(D'+D)\sigma^n_{ab} + (DD')^{1/2}\sigma^{n-1}_{ab}. \tag{3.3c}$$

Here the index n denotes the number of electrons arriving to the collector at time t, and $D(D')$ is the transition rate of an electron hopping from the left to the right detector reservoirs, $D=T(\mu_L-\mu_R)/2\pi$ [Eqs. (2.1)]. Notice that the presence of the detector results in additional terms in the rate equations in comparison with Eqs. (3.1). These terms are generated by transitions of an electron from the left to right detector reservoirs with the rates D and D', respectively. The equation for the nondiagonal density-matrix elements σ^n_{ab} is slightly different from the standard Bloch equations due to the last term, which describes the transition between different coherences, σ^{n-1}_{ab} and σ^n_{ab}. This term appears in the Bloch equations for coherences whenever the same hopping $(n-1\to n)$ takes place in *both* states of the off-diagonal density-matrix element (a and b) (see Refs. 8 and 9 and Appendix B). The rate of such transitions is determined by a product of the corresponding *amplitudes* ($T^{1/2}$ and $T'^{1/2}$).

It follows from Eqs. (2.3a) and (3.3) that the variation of the point-contact current $\Delta I_d(t)=I^{(0)}-I_d(t)$ measures directly the charge in the first dot. Indeed, for the detector current one obtains

$$I_d(t)=\sum_n n[\sigma^n_{aa}(t)+\sigma^n_{bb}(t)]=D'\sigma_{aa}(t)+D\sigma_{bb}(t), \tag{3.4}$$

where $\sigma_{ij}=\Sigma_n\sigma^n_{ij}$. Therefore $\Delta I_d(t)$ is given by Eq. (2.5), where $\sigma(t)\equiv\sigma_{aa}(t)$.

In order to determine the influence of the detector on the double-well system, we trace out the detector states in Eqs. (3.3), thus obtaining

$$\dot{\sigma}_{aa}=i\Omega_0(\sigma_{ab}-\sigma_{ba}), \tag{3.5a}$$

$$\dot{\sigma}_{bb}=i\Omega_0(\sigma_{ba}-\sigma_{ab}), \tag{3.5b}$$

$$\dot{\sigma}_{ab}=i\epsilon\sigma_{ab}+i\Omega_0(\sigma_{aa}-\sigma_{bb})-\tfrac{1}{2}(\sqrt{D}-\sqrt{D'})^2\sigma_{ab}, \tag{3.5c}$$

where $\sigma_{ij}=\Sigma_n\sigma^n_{ij}(t)$.

These equations coincide with Eqs. (3.1), describing the electron oscillations without a detector, except for the last term in Eq. (3.5c). The latter generates the exponential damping of the nondiagonal density-matrix element with the ''dephasing'' rate[15-17]

$$\Gamma_d=(\sqrt{D}-\sqrt{D'})^2=(\sqrt{T}-\sqrt{T'})^2\frac{V_d}{2\pi}. \tag{3.6}$$

This implies that $\sigma_{ab}\to 0$ for $t\to\infty$. We can check this by looking for the stationary solutions of Eqs. (3.5) in the limit $t\to\infty$. In this case $\dot{\sigma}_{ij}(t\to\infty)=0$, and Eqs. (3.5) become linear algebraic equations, which can be easily solved. One finds that the electron-density matrix becomes the statistical mixture

$$\sigma(t)=\begin{pmatrix}\sigma_{aa}(t) & \sigma_{ab}(t)\\ \sigma_{ba}(t) & \sigma_{bb}(t)\end{pmatrix}\to\begin{pmatrix}\frac{1}{2} & 0\\ 0 & \frac{1}{2}\end{pmatrix} \quad \text{for} \quad t\to\infty. \tag{3.7}$$

Notice that the damping of the nondiagonal density-matrix elements comes entirely from the possibility of disclosing the electron in one of the wells. Indeed, if the detector does not distinguish which of the wells is occupied, i.e., $T=T'$, then $\Gamma_d=0$.

The Bloch equations (3.3) and (3.5) display explicitly the mechanism of the dephasing during a noninvasive measurement, i.e., that which does not distort the energy levels of the measured system.[18] The dephasing appears in the reduced density-matrix as the ''dissipative'' term in the nondiagonal density matrix elements only, as a result of tracing out the detector variables. All other terms related to the detector are canceled after tracing out the detector variables. It is important to note that such a dephasing term in Eq. (3.5c) generates the ''collapse'' of the electron-density matrix into the statistical mixture, Eq. (3.7), without explicit use of the measurement reduction postulate.[19] The collapse is fully described by the Bloch-type equations, derived from the Schrödinger equation (Appendix B).

In fact, the idea that the dissipative interaction of a measured system with a detector can be responsible for the density-matrix collapse is not new. It was discussed in many publications, as for instance in works of Zurek,[20] which stressed conceptual points, or in detailed studies of more specific examples of atomic transitions.[21] Yet the present study of mesoscopic systems elaborates additional aspects of the dephasing problem. These are the dephasing mechanism due to continuous observation with a noninvasive detector, and the striking difference between the continuous and rapidly repeated measurements. The latter is discussed below.

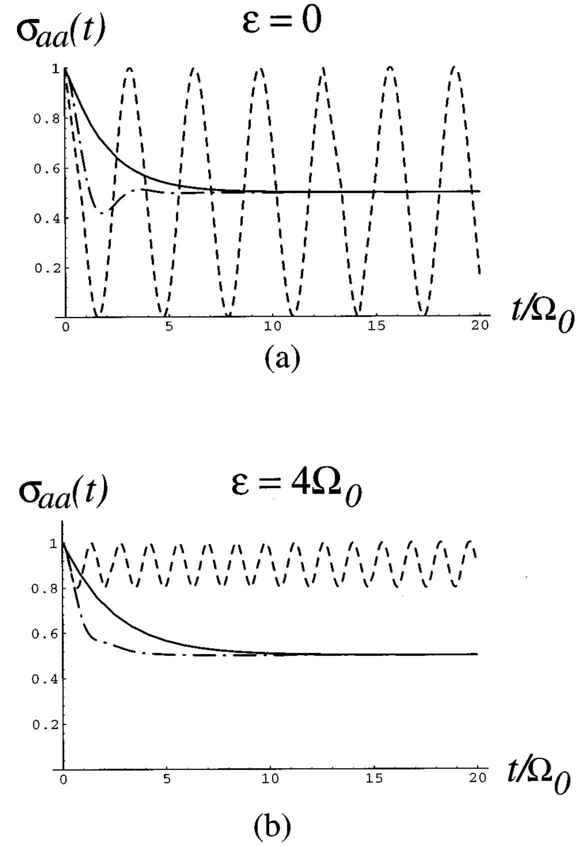

FIG. 3. The occupation of the first well as a function of time [Eqs. (3.5)]: (a) the levels are aligned ($\epsilon=0$); (b) the levels are displaced ($\epsilon=4\Omega_0$). The curves correspond to different values of the dephasing rate: $\Gamma_d=0$ (dashed), $\Gamma_d=4\Omega_0$ (dot-dashed), and $\Gamma_d=16\Omega_0$ (solid).

Continuous measurement and Zeno effect

The most surprising phenomenon which displays Eq. (3.7) is that the transition to the statistical mixture takes place even for a large displacement of the energy levels, $\epsilon \gg \Omega_0$, irrespectively of the initial conditions. This means that an electron initially localized in one of the wells would be always *delocalized* at $t\to\infty$. This would happen even if the electron was initially localized at the lower level. (Of course, this does not violate the energy conservation, since the double well is not isolated.) Such a behavior is not expected, because the amplitude of electron oscillations is very small for large level displacement [Eq. (3.2)]. Thus the electron should stay localized in one of the wells. One could expect that the continuous observation of this electron by a detector could only increase its localization. This can be inferred from the so-called Zeno effect.[22] The latter tells us that repeated observations of the system slow down transitions between quantum states due to the collapse of the wave function into the observed state. Since in our case the change of the detector current, $\Delta I_s(t)$ monitors $\overline{\sigma(t)}$ in the left well [Eqs. (2.5) and (3.4)], it represents the continuous measurement of the charge in this well. Nevertheless the effect is just the opposite—the continuous measurement delocalizes the system.[23]

In fact, our results for small t seem to be in an agreement with the Zeno effect; even so, we have not explicitly implied the projection postulate. For instance, Fig. 3(a) shows the time-dependence of the probability to find an electron in the left dot, as obtained from the solution of Eqs. (3.5) for the aligned levels ($\epsilon=0$), and $\Gamma_d=0$ (dashed curve), $\Gamma_d=4\Omega_0$ (dot-dashed curve), and $\Gamma_d=16\Omega_0$ (solid curve). One finds that for small t the rate of transition from the left to the right well slows down with the increase of Γ_d.

We find the same slowing down of the transition rate for small t for the disaligned levels ($\epsilon=4\Omega_0$) in Fig. 3(b). This implies that very frequent repeated measurements would indeed localize the system. In that sense the Bloch equations reproduce the Zeno effect without explicit use of the projection postulate. Actually, this result was found earlier by an analysis of atomic transitions by using the Bloch equation for a three-level system.[24,25] It was shown that the repeated measurement with short intervals $\Delta t=t/n$ localizes the system in the limit $n\to\infty$. Yet in our case the continuous measurement leads to an electron *delocalization*, whereas in the absence of detector an electron would stay localized in the left well [the dashed curve in Fig. 3(b)]. Thus the continuous and very frequent repeated measurements affect the system in opposite ways.

Our microscopic treatment allows us to determine the origin of the difference in both treatments. One easily finds that the derivation of the Bloch-type equations, describing the measured system [Eqs. (3.5)] implies the tracing of the detector variables [Eq. (1.1)]. Since this procedure is outside the Schrödinger equation, it could distort the time development of the system. In our case of continuous measurement the tracing is done at the time t, whereas the frequent repeated measurement with the intervals $\Delta t=t/n$ implies that the tracing of the detector variables takes part at the end of

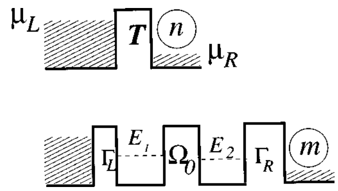

FIG. 4. Resonant tunneling through the double dot. $\Gamma_{L,R}$ denote the corresponding rate for the tunneling from (to) the left (right) reservoirs. The penetration coefficient of the point contact is T for the empty double-dot system or for the occupied second dot, and it is T' for the occupied first dot. The indices m and n denote the number of electrons penetrating to the right reservoirs at time t.

each interval Δt. As a result, the limit of $n\to\infty$, the measured system, stays localized.[24]

IV. MEASUREMENT OF RESONANT CURRENT IN COUPLED DOTS

In spite of great progress made in microfabrication techniques, the direct measurement of single-electron oscillations in a coupled-dot system is still a complicated problem. However, it is much easier to measure similar quantum coherence effects in electrical current flowing through coupled-dot systems. We therefore consider the same coupled dot of Sec. III, but now connected with two reservoirs (emitter and collector). As in the previous example the point-contact detector measures the occupation of the left dot (Fig. 4). For the sake of simplicity we assume strong inner and interdot Coulomb repulsion, so only one electron can occupy this system.[13] Then there are only three available states of the coupled-dot system: the dots are empty (a), the first dot is occupied (b), and the second dot is occupied (c). In analogy with Eqs. (2.1) and (3.3) we write the following Bloch equations for the density matrix $\sigma_{ij}^{m,n}(t)$ describing the entire system:[8,9]

$$\dot{\sigma}_{aa}^{m,n}=-(\Gamma_L+D)\sigma_{aa}^{m,n}+\Gamma_R\sigma_{cc}^{m-1,n}+D\sigma_{aa}^{m,n-1}\,, \tag{4.1a}$$

$$\dot{\sigma}_{bb}^{m,n}=-D'\sigma_{bb}^{m,n}+D'\sigma_{bb}^{m,n-1}+\Gamma_L\sigma_{aa}^{m,n}+i\Omega_0(\sigma_{bc}^{m,n}-\sigma_{cb}^{m,n}), \tag{4.1b}$$

$$\dot{\sigma}_{cc}^{m,n}=-(\Gamma_R+D)\sigma_{cc}^{m,n}+D\sigma_{cc}^{m,n-1}-i\Omega_0(\sigma_{bc}^{m,n}-\sigma_{cb}^{m,n}), \tag{4.1c}$$

$$\dot{\sigma}_{bc}^{m,n}=i\epsilon\sigma_{bc}^{m,n}+i\Omega_0(\sigma_{bb}^{m,n}-\sigma_{cc}^{m,n})-\frac{1}{2}(\Gamma_R+D'+D)\sigma_{bc}^{m,n}+(DD')^{1/2}\sigma_{bc}^{m,n-1}\,, \tag{4.1d}$$

where the indices n and m denote the number of electrons that arrive at time t to the upper and the lower collector reservoirs, respectively. Here Γ_L and Γ_R are the rates of electron transitions from the left reservoir to the first dot, and from the second dot to the right reservoir, and Ω_0 is the amplitude of hopping between two dots.

The currents in the double-dot system (I_s) and in the detector (I_d) are given by the following expressions [cf. Eqs. (2.2) and (2.3b)]:

$$I_s=\sum_{m,n} m(\dot{\sigma}_{aa}^{m,n}+\dot{\sigma}_{bb}^{m,n}+\dot{\sigma}_{cc}^{m,n})=\Gamma_R\sigma_{cc}\,, \tag{4.2a}$$

$$I_d=\sum_{m,n} n(\dot{\sigma}_{aa}^{m,n}+\dot{\sigma}_{bb}^{m,n}+\dot{\sigma}_{cc}^{m,n})=D-(D-D')\sigma_{bb}, \tag{4.2b}$$

where $\sigma_{ij}=\Sigma_{m,n}\sigma_{ij}^{m,n}$. It follows from Eq. (4.2b) that the variation of the detector current $\Delta I_d=I_d^{(0)}-I_d$ is given by Eq. (2.5), where $\overline{\sigma}=\sigma_{bb}$. Thus the point contact measures the occupation of the left dot directly.

Performing summation in Eqs. (4.1) over the number of electrons arriving at the collectors (m,n), we obtain the following Bloch equations for the reduced density matrix of the double-dot system:

$$\dot{\sigma}_{aa}=-\Gamma_L\sigma_{aa}+\Gamma_R\sigma_{cc}\,, \tag{4.3a}$$

$$\dot{\sigma}_{bb}=\Gamma_L\sigma_{aa}+i\Omega_0(\sigma_{bc}-\sigma_{cb}), \tag{4.3b}$$

$$\dot{\sigma}_{cc}=-\Gamma_R\sigma_{cc}-i\Omega_0(\sigma_{bc}-\sigma_{cb}), \tag{4.3c}$$

$$\dot{\sigma}_{bc}=i\epsilon\sigma_{bc}+i\Omega_0(\sigma_{bb}-\sigma_{cc})-\tfrac{1}{2}(\Gamma_R+\Gamma_d)\sigma_{bc}\,, \tag{4.3d}$$

where Γ_d is the dephasing rate generated by the detector [Eq. (3.6)]. These equations can be compared with those describing electron transport through the same system, but without a detector.[8,9,13] We find that the difference appears only in the nondiagonal density-matrix elements [Eq. (4.3d)]. The latter includes an additional dissipation rate Γ_d generated by the detector.

Solving Eqs. (4.3) in the limit $t\to\infty$, we find the following expression for the current I_s, Eq. (4.2b), flowing through the double-dot system:

$$I_s=\frac{(\Gamma_R+\Gamma_d)\Omega_0^2}{\epsilon^2+\dfrac{(\Gamma_R+\Gamma_d)^2}{4}+\Omega_0^2(\Gamma_R+\Gamma_d)\left(\dfrac{2}{\Gamma_R}+\dfrac{1}{\Gamma_L}\right)}. \tag{4.4}$$

By analyzing Eq. (4.4), one finds that all the measurement effects, discussed in Sec. III are reflected in the behavior of the resonant current I_s as a function of the level displacement ϵ and the dephasing rate Γ_d. As an example, in Fig. 5 we show the resonant current $I_s(\epsilon)$ for three values of the dephasing rate: $\Gamma_d=0$, $\Gamma_d=4\Omega_0$, and $\Gamma_d=16\Omega_0$. We find that for small ϵ the current decreases with Γ_d. However, for larger values of ϵ the current *increases* with Γ_d. This reflects an electron delocalization in a double-well system [Fig. 3(b)], due to continuous monitoring of the charge in the left dot. In contrast, rapidly repeated measurements[21,24] would always localize an electron and therefore diminish the current I_s.

V. SUMMARY

In this paper we studied the mechanism of decoherence generated by continuous observation of one of the states out of the coherent superposition in experiments with mesos-

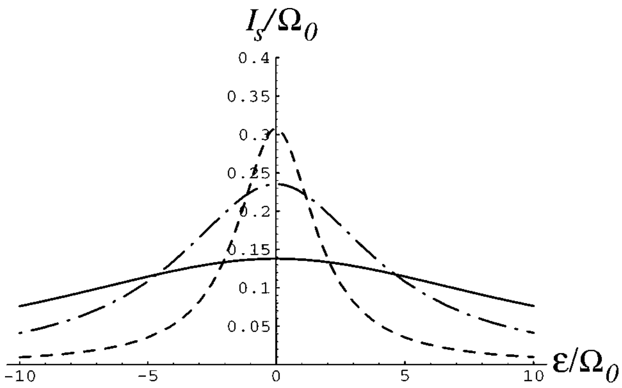

FIG. 5. Electron current through the double dot, Eq. (4.4), for $\Gamma_L = \Gamma_R = \Omega_0$ as a function of the level displacement $\epsilon = E_2 - E_1$. The curves correspond to different values of the dephasing rate: $\Gamma_d = 0$ (dashed), $\Gamma_d = 4\Omega_0$ (dot dashed), and and $\Gamma_d = 16\Omega_0$ (solid).

copic systems. As an example, we considered a coupled quantum-dot system, which is simple enough for a detailed theoretical treatment of the measured object and the detector together. On the other hand, it bears all the essential physics of the measurement process. For a description of the entire system, we applied Bloch-type equations, which are obtained from the many-body Schrödinger equation and provide the most simple and transparent treatment of quantum coherence effects.

As the detector, we used the point contact in close proximity to one of the dots. We demonstrated that the variation of the point-contact current due to electrostatic interaction with electrons in the dot measures directly the occupation of this dot.

We started with quantum oscillations of an electron in coupled quantum dots. It appears that the presence of the point-contact detector near one of the dots generates the dephasing rate in the Bloch equations for the off-diagonal density-matrix elements. We found that the dephasing rate is proportional to the variation of the point-contact transmission amplitude squared [Eq (3.6)]. The Bloch equations for the diagonal density-matrix elements are not affected by the detector, providing that it does not distort the energy levels of the double-dot system.

The appearance of the dephasing rate Γ_d in the Bloch equation leads to the collapse of the density matrix into the statistical mixture at $t\to\infty$ [Eq. (3.7)]. The collapse happens even for a large disalignment of the energy levels. In this case the measurement process results in an electron delocalization inside the double dot (after some critical time $t > t_0$), which otherwise would stay localized in one of the dots. This contradicts a common opinion that the continuous measurement always leads to a localization due to the wave-packet reduction (Zeno effect). In fact, the localization would take place if we considered the continuous measurement as rapidly repeated measurements with intervals $\Delta t = t/n$ for $n\to\infty$. The reason for such a different behavior of the measured system stems from the different procedure of tracing out of the detector variables from the total density matrix.

The same measurement effects appear in dc current flowing through coupled dots. We found that the dc current vanishes for $\Gamma_d\to\infty$, which can be interpreted in terms of an electron localization due to the Zeno effect. Nevertheless, for a finite Γ_d, and for disaligned energy levels ($E_1 \neq E_2$), the dc current *increases* with Γ_d. Here again, the situation is opposite to that of rapidly repeated measurement, where the current always *decreases* with Γ_d.

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APPENDIX A: RATE EQUATIONS FOR A POINT-CONTACT DETECTOR

Here we present the microscopic derivation of the rate equations describing electron transport in the point contact. The point contact is considered as a barrier, separating two reservoirs (the emitter and the collector) (Fig. 1). All the levels in the emitter and the collector are initially filled up to the Fermi energies μ_L and μ_R, respectively. We call it the ''vacuum'' state $|0\rangle$. The tunneling Hamiltonian $\mathcal{H}_{\rm PC}$ describing this system can be written as

$$\mathcal{H}_{\rm PC} = \sum_l E_l a_l^\dagger a_l + \sum_r E_r a_r^\dagger a_r + \sum_{l,r} \Omega_{lr}(a_l^\dagger a_r + \text{H.c.}), \tag{A1}$$

where $a_l^\dagger(a_l)$ and $a_r^\dagger(a_r)$ are the creation (annihilation) operators in the left and right reservoirs, respectively, and $\Omega_{\rm lr}$ is the hopping amplitude between the states E_l and E_r in the right and left reservoirs. (We choose the the gauge, where $\Omega_{\rm lr}$ is real.) The Hamiltonian equation (A1) requires the vacuum state $|0\rangle$ to decay exponentially to continuum states having the following forms: $a_r^\dagger a_l|0\rangle$, with an electron in the collector continuum and a hole in the emitter continuum; $a_r^\dagger a_{r'}^\dagger a_l^\dagger a_{l'}|0\rangle$ with two electrons in the collector continuum and two holes in the emitter continuum, and so on. The many-body wave function describing this system can be written in the occupation number representation as

$$|\Psi(t)\rangle = \Big[b_0(t) + \sum_{l,r} b_{lr}(t) a_r^\dagger a_l + \sum_{l<l',r<r'} b_{ll'rr'}(t) a_r^\dagger a_{r'}^\dagger a_l a_{l'} + \cdots \Big] |0\rangle, \tag{A2}$$

where $b(t)$ are the time-dependent probability amplitudes to find the system in the corresponding states with the initial condition $b_0(0)=1$, and all the other $b(0)$'s being zeros. Substituting Eq. (A2) into the Schrödinger equation $i|\dot{\Psi}(t)\rangle = \mathcal{H}_{\rm PC}|\Psi(t)\rangle$ and performing the Laplace transform,

$$\tilde{b}(E) = \int_0^\infty e^{iEt} b(t) dt, \tag{A3}$$

we obtain an infinite set of the coupled equations for the amplitudes $\tilde{b}(E)$:

$$E\tilde{b}_0(E) - \sum_{l,r} \Omega_{lr}\tilde{b}_{lr}(E) = i, \tag{A4a}$$

$$(E+E_l-E_r)\tilde{b}_{lr}(E)-\Omega_{lr}\tilde{b}_0(E)-\sum_{l',r'}\Omega_{l'r'}\tilde{b}_{ll'rr'}(E)=0, \tag{A4b}$$

$$(E+E_l+E_{l'}-E_r-E_{r'})\tilde{b}_{ll'rr'}(E)-\Omega_{l'r'}\tilde{b}_{lr}(E)+\Omega_{lr}\tilde{b}_{l'r'}(E)-\sum_{l'',r''}\Omega_{l''r''}\tilde{b}_{ll'l''rr'r''}(E)=0,$$

$$\cdots. \tag{A4c}$$

Equations (A4) can be substantially simplified by replacing the amplitude $\tilde{b}$ in the term $\Sigma\Omega\tilde{b}$ of each of the equations by its expression obtained from the subsequent equation.[8,9] For example, substituting $\tilde{b}_{lr}(E)$ from Eq. (A4b) into Eq. (A4a), one obtains

$$\left[E-\sum_{l,r}\frac{\Omega^2}{E+E_l-E_r}\right]\tilde{b}_0(E)-\sum_{ll',rr'}\frac{\Omega^2}{E+E_l-E_r}\tilde{b}_{ll'rr'}(E)=i, \tag{A5}$$

where we assumed that the hopping amplitudes are functions weakly dependent on the energies $\Omega_{lr}\equiv\Omega(E_l,E_r)=\Omega$. Since the states in the reservoirs are very dense (continuum), one can replace the sums over l and r by integrals, for instance $\Sigma_{l,r}\rightarrow\int\rho_L(E_l)\rho_R(E_r)dE_ldE_r$, where $\rho_{L,R}$ are the density of states in the emitter and collector. Then the first sum in Eq. (A5) becomes an integral which can be split into a sum of the singular and principal value parts. The singular part yields $i\pi\Omega^2\rho_L\rho_RV_d$, and the principal part is merely included in a redefinition of the energy levels. The second sum in Eq. (A5) can be neglected. Indeed, by replacing $\tilde{b}_{ll'rr'}(E)\equiv\tilde{b}(E,E_l,E_{l'},E_r,E_{r'})$ and the sums by the integrals, we find that the integrand has poles on the same sides of the integration contours. This means that the corresponding integral vanishes, providing $V_d\gg\Omega^2\rho$.

Applying analogous considerations to the other equations of system (A4), we finally arrive at the following set of equations:

$$(E+iD/2)\tilde{b}_0=i, \tag{A6a}$$

$$(E+E_l-E_r+iD/2)\tilde{b}_{lr}-\Omega\tilde{b}_0=0, \tag{A6b}$$

$$(E+E_l+E_{l'}-E_r-E_{r'}+iD/2)\tilde{b}_{ll'rr'}-\Omega\tilde{b}_{lr}+\Omega\tilde{b}_{l'r'}=0,$$

$$\cdots, \tag{A6c}$$

where $D=2\pi\Omega^2\rho_L\rho_RV_d$.

The charge accumulated in the collector at time t is

$$N_R(t)=\langle\Psi(t)|\sum_r a_r^\dagger a_r|\Psi(t)\rangle=\sum_n n\sigma^{(n)}(t), \tag{A7}$$

where

$$\sigma^{(0)}(t)=|b_0(t)|^2,\quad \sigma^{(1)}(t)=\sum_{l,r}|b_{lr}(t)|^2,$$

$$\sigma^{(2)}(t)=\sum_{ll',rr'}|b_{ll'rr'}(t)|^2,\cdots \tag{A8}$$

are the probabilities to find n electrons in the collector. These probabilities are directly related to the amplitudes $\tilde{b}(E)$ through the inverse Laplace transform

$$\sigma^{(n)}(t)=\sum_{l\ldots,r\ldots}\int\frac{dEdE'}{4\pi^2}\tilde{b}_{l\ldots r\ldots}(E)\times\tilde{b}^*_{l\ldots r\ldots}(E')e^{i(E'-E)t}. \tag{A9}$$

Using Eq. (A9), one can transform Eqs. (A6) into the rate equations for $\sigma^{(n)}(t)$ (cf. Refs. 8 and 9). We find

$$\dot{\sigma}^{(0)}(t)=-D\sigma^{(0)}(t), \tag{A10a}$$

$$\dot{\sigma}^{(1)}(t)=D\sigma^{(0)}(t)-D\sigma^{(1)}(t), \tag{A10b}$$

$$\dot{\sigma}^{(2)}(t)=D\sigma^{(1)}(t)-D\sigma^{(2)}(t), \tag{A10c}$$

$$\cdots.$$

The operator, which defines the current flowing in this system, is

$$\hat{I}=i\left[\mathcal{H}_{\rm PC},\sum_r a_r^\dagger a_r\right]=i\sum_{l,r}\Omega_{lr}(a_l^\dagger a_r-a_r^\dagger a_l). \tag{A11}$$

Using Eqs. (A2), (A10), and (A11), we find, for the current,

$$I=\langle\Psi(t)|\hat{I}|\Psi(t)\rangle=D\sum_n\sigma^{(n)}(t)=D. \tag{A12}$$

Since $D=(2\pi)^2\Omega^2\rho_L\rho_R=T$,[26] where T is the transmission probability, the current can be rewritten as $I=TV_d/(2\pi)$, which is the well-known Landauer formula.

APPENDIX B: POINT-CONTACT DETECTOR NEAR DOUBLE WELL

Now we present the microscopic derivation of the Bloch equations (3.3) describing electron oscillations in a double well with a point contact in close proximity to one of the wells (Fig. 2). We start with the many-body Schrödinger equation $i|\dot{\Psi}(t)\rangle=\mathcal{H}|\Psi(t)\rangle$ for the entire system. Here $\mathcal{H}$ is the tunneling Hamiltonian, which can be written as $\mathcal{H}=\mathcal{H}_{\rm PC}+\mathcal{H}_{\rm DD}+\mathcal{H}_{\rm int}$. Here $\mathcal{H}_{\rm PC}$ is the tunneling Hamiltonian for the point-contact detector [Eq. (A1)]; $\mathcal{H}_{\rm DD}$ is the tunneling Hamiltonian for the measured double-dot system,

$$\mathcal{H}_{\rm DD}=E_1c_1^\dagger c_1+E_2c_2^\dagger c_2+\Omega_0(c_2^\dagger c_1+c_1^\dagger c_2); \tag{B1}$$

and $\mathcal{H}_{\rm int}$ describes the interaction between the detector and the measured system. Since the presence of an electron in the left well results in an effective increase of the point-contact barrier ($\Omega_{lr}\rightarrow\Omega_{lr}+\delta\Omega_{lr}$), we can represent the interaction term as

$$\mathcal{H}_{\rm int}=\sum_{l,r}\delta\Omega_{lr}c_1^\dagger c_1(a_l^\dagger a_r+{\rm H.c.}). \tag{B2}$$

The many-body wave function for the entire system can be written as

$$\begin{aligned}|\Psi(t)\rangle=\bigg[&b_1(t)c_1^\dagger+\sum_{l,r}b_{1lr}(t)c_1^\dagger a_r^\dagger a_l\\&+\sum_{l<l',r<r'}b_{1ll'rr'}(t)c_1^\dagger a_r^\dagger a_{r'}^\dagger a_l a_{l'}\\&+b_2(t)c_2^\dagger+\sum_{l,r}b_{2lr}(t)c_2^\dagger a_r^\dagger a_l\\&+\sum_{l<l',r<r'}b_{2ll'rr'}(t)c_2^\dagger a_r^\dagger a_{r'}^\dagger a_l a_{l'}+\cdots\bigg]|0\rangle,\end{aligned} \tag{B3}$$

where $b(t)$ are the probability amplitudes to find the entire system in the states defined by the corresponding creation and annihilation operators. Notice that Eq. (B3) has the same form as Eq. (A2), where only the probability amplitudes $b(t)$ acquire an additional index (1 or 2) that denotes the well, occupied by an electron. Proceeding in the same way as in Appendix A, we arrive at an infinite set of the coupled equations for the amplitudes $\tilde{b}(E)$, which are the Laplace transform of the amplitudes $b(t)$ [Eq. (A3)]:

$$(E-E_1)\tilde{b}_1(E)-\Omega_0\tilde{b}_2(E)-\sum_{l,r}\Omega'_{lr}\tilde{b}_{1lr}(E)=i, \tag{B4a}$$

$$(E-E_2)\tilde{b}_2(E)-\Omega_0\tilde{b}_1(E)-\sum_{l,r}\Omega_{lr}\tilde{b}_{2lr}(E)=0, \tag{B4b}$$

$$\begin{aligned}(E+E_l-E_1-E_r)\tilde{b}_{1lr}(E)-\Omega'_{lr}\tilde{b}_1(E)-\Omega_0\tilde{b}_{2lr}(E)\\-\sum_{l',r'}\Omega_{l'r'}\tilde{b}_{1ll'rr'}(E)=0,\end{aligned} \tag{B4c}$$

$$\begin{aligned}(E+E_l-E_2-E_r)\tilde{b}_{2lr}(E)-\Omega_{lr}\tilde{b}_2(E)-\Omega_0\tilde{b}_{1lr}(E)\\-\sum_{l',r'}\Omega_{l'r'}\tilde{b}_{2ll'rr'}(E)=0,\end{aligned} \tag{B4d}$$

$$\cdots.$$

The same algebra as that used in Appendix A and in Refs. 8 and 9 allows us to simplify these equations, which then become

$$(E-E_1+iD'/2)\tilde{b}_1-\Omega_0\tilde{b}_2=i, \tag{B5a}$$

$$(E-E_2+iD/2)\tilde{b}_2-\Omega_0\tilde{b}_1=0, \tag{B5b}$$

$$(E+E_l-E_1-E_r+iD'/2)\tilde{b}_{1lr}-\Omega'\tilde{b}_1-\Omega_0\tilde{b}_{2lr}=0, \tag{B5c}$$

$$(E+E_l-E_2-E_r+iD/2)\tilde{b}_{2lr}-\Omega\tilde{b}_2-\Omega_0\tilde{b}_{1lr}=0, \tag{B5d}$$

$$\cdots,$$

where $D=TV_d/2\pi$. [We assumed for simplicity that the hopping amplitude of the point contact is weakly dependent on the energies, so that $\Omega_{lr}\equiv\Omega(E_l,E_r)=\Omega$.]

Using the inverse Laplace transform (A9) we can transform Eqs. (B5) into differential equations for the density-matrix elements $\sigma_{ij}^{(n)}(t)$ $(i,j=1,2)$,

$$\begin{aligned}&\sigma_{ij}^{(0)}(t)=b_i(t)b_j^*(t),\quad \sigma_{ij}^{(1)}(t)=\sum_{l,r}b_{ilr}(t)b_{jlr}^*(t),\\&\sigma_{ij}^{(2)}(t)=\sum_{ll',rr'}b_{ill'rr'}(t)b_{jll'rr'}^*(t),\\&\cdots,\end{aligned} \tag{B6}$$

where n denotes the number of electrons accumulated in the collector. Consider, for instance the off-diagonal density-matrix element $\sigma_{12}^{(1)}(t)$. The corresponding differential equation for this term can by obtained by multiplying Eq. (B5c) by $\tilde{b}_{2lr}^*(E')$ and subtracting the complex conjugated Eq. (B5d) multiplied by $\tilde{b}_{1lr}(E)$. We then obtain

$$\begin{aligned}\int\frac{dEdE'}{4\pi^2}\sum_{l,r}\bigg\{&\left(E'-E-\epsilon-i\frac{D+D'}{2}\right)\tilde{b}_{1lr}(E)\tilde{b}_{2lr}^*(E')\\&-[\Omega\tilde{b}_{1lr}(E)\tilde{b}_2^*(E')-\Omega'\tilde{b}_{2lr}^*(E')\tilde{b}_1(E)]\\&-\Omega_0[\tilde{b}_{1lr}(E)\tilde{b}_{1lr}^*(E')-\tilde{b}_{2lr}^*(E')\tilde{b}_{2lr}(E)]\bigg\}e^{i(E'-E)t}\\=0.&\end{aligned} \tag{B7}$$

One easily finds that the first term in this equation equals $-i\dot{\sigma}_{12}^{(1)}-[\epsilon+i(D+D')/2]\sigma_{12}^{(1)}$ and the third term equals $-\Omega_0(\sigma_{11}^{(1)}-\sigma_{22}^{(1)})$. In order to evaluate the second term in Eq. (B7), we replace $\Sigma_{l,r}$ by the integrals, and substitute

$$\begin{aligned}\tilde{b}_{1lr}(E)&=\frac{\Omega'\tilde{b}_1(E)+\Omega_0\tilde{b}_{2lr}(E)}{E+E_l-E_1-E_r+iD'/2},\\\tilde{b}_{2lr}^*(E')&=\frac{\Omega\tilde{b}_2^*(E')+\Omega_0\tilde{b}_{1lr}^*(E')}{E'+E_l-E_2-E_r-iD/2},\end{aligned} \tag{B8}$$

obtained from Eqs. (B5c) and (B5d), into Eq. (B7). Then integrating over E_l and E_r, we find that the second term in Eq. (B7) becomes $2i\pi\Omega\Omega'\rho_L\rho_RV_d\sigma_{12}^{(0)}$. Thus Eq. (B7) can be rewritten as

$$\begin{aligned}\dot{\sigma}_{12}^{(1)}=&\,i\epsilon\sigma_{12}^{(1)}+i\Omega_0(\sigma_{11}^{(1)}-\sigma_{22}^{(1)})-\tfrac{1}{2}(D'+D)\sigma_{12}^{(1)}\\&+(DD')^{1/2}\sigma_{12}^{(0)}\end{aligned} \tag{B9}$$

which coincides with the Bloch equation (3.3c) for $n=1$ and $\sigma_{aa}\equiv\sigma_{11}$, $\sigma_{bb}\equiv\sigma_{22}$, and $\sigma_{ab}\equiv\sigma_{12}$. Applying the same procedure to each of the Eqs. (B5), we arrive at the Bloch equations (3.3) for density-matrix elements $\sigma_{ij}^{(n)}$.

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