Rabi oscillations in the four-level double-dot structure

under the influence of the resonant pulse

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Abstract

We study theoretically the quantum dynamics of an electron in the symmetric four-level double-dot structure under the influence of the monochromatic resonant pulse. The probability amplitudes of the eigenstates relevant for the quantum dynamics are found from the solution of the non-stationary Schrödinger equation. The first-order correction term to the solution obtained through the rotating wave approximation is calculated. The three-level double-dot dynamics and the two-level single-dot dynamics, as well as the off-resonant excitation process, are derived from the general formulae for corresponding choices of the pulse and structure parameters. The results obtained may be applied to the solid-state qubit design.

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I. INTRODUCTION

Recently the low-dimensional semiconductor structures containing a small number of electrons in the size-quantized conduction band have attracted much attention. One of the reasons for that interest consists in their potential applicability to the quantum information processing [1]. It is commonly believed that those structures may be scaled up to form the quantum register with appropriate number of qubits. During the past decade many proposals for the semiconductor qubit realization were made. Here we mention the phosphorous donors embedded in a silicon host [2, 3] and a wide class of the systems based on the quantum dots (QDs) (see, e.g., [4, 5, 6, 7, 8, 9, 10, 11, 12, 13, 14]).

One of the features characterizing the coherent evolution of a quantum system is the Rabi oscillations induced by an external field. If the system was initially in one of the eigenstates of unperturbed Hamiltonian, it starts to oscillate under perturbation between its eigenstates giving rise to a broad class of different phenomena observed in the experiments. In view of the quantum algorithm realization, the Rabi oscillations are often considered as the process generating the desired qubit-state evolution. The most popular setup to study the Rabi oscillations of the electron population in the semiconductor nanostructures is based on the double-dot structure containing a single electron in the conduction band [15, 16] or, alternatively, on the artificial $H_2^+$ molecule formed by the donor implantation techniques (the so-called ”charge qubit” [17]). In what follows we shall consider the QDs keeping in mind the general
character of the conclusions made below. The standard method usually exploited for an electron charge manipulations in such structures is based on the electrostatic field control. By applying an adiabatically switched voltage one can lower or raise the potential barrier that separates the QDs thus allowing an electron to tunnel between the QDs or localizing it in one of the QDs. The quantum evolution of an electron clearly demonstrates the Rabi-like behavior that originates from the coherent electron tunnelling between the QDs [15, 16].

There is another scheme for the manipulations with a single electron confined in the double-dot structure. It is based on the resonant interaction of an electron with the coherent electromagnetic pulse [18]. The laser field, instead of electrostatic one, is applied to the double-dot structure, and the quantum state engineering is realized via the optically-induced transitions between the size-quantized electron levels. As it was shown, the pulse parameters (the frequency, duration, and amplitude) can be chosen so as to drive the electron, localized initially in the ground state in one of the QDs, to the ground state of another QD via the state delocalized over QDs and used here as a “transport” state. If the states localized in different QDs are viewed as the Boolean states 0 and 1, then, e.g., the electron transfer between them may be considered as the unitary operation NOT. The idea was initially proposed by Openov [18] and then developed further in the works [19, 20, 21, 22, 23, 24]. The influence of strong electromagnetic fields on the tunnelling phenomena in several-level nanostructures was studied in Refs. [25, 26]. It was shown that a laser with
appropriate power and frequency can drive the electron between quantum wells in a finite-size quantum well nanostructure or localize it in one of them. In Ref. [18], the opposite effect - i.e., a weak-laser-induced electron transfer between two quantum dots, was considered. In that model of the one-electron quantum dynamics, the assumption of instantaneous spreading of an excited electron over the QD structure was made or, equivalently, the matrix element of optical dipole transition between the ground and the "transport" states was supposed to be much less than the matrix element for tunnelling of an excited electron between the QDs. This allowed to describe the dynamics within the framework known in the atomic optics as the three-level (or Λ) scheme. The probability amplitudes to find an electron in the states localized in each of the QDs and in the "transport" state delocalized over the QD structure, were found as functions of the time, the pulse parameters, and the structure parameters. Provided that the quantum evolution is coherent, this process describes the three-level Rabi oscillations of the electron population. However, in all of these studies, the rigorous quantitative analysis of the assumption concerning the instantaneous spreading of an excited electron over the double-dot structure has not been presented.

In this work we focus on the quantitative study of the coherent quantum dynamics of an electron in the symmetric four-level double-dot structure under the influence of the resonant laser pulse for arbitrary tunnelling rates between the excited states of the QDs. Here we give the detailed derivation of the analytical expressions for the
probability amplitudes of the electron eigenstates relevant for the quantum dynamics. The results will be presented in terms of the basis states of isolated QDs. In this picture, the probability amplitudes are the explicit functions of both the tunnelling matrix element and the matrix element of the electron optical dipole transition. We show that the character of the system evolution is determined by the ratio between these matrix elements. The three-level double-dot and the two-level single-dot dynamics are derived from the general formulae as the limiting cases. We examine also the off-resonant excitation scheme that is very promising for the experimental realization of the proposed method of electron state manipulation.

The paper is organized as follows. In Section II we present the description of the model and obtain the general solution for the electron dynamics in the four-level double-dot structure. The important particular cases of the three- and two-level dynamics as well as the off-resonant Raman-like excitation are considered in Section III. Section IV contains the results of numerical simulations. The conclusions are summarized in Section V.

II. MODEL AND GENERAL SOLUTION

Let us consider the double-dot structure (see Fig.1) containing an electron in the size-quantized conduction band. For the sake of simplicity, we suppose the dots \( A \) and \( B \) to be identical. The existence of at least two one-electron orbital states \( |A(B)0\rangle \) and \( |A(B)1\rangle \) (ground and excited) in each of the QDs is assumed, with the one-electron wave functions \( \varphi_{A(B)0}(\mathbf{r}) = \langle \mathbf{r} | A(B)0 \rangle \) and \( \varphi_{A(B)1}(\mathbf{r}) = \langle \mathbf{r} | A(B)1 \rangle \),
respectively. Provided that the distance between the QDs is sufficiently large, the wave functions of the QD ground states are localized in corresponding QDs, and their overlap can be neglected. The overlap between the ground state and the excited state belonging to different QDs will be neglected as well: \( \langle A(B)0 \mid B(A)1 \rangle \approx 0 \). The excited levels are chosen to be close to the edge of the potential barrier separating the QDs. They couple through the electron tunnelling [18].

The Hamiltonian of an electron confined in the symmetric four-level double-dot structure is

\[
H_0 = \varepsilon_0 (|A0\rangle \langle A0| + |B0\rangle \langle B0|) + \varepsilon_1 (|A1\rangle \langle A1| + |B1\rangle \langle B1|) + \\
+ [-V(t) |A1\rangle \langle B1| + H.c.] ,
\]

where \( \varepsilon_0 \) and \( \varepsilon_1 \) are the one-electron energies of the ground and excited states, respectively (the same for both QDs); \( V(t) > 0 \) is the matrix element for the electron tunnelling between the excited states of the QDs, that, in general, may be a time-dependent function.

We consider the quantum evolution of an electron under the influence of the electromagnetic field that induces the optical transitions between the ground and excited states in each of the QDs (\( |A0\rangle \leftrightarrow |A1\rangle \) and \( |B0\rangle \leftrightarrow |B1\rangle \)). It is convenient to examine that evolution as a complex process including both the optical excitation of an electron in one of the QDs and the tunnelling of the excited electron into the other QD. The model Hamiltonian has the form

\[
H = H_0 + [\lambda(t) (|A0\rangle \langle A1| + |B0\rangle \langle B1|) + H.c.] ,
\]
where $\lambda(t)$ is the matrix element of the electron-field interaction. In what follows we shall suppose the matrix elements $V$ and $\lambda$ to be real and not show explicitly that they are time-dependent. The criterion of applicability of the model Hamiltonian (2) is expressed by the inequalities $|\lambda| \ll \omega_{10}, \ V \ll \omega_{10},$ where $\omega_{10} = \varepsilon_1 - \varepsilon_0$. Besides, we shall suppose that there are no additional levels localized in the near neighborhood of the energy $\varepsilon_1$.

The state vector of the system may be represented in terms of the eigenstates of isolated QDs as

$$|\Psi(t)\rangle = \sum_{n=A_0,B_0,A_1,B_1} c_n(t) \ |n\rangle. \quad (3)$$

Let an electron be initially localized in the ground state of the QD $A$, i.e. $|\Psi(0)\rangle = |A_0\rangle$. The quantum evolution of the state vector is governed by the non-stationary Schrödinger equation

$$i \frac{\partial |\Psi(t)\rangle}{\partial t} = H \ |\Psi(t)\rangle, \quad (4)$$

or, in the matrix form,

$$i \frac{\partial}{\partial t} \begin{pmatrix} c_{A_0} \\ c_{B_0} \\ c_{A_1} \\ c_{B_1} \end{pmatrix} = \begin{pmatrix} \varepsilon_0 & 0 & \lambda & 0 \\ 0 & \varepsilon_0 & 0 & \lambda \\ \lambda & 0 & \varepsilon_1 & -V \\ 0 & \lambda & -V & \varepsilon_1 \end{pmatrix} \begin{pmatrix} c_{A_0} \\ c_{B_0} \\ c_{A_1} \\ c_{B_1} \end{pmatrix} \quad (5)$$

with the initial conditions $c_n(0) = \delta_{n,A_0}$ (hereafter $\hbar = 1$).

The straightforward diagonalization of the Hamiltonian matrix amounts to the
set of eigenvectors

\[ |1\rangle = u_- (|A0\rangle + |B0\rangle) / \sqrt{2} + v_- (|A1\rangle + |B1\rangle) / \sqrt{2}, \]
\[ |2\rangle = u_+ (|A0\rangle - |B0\rangle) / \sqrt{2} + v_+ (|A1\rangle - |B1\rangle) / \sqrt{2}, \]
\[ |3\rangle = v_- (|A0\rangle + |B0\rangle) / \sqrt{2} - u_- (|A1\rangle + |B1\rangle) / \sqrt{2}, \]
\[ |4\rangle = v_+ (|A0\rangle - |B0\rangle) / \sqrt{2} - u_+ (|A1\rangle - |B1\rangle) / \sqrt{2}, \] (6)

and to the corresponding eigenenergies

\[ E_1 = \varepsilon_0 + \frac{1}{2} \left( \omega_- + \sqrt{\omega_-^2 + 4\lambda^2} \right), \quad E_2 = \varepsilon_0 + \frac{1}{2} \left( \omega_+ + \sqrt{\omega_+^2 + 4\lambda^2} \right), \]
\[ E_3 = \varepsilon_0 + \frac{1}{2} \left( \omega_- - \sqrt{\omega_-^2 + 4\lambda^2} \right), \quad E_4 = \varepsilon_0 + \frac{1}{2} \left( \omega_+ - \sqrt{\omega_+^2 + 4\lambda^2} \right), \] (7)

where \( u_\pm \approx \frac{1}{\omega_\pm} \left( 1 - \frac{3}{2} \frac{\lambda^2}{\omega_\pm^2} \right), \quad v_\pm \approx 1 - \frac{1}{2} \frac{\lambda^2}{\omega_\pm^2} \) and \( \omega_\pm = \omega_{10} \pm V \). (Here we restrict ourselves by the third-order terms in the parameters \( \frac{|\lambda|}{\omega_\pm} \ll 1 \). Note that the expansion of (7) over the small parameters \( \frac{|\lambda|}{\omega_\pm} \ll 1 \) gives rise to the Bloch-Siegert term \( \frac{\lambda^2}{\omega_{10}} \) in the eigenenergies \( E_i, i = 1 - 4 \).) The matrix diagonalizing the right-hand side of Eq. (5) has the form

\[ D = \begin{pmatrix} U_1 & U_2 \\ U_2 & -U_1 \end{pmatrix}, \] (8)

where

\[ U_1 = \frac{1}{\sqrt{2}} \begin{pmatrix} u_- & u_+ \\ u_- & -u_+ \end{pmatrix}, \quad U_2 = \frac{1}{\sqrt{2}} \begin{pmatrix} v_- & v_+ \\ v_- & -v_+ \end{pmatrix} \] (9)

(The columns of \( D \) are the transposed eigenvectors \( |k\rangle \), \( k = 1 - 4 \)).

The electron dynamics is most easily revealed in the representation of the instantaneous eigenstates, Eqs. (6), of the system Hamiltonian. In this basis, the time-dependent state vector of the system takes the form

\[ |\Phi(t)\rangle = \sum_{k=1}^{4} a_k(t) |k\rangle, \] (10)
where the instantaneous probability amplitudes \( a_k(t), k = 1 - 4 \), are related to the probability amplitudes \( c_n(t) \) by the matrix \( D \):

\[
c(t) = Da(t),
\]

where \( c(t) = [c_{A0}(t), c_{B0}(t), c_{A1}(t), c_{B1}(t)]^T \) and \( a(t) = [a_1(t), a_2(t), a_3(t), a_4(t)]^T \).

The Schrödinger equation in the new basis reads

\[
i \frac{\partial \Phi(t)}{\partial t} = \left\{ D^\dagger HD - iD^\dagger \frac{\partial D}{\partial t} \right\} \Phi(t)
\]

or, in the matrix form,

\[
i \frac{\partial}{\partial t} \begin{pmatrix} a_1 \\ a_2 \\ a_3 \\ a_4 \end{pmatrix} = \begin{pmatrix} E_1 & 0 & \mu_- & 0 \\ 0 & E_2 & 0 & \mu_+ \\ \mu_-^* & 0 & E_3 & 0 \\ 0 & \mu_+^* & 0 & E_4 \end{pmatrix} \begin{pmatrix} a_1 \\ a_2 \\ a_3 \\ a_4 \end{pmatrix},
\]

with the initial conditions \( a_{1(2)}(0) = u_{-(+)}/\sqrt{2}, \ a_{3(4)}(0) = v_{-(+)}/\sqrt{2} \). Here \( E_{1(2)} \approx \varepsilon_0 + \omega_{-(+)} \left( 1 + \frac{\lambda^2}{\omega_{-(+)}^2} \right) \), \( E_{3(4)} \approx \varepsilon_0 - \frac{\lambda^2}{\omega_{-(+)}^2} \) and \( \mu_\pm \approx i \frac{\partial}{\partial t} \left[ \frac{\lambda}{\omega_\pm} \left( 1 - \frac{4\lambda^2}{\omega_\pm^2} \right) \right] \), in agreement with the approximations for \( u_\pm \) and \( v_\pm \). Besides, we have neglected the terms \( \sim \frac{\lambda^3}{\omega_\pm^3} \frac{\partial}{\partial t} \left( \frac{\lambda}{\omega_\pm} \right) \) in \( E_i \) assuming them much smaller than the Bloch-Siegert term. This approximation requires that \( \left| \frac{\lambda^4}{\omega_\pm^4} \frac{\partial}{\partial t} \left( \frac{\lambda}{\omega_\pm} \right) \right| \ll 1 \) and imposes the constrain on the speed of changes in both the tunnelling matrix element \( V \) and the matrix element of the electron-field interaction \( \lambda \). The terms \( \sim \frac{\lambda^4}{\omega_\pm^4} \frac{\partial}{\partial t} \left( \frac{\lambda}{\omega_\pm} \right) \) in \( \mu_\pm \) have been rejected as well. At this stage of consideration we see that the pairs of coefficients \( (a_1, a_3) \) and \( (a_2, a_4) \) evolve independently from each other. The four-level problem of Eq. (5) thus reduces to the two-level ones that can, in principle, be solved by one of the standard methods developed earlier [27].
To proceed further we have to concretize the electron-field interaction term $\lambda$ in Eq. (2). In what follows we shall consider the situation where the tunnelling matrix element $V$ is time-independent. It is worth to mention that for the adiabatically switched voltages the Hamiltonian matrix in Eq. (13) is diagonal and its solution is straightforward. As an example we give here the solution describing the electron behavior under the influence of the electrostatic field $E_0 \neq E_0(t)$ ($\mu_{\pm} = 0$). The expression for $c$ follows directly from Eqs. (11) and (13). The probability amplitudes for an electron to be in the ground states of the QDs have the form

$$c_{A0} \approx e^{-i\left(\varepsilon_0 - \frac{\lambda_0^2}{\omega_{10}}\right)t} \cos\left(\frac{\lambda_0^2 V}{\omega_{10}} t\right), \quad c_{B0} \approx i e^{-i\left(\varepsilon_0 - \frac{\lambda_0^2}{\omega_{10}}\right)t} \sin\left(\frac{\lambda_0^2 V}{\omega_{10}} t\right),$$  

where $\lambda_0 = eE_0 \int \varphi_{A(B)0}^{\ast}(r) \varphi_{A(B)1}(r) dr$ is the matrix element of optical dipole transition between the states $|A(B)0\rangle$ and $|A(B)1\rangle$. The population is therefore localized (up to the $\lambda_0^2/\omega_{10}^2 \ll 1$ terms) in the ground-state subspace and exhibits the two-level Rabi oscillations at the frequency $\Omega_0 = \lambda_0^2 V/\omega_{10}^2$. The probabilities to find an electron in the states $|A1\rangle$ and $|B1\rangle$ are of order of $\lambda_0^2/\omega_{10}^2 \ll 1$ and oscillate at the frequency $V$. One sees that even for the static fields, a substantial electron state evolution in the double-dot structure occurs for the characteristic time $T \sim 1/\Omega_0$. In principle, the electron oscillations may be utilized for the qubit-state engineering but this process seems to be too slow in comparison with the resonant optical driving and unviable in view of the decoherence. However, this effect should be taken into account if quantum operations are performed through the sequence of electrostatic voltages since in this case the electron transitions
between the localized and delocalized states induced by the static fields bring about an unwanted qubit dynamics, i.e., a computational error.

The central part of our investigation will be devoted to the interaction of an electron with the time-dependent resonant pulses. For the sake of simplicity we consider a monochromatic square pulse of the amplitude $E_0$, the duration $T$, and the frequency $\omega$:

$$E(t) = E_0 \cos(\omega t) \left[ \theta(t) - \theta(t - T) \right].$$  \hspace{1cm} (15)

In this case

$$\lambda = \lambda_0 \cos(\omega t) \left[ \theta(t) - \theta(t - T) \right],$$  \hspace{1cm} (16)

where $\theta(t)$ is the step function. The frequency $\omega$ of the laser pulse may be detuned from the resonant frequencies $\omega_{\pm}$ by the values $\delta_{\pm} = \omega - \omega_{\pm}$, where $|\delta_{\pm}| \ll \omega$.

As we see from Eq. (13), it is sufficient to analyze the dynamics of just one pair of the coefficients, e.g., $a_1$ and $a_3$ (the dynamics for $a_2$ and $a_4$ is then revealed by the substitution $\omega_- \rightarrow \omega_+$ in the results obtained for $a_1$ and $a_3$). Transforming the coefficients $a_1$ and $a_3$ according to the formulae $a_1 = \tilde{a}_1 e^{-i\varepsilon_0 t}$ and $a_3 = \tilde{a}_3 e^{-i\varepsilon_0 t}$ and inserting the expression (16) for $\lambda$ into Eq. (13), we arrive at the set of two coupled linear differential equations for the coefficients $\tilde{a}_1$ and $\tilde{a}_3$ (an analogous set is obtained for the coefficients $\tilde{a}_2$ and $\tilde{a}_4$):

$$\begin{cases}
    i\dot{\tilde{a}}_1 = \frac{\lambda_0^2}{\omega} \cos(\omega t) \tilde{a}_1 + \frac{\lambda_0}{2} \left( 1 - e^{i2\omega t} \right) e^{-i\delta_{\pm} t} \tilde{a}_3, \\
    i\dot{\tilde{a}}_3 = -\frac{\lambda_0^2}{\omega} \cos(\omega t) \tilde{a}_3 + \frac{\lambda_0}{2} \left( 1 - e^{-i2\omega t} \right) e^{i\delta_{\pm} t} \tilde{a}_1,
\end{cases} \quad \lambda_0' = \lambda_0 \frac{\omega}{\omega_-}. \hspace{1cm} (17)$$

Here we restrict ourselves to the terms linear on the small parameter $\lambda_0/\omega$ in $\mu_{\pm}$ and
retain the Bloch-Siegert term $\sim \lambda_0^2 / \omega$. As we shall see below, the account of this term is necessary for obtaining the correct result within the first-order approximation on the parameter $\lambda_0 / \omega \ll 1$.

Usually, at this point the rotating wave approximation (RWA) is made and the fast oscillating exponents $e^{\pm i2\omega t}$ in Eqs. (17) are omitted. The solution thus accounts only of the one-photon processes conserving the energy of the system. This approximation is valid if the frequency $\omega_{10} \approx \omega \approx \omega_{\pm}$ dominates the Rabi frequency that is of the order of $|\lambda_0|$. Since this requirement on the pulse parameters is inherent to many optics setups, the theoretical predictions based on the RWA are in excellent agreement with the experimental data. However, any possible extension of the RWA seems to be very instructive in view of the quantitative estimate of corrections to the results obtained by the RWA. In Ref. [28] the authors suggested a simple and clear way of how to calculate the first-order correction term to the RWA solution. Using the adiabatic elimination procedure for the virtual two-photon states, they were able to find the probability amplitudes for a two-level system beyond the RWA. Here we shall utilize their method to solve the set of Eqs. (17).

According to Ref. [28], the coefficients $\tilde{a}_1$ and $\tilde{a}_3$ can be sought in the form:

$$\tilde{a}_1 \approx a_1^{(0)} + a_1^{(-)} e^{-i2\omega t} + a_1^{(+)} e^{i2\omega t},$$

$$\tilde{a}_3 \approx a_3^{(0)} + a_3^{(-)} e^{-i2\omega t} + a_3^{(+)} e^{i2\omega t},$$

(18)

where the higher-order terms proportional to $e^{\pm i2m\omega t}$, $m > 1$, are dropped. From
the Eqs. (17) and (18) we obtain the set of six equations:

\[
\begin{align*}
    i\dot{a}_1^{(0)} &= \frac{\lambda_0^2}{\omega} a_1^{(0)} + \frac{\lambda_0^2}{2} e^{-i\delta - t} (a_3^{(0)} - a_3^{(-)}) \\
    i\dot{a}_3^{(0)} &= -\frac{\lambda_0^2}{\omega} a_3^{(0)} + \frac{\lambda_0^2}{2} e^{i\delta - t} (a_1^{(0)} - a_1^{(+)}) \\
    i\dot{a}_1^{(-)} &= -2\omega a_1^{(-)} + \frac{\lambda_0^2}{2} e^{-i\delta - t} a_3^{(-)} \\
    i\dot{a}_3^{(-)} &= -2\omega a_3^{(-)} + \frac{\lambda_0^2}{2} e^{i\delta - t} (a_1^{(-)} - a_1^{(0)}) \\
    i\dot{a}_1^{(+)} &= 2\omega a_2^{(+)} + \frac{\lambda_0^2}{2} e^{-i\delta - t} (a_3^{(+)} - a_3^{(0)}) \\
    i\dot{a}_3^{(+)} &= 2\omega a_3^{(+)} + \frac{\lambda_0^2}{2} e^{i\delta - t} a_1^{(0)}
\end{align*}
\]

(19)

Adiabatic elimination yields

\[
a_1^{(-)} \approx 0, \ a_3^{(+)} \approx 0
\]

(20)

and

\[
a_1^{(+)} \approx \frac{\lambda_0}{4\omega} a_3^{(0)} e^{-i\delta - t}, \ a_3^{(-)} \approx -\frac{\lambda_0}{4\omega} a_1^{(0)} e^{i\delta - t},
\]

(21)

which in turn amount to the set

\[
\begin{align*}
    i\dot{\tilde{a}}_1^{(0)} &= \frac{5\lambda_0^2}{8\omega} a_1^{(0)} + \frac{\lambda_0^2}{2} e^{-i\delta - t} a_3^{(0)} \\
    i\dot{\tilde{a}}_3^{(0)} &= -\frac{5\lambda_0^2}{8\omega} a_3^{(0)} + \frac{\lambda_0^2}{2} e^{i\delta - t} a_1^{(0)}
\end{align*}
\]

(22)

Note that this procedure enables us to calculate only the first-order correction term \(\sim \lambda_0^2/\omega\) to the RWA since account of the higher order terms makes the set (22) incompatible.

The substitution \(a_1^{(0)} = \tilde{a}_1^{(0)} e^{-\frac{5\lambda_0^2}{4\omega} t}, \ a_3^{(0)} = \tilde{a}_3^{(0)} e^{\frac{5\lambda_0^2}{4\omega} t}\) transforms the set of Eqs. (22) into

\[
\begin{align*}
    i\dot{\tilde{a}}_1^{(0)} &= \frac{\lambda_0^2}{2} e^{-i\left(\delta - \frac{5\lambda_0^2}{4\omega} t\right)} \tilde{a}_3^{(0)} \\
    i\dot{\tilde{a}}_3^{(0)} &= \frac{\lambda_0^2}{2} e^{i\left(\delta - \frac{5\lambda_0^2}{4\omega} t\right)} \tilde{a}_1^{(0)}
\end{align*}
\]

(23)

that is equivalent to the following second-order differential equation:

\[
\ddot{\tilde{a}}_1^{(0)} + i \left(\delta - \frac{5\lambda_0^2}{4\omega}\right) \dot{\tilde{a}}_1^{(0)} + \frac{\lambda_0^2}{4} \tilde{a}_1^{(0)} = 0
\]

(24)
with the initial conditions $\tilde{a}^{(0)}_1(0) = \frac{3\lambda_0}{4\sqrt{2}\omega}$, $\tilde{a}^{(0)}_3(0) = -i\lambda_0 \frac{\sqrt{2}}{2}$.

Of course, we have to justify the adiabatic approximation used in Eqs. (20) - (21) by imposing the requirement on the pulse switching time $\tau_0$:

$$\omega^{-1} \ll \tau_0 \ll T. \tag{25}$$

In what follows, however, we shall continue to handle the ramp pulses since the accurate calculation carried out for an adiabatically switched pulse brings about simple renormalization of the matrix element $\lambda_0$ conserving the total character of the ramp-pulsed dynamics [28]. Note that the first of inequalities (25) ensures the applicability of Eq. (13) at $t \leq \tau_0$. Besides, we assume that $\tau_0 \ll |\lambda_0|^{-1}$. This unnecessary but very useful condition minimizes the influence of the pulse shape on the Rabi oscillation pattern.

The solution of Eq. (24) is straightforward; transforming it back to the coefficients $a^{(0)}_1$ and $a^{(0)}_3$, one has

$$\begin{align*}
\tilde{a}^{(0)}_1 &= \frac{1}{\sqrt{2}} e^{-i\frac{\delta_+}{2}} \left[ \frac{3\lambda_0}{4\omega} \cos (2\Omega'_- t) - i \frac{\lambda_0}{4\Omega'_-} \sin (2\Omega'_- t) \right], \\
\tilde{a}^{(0)}_3 &= \frac{1}{\sqrt{2}} e^{i\frac{\delta_+}{2}} \left[ \cos (2\Omega'_- t) - i \frac{\lambda_0}{4\Omega'_-} \sin (2\Omega'_- t) \right],
\end{align*} \tag{26}$$

where $\Omega'_- = \sqrt{\lambda'_0^2 + \delta'_-^2} / 4$, $\delta'_- = \delta_- - 5\lambda'_0^2 / 4\omega$. It is easy to verify that these expressions satisfy the set of Eqs. (19).

As a result, for the coefficients in the laboratory frame we obtain

$$\begin{align*}
\begin{align*}
c_{A0(B0)} &= \frac{1}{2} e^{-i\epsilon_0 t} \left[ f_0^- (t) \pm f_0^+ (t) \right], \\
\tilde{f}_0^\pm (t) &= e^{i\frac{\delta_+}{2}} \left[ \cos (2\Omega_{\pm} t) - i \frac{\delta_+}{4\Omega_{\pm}} \sin (2\Omega_{\pm} t) - i \frac{\lambda_0}{8\omega \Omega_{\pm}} e^{-i2\omega t} \sin (2\Omega_{\pm} t) \right]; \\
\begin{align*}
c_{A1(B1)} &= \frac{1}{2} \left[ e^{-i\epsilon t} f_1^- (t) \pm e^{-i\epsilon t} f_1^+ (t) \right], \\
\tilde{f}_1^\pm (t) &= e^{-i\frac{\delta_+}{2}} \frac{\lambda_0}{2\sqrt{2}\Omega_{\pm}^2} \left[ -i \sin (2\Omega_{\pm} t) + \frac{\Omega_{\pm}}{\omega} (1 - e^{i2\omega t}) \cos (2\Omega_{\pm} t) \right].
\end{align*}
\end{align*} \tag{27, 28}
\end{align*}$$

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where
\[
\Omega_{\pm} = \frac{1}{4} \sqrt{2\lambda^2 + \delta_\pm^2}
\] (29)

are the Rabi frequencies and \( \tilde{\lambda} = \frac{\lambda_0}{\sqrt{2}}, \varepsilon_\pm = \varepsilon_0 + \omega_\pm. \)

The formulae (27) - (29) describe the general type of the coherent one-electron evolution in the symmetric four-level double-dot structure driven by the resonant monochromatic pulse. One can check that the normalization condition
\[
\sum_{n=A0,B0,A1,B1} |c_n(t)|^2 = 1 + O\left(\frac{\lambda_0^2}{\omega^2}, \frac{\delta_{\pm}}{\omega}\right)
\] (30)
is fulfilled with the accuracy adopted through the calculations.

We couldn’t reveal any noticeable effect on the dynamics that would be produced by the small parameters \( \delta_\pm/\omega. \) They appear in the expression for the Rabi frequencies \( \Omega_{\pm}', \) Eq. (26), and may compete with the terms \( \delta_\pm^2/\lambda_0^2 \) if \( |\lambda_0|/\omega \geq |\delta_\pm|/\lambda_0|. \) Since we are not interested in detailed consideration of the system dynamics in that range of parameters we have neglected the terms \( \sim \delta_\pm/\omega \) in Eqs. (27) - (29). Note that the terms \( \lambda_0\delta_\pm/\omega V, \lambda_0\delta_\pm/\omega^2 \) which also contain this small parameter may be comparable to the terms of the order of \( \lambda_0^2/\omega^2 \) that have been omitted in the solution, and hence must be omitted as well.

III. QUANTUM DYNAMICS IN THE STRONG AND WEAK TUNNELING REGIMES.

In this Section we analyze the results obtained above for various choices of the pulse and structure parameters. From the general formulae (27) - (29), we derive
the expressions for the probability amplitudes corresponding to the situations where the characteristic tunnelling energy $V$ is either much greater or much less than the matrix element of optical dipole transition $|\lambda_0|$. Besides, we investigate also the electron dynamics in the strongly detuned Raman-like regime.

A. The three-level quantum dynamics.

If an electron being excited in one of the QDs tunnels into the other QD during a time much shorter than the Rabi oscillation period, one can speak about the simultaneous electronic excitation in both of QDs. In other words, one of the hybridized states, i.e. $(|A1\rangle + |B1\rangle)/\sqrt{2}$ or $(|A1\rangle - |B1\rangle)/\sqrt{2}$ (equally-weighted in each of the QDs), is excited. Both of these states are the eigenstates of the stationary Hamiltonian $H_0$, Eq. (1). It seems then preferable to expand the state vector $|\Psi (t)\rangle$ over the eigenstates of $H_0$ instead than over the states of isolated QD basis. Doing so and using the resonant approximation we may consider only one of hybridized states that is formally equivalent to the setting $V \gg |\lambda_0|$. The quantum dynamics of our system thus coincides with that of the three-level nanostructure. This situation was studied in the works [18] - [24] without, however, paying enough attention to the mathematical proof of that proposal. In what follows we shall consider this case in detail and show to what extent the electron dynamics may correspond to the scheme just sketched.

The condition of the strong tunnel coupling between the excited orbital states of the QDs, as compared to the optical dipole coupling between the ground and excited
orbital states of the single QD, is

\[ 2V = |\delta_+ - \delta_-| \gg |\lambda_0|. \]  (31)

We are interested in the resonant electron-pulse interaction, when the pulse frequency \( \omega \) matches one of the resonant frequencies \( \omega_\pm \) and is strongly detuned from the other one. For definiteness, let the pulse frequency \( \omega \) to be close to the frequency \( \omega_- \) so that \( |\delta_-| \ll |\delta_+| \) and, as it follows from Eq. (31), \( |\delta_+| \gg |\lambda_0| \). This choice of the pulse frequency corresponds to electron transition from \( |A(B)0\rangle \) to \( (|A1\rangle + |B1\rangle)/\sqrt{2} \). Obviously, there are two different situations concerning the mutual relation between the pulse detuning \( \delta_- \) and the value of \( |\lambda_0| \), i.e. \( |\delta_-| \ll |\lambda_0| \) and \( |\delta_-| \gg |\lambda_0| \). The first inequality corresponds to the resonant electron-pulse interaction, whereas the second one describes the off-resonant Raman-like coupling (see Sec. III C).

Since here we consider the resonant case, all of the conditions imposed on the system parameters may be summarized in the following inequalities:

\[ |\delta_-| \ll |\lambda_0| \ll |\delta_+|, \]  (32)

where \( \delta_+ \approx -2V \) and \( \omega_- = \omega_+ - 2V \).

Making use of Eqs. (27) - (29) and taking into account that

\[ \cos (2\Omega \pm t) - i \frac{\delta_\pm}{4\Omega_\pm} \sin (2\Omega \pm t) \approx e^{-i\frac{\delta_\pm}{4\Omega_\pm} \left(1 + \frac{\lambda_\pm^2}{\delta_\pm^2}\right)t} + isgn (\delta_\pm) \frac{\lambda_\pm^2}{\delta_\pm^2} \sin \left(\frac{|\delta_\pm|}{2}t\right) \]  (33)
at $|\lambda_0| \ll |\delta_\pm|$, let us rewrite the Eq. (3) in the form

$$
\begin{align*}
|\Psi\rangle &= \frac{1}{2} e^{-i\lambda_0 t} \left[ 1 + f_0^- (t) - \frac{\lambda^2}{4V^2} e^{-iVt} \sin(Vt) \right] |A0\rangle \\
&+ \frac{1}{\sqrt{2}} e^{-i\epsilon_0 t} f_1^- (t) \frac{|\{A1\} + |B1\rangle}{\sqrt{2}} - \frac{\lambda}{2V} e^{-i(\epsilon + V)t} \sin(Vt) \frac{|\{A1\} - |B1\rangle}{\sqrt{2}}.
\end{align*}
$$

(34)

It is easy to calculate from Eq. (34) the probability of the state inversion $p_{B0} (t) = |\langle B0 | \Psi(t) \rangle|^2$ after the pulse of the duration $T_\pi = \pi/\lambda_0$ (the so-called $\pi$-pulse) is off:

$$
p_{B0} (T_\pi) = 1 - \frac{\delta^2 \pi^2}{64\Omega_-^2} - \frac{\lambda^2}{4V^2} \sin^2 \left( \frac{V\pi}{2\Omega_-} \right).
$$

(35)

The first two terms in Eq. (35) correspond to the results of Ref. [18] where the off-resonant electron transitions to the state $(|A1\rangle - |B1\rangle)/\sqrt{2}$ were completely neglected. The third term arises due to account for such transitions. Their contribution to the electron state evolution (34) is proportional to the small parameter $\frac{\lambda}{V} \ll 1$ and results in the relative phase and amplitude shifts between the coefficients $c_{A1}$ and $c_{B1}$ that indicate on the finite tunnelling time $\tau_{tunn} \sim 1/V$ between the QDs. Moreover, the oscillations at the frequency $2\omega$ (so-called Bloch-Siegert oscillations) affect, to some extent, the ideal three-level oscillation picture.

The results obtained show that the three-level scheme can be used for the description of electron dynamics if the conditions $\frac{\lambda}{V} \ll 1$, $\frac{\Delta}{\omega} \ll 1$ are satisfied. The errors introduced due to the presence of a nearby forth level are of the order of $\frac{\lambda}{V}$.

**B. The two-level quantum dynamics.**

Next we study the opposite case of small $V$ when a substantial spreading of an excited electron between the QDs occurs after many Rabi oscillations in single QD.
have completed. Such an electron dynamics is realized in the double-dot structure where the tunnel coupling between the excited states of the QDs is rather small as compared to the electron-pulse coupling:

$$2V = |\delta_+ - \delta_-| \ll |\lambda_0|.$$  \hspace{1cm} (36)

The condition (36) may be satisfied for two different pulse designs, i.e. for both $|\lambda_0| \ll |\delta_-|, |\delta_+|$ and $|\delta_-|, |\delta_+| \ll |\lambda_0|$. The first inequality corresponds to the off-resonant single QD excitation whereas the second one characterizes the resonant two-level Rabi oscillations in the same QD. The off-resonant case does not reveal significant two-level dynamics since an electron stays predominantly localized in the state $|A_0\rangle$ (the population of the state $|A_1\rangle$ is of the order of $\lambda_0^2/\delta^2 \ll 1$). We shall focus our attention on the resonant transition for which the conditions

$$V, |\delta_-|, |\delta_+| \ll |\lambda_0|$$  \hspace{1cm} (37)

are satisfied.

In this case an electron oscillates between the ground and excited states of the QD A for the pulse durations $T \ll 1/V$. To prove this statement we make use of the approximations $\Omega_+ - \Omega_- \approx 2V|\delta|/|\lambda|$ and $e^{-iVt} \approx 1 - iT$ that holds for the time domain $t \ll 1/V$. Inserting them into the Eqs. (27)-(29) and retaining in the time dependencies the terms up to the first order in $Vt$, we have the following expression

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for the state vector:

\[
|\Psi\rangle = e^{-i\epsilon_0 t} \left\{ \left[ f_0^- (t) + i\frac{V}{2} e^{i\frac{\delta_-}{2} t} \cos (2\Omega_- t) \right] |A0\rangle + i\frac{V}{2} e^{i\frac{\delta_-}{2} t} \cos (2\Omega_- t) |B0\rangle \right\} + \\
+ e^{-i\epsilon_- t} \left\{ \left[ f_1^- (t) + \frac{1}{2\sqrt{2\Omega_-}} \frac{V}{2} e^{i\frac{\delta_-}{2} t} \sin (2\Omega_- t) \right] |A1\rangle - \frac{1}{2\sqrt{2\Omega_-}} \frac{V}{2} e^{i\frac{\delta_-}{2} t} \sin (2\Omega_- t) |B1\rangle \right\} 
\]

that demonstrates the two-level electron evolution in the QD A is slightly distorted by the excitations in the QD B. Setting in Eq. (38) \( V = 0, \delta_- = 0 \) we find the state-vector evolution for the two-level system being in the exact resonance with the external pulse

\[
|\Psi\rangle = e^{-i\epsilon_0 t} \left[ \cos \left( \frac{\lambda_0 t}{2} \right) - i\frac{1}{4\lambda_0} e^{-i2\omega t} \sin \left( \frac{\lambda_0 t}{2} \right) \right] |A0\rangle + \\
+ e^{-i\epsilon_- t} \left[ -i \sin \left( \frac{\lambda_0 t}{2} \right) + \frac{1}{4\lambda_0} \left( 1 - e^{-i2\omega t} \right) \cos \left( \frac{\lambda_0 t}{2} \right) \right] |A1\rangle.
\]

If one adopts the scheme where the quantum information is encoded into the ground and first excited electron (or exciton) states of the single QD (\([6], [8], [29]\)), the population transfer into neighboring QD should be considered as the information leakage from the computational subspace and the corresponding error probability may be evaluated with the help of Eq. (38). The computational error introduced by the Bloch-Siegert oscillations alone is deduced from Eq. (39).

**C. The electron excitation driven by the strongly detuned pulse.**

Finally, we shall examine the case of the Raman-like off-resonant excitation of an electron in the double-dot structure (\(|\lambda_0| \ll |\delta_\pm|\)). This mechanism of the optical quantum-state engineering is currently under extensive investigations because of important properties that distinguish it from the other optical schemes (see, e.g., \([30]\)). First, the excited (auxiliary) states are populated only virtually that allows
one to localize the electron population almost completely in the ground-state subspace \( \{ |A0\rangle, |B0\rangle \} \). Provided that the states \( |A0\rangle \) and \( |B0\rangle \) constitute the qubit basis one can therefore operate with the quantum information trapped in the logical subspace for any time. This, in its turn, simplifies the state evolution design and prevents the qubit from the decoherence induced by spontaneous photon emission from the excited levels. Second, the population transfer realized via the off-resonant excitations is quite robust against the pulse imperfections such as the uncontrollable detunings and the timing errors. The quantum optics provides one with wide class of schemes specially developed for those purposes. Quite recently several attempts have been made to adopt those schemes for the solid-state objects possessing of the atomic-like spectrum, e.g. the QDs \([21]\), the QDs combined with cavity QEDs \([7]\), the rf-SQUIDs \([31]\).

Some features of the qubit state evolution based on one-electron quantum dynamics in the symmetric double-dot structure driven by the strongly detuned pulse have been outlined in the work \([21]\). Here we consider this effect as the particular case of the four-level double-dot dynamics studied in Sec. II. Choosing the system parameters to satisfy the inequalities

\[
|\lambda_0| \ll |\delta_-|, |\delta_+| \ll \omega, \quad |\lambda_0| \ll V,
\]

we get from the general formulae (27)-(29) the following expression for the state
vector

\[ |\Psi\rangle \approx e^{-i\left(\varepsilon_0 + \frac{\delta_- + \delta_+}{4\delta_- \delta_+} \tilde{\lambda}^2\right)t} \left[ \cos \left( \frac{V}{2\delta_- \delta_+} \tilde{\lambda}^2 t \right) |A0\rangle + i \sin \left( \frac{V}{2\delta_- \delta_+} \tilde{\lambda}^2 t \right) |B0\rangle \right]. \]  

(41)

The Eq. (41) describes the two-level Rabi oscillations at the frequency \( \Omega_\delta = V \tilde{\lambda}^2 / 2\delta_- \delta_+ \) similar to those induced by the electrostatic field (see Eq. (14)). The Rabi frequencies of these processes are very different from each other, viz. \( \Omega_0 / \Omega_\delta \sim |\delta_- \delta_+|/\omega^2_{10} \ll 1 \). This makes the optically-driven oscillations more preferable for a qubit-state engineering due to their higher speed as compared with that of the electrostatic driving. As it follows from Eqs. (14) and (41) the quantum dynamics in both cases is frozen if \( V/|\lambda_0| \) approaches zero. This effect may be explained in terms of the destructive interference between the probability amplitudes of the hybridized states represented by the symmetric and antisymmetric superpositions of the excited states of isolated QDs. Since those states become nearly-degenerate with the decrease of the parameter \( V/|\lambda_0| \), their probability amplitudes sum up in QD A whereas they cancel one another in QD B. Despite of these states are empty during the population transfer, the importance of their assistance to the process becomes more clear from this analysis.

It is worth to note that the system evolution described by Eq. (41) cannot result in an arbitrary rotation of the qubit-state vector on the Bloch sphere since it contains only one time-dependent parameter \( \theta = \Omega_\delta t \) corresponding to the polar angle (the azimuthal angle is fixed and equals to \( \pi/2 \)). To overcome this obstacle one should break the symmetry of the structure and use at least two driving pulses with different
parameters to implement the desired rotation [20].

IV. NUMERICS

To illustrate the analytical results of Sec. II and Sec. III we have performed the numerical simulations of electron dynamics in our structure. The Eq. (5) with \( \lambda \) defined by the Eq. (16) was integrated within the time interval \( 0 \leq t \leq 3T \) (where \( T = \pi/\lambda_0 \)) for \( |\lambda_0|/\omega = 10^{-3} \) and \( \delta_\perp = 0 \). This choice of the pulse-structure parameters corresponds to that usually realized in the QD systems where \( \omega_{10} \sim 10^{-2} \) eV and \( |\lambda_0| \sim 10^{-5} \) eV for the pulse strength \( E_0 \sim 1 \div 10 \) V/cm. Since we are interested in demonstration of the transition from the three-level double-dot scheme to the two-level single-dot scheme, the ratio \( V/|\lambda_0| \) was varied from 0.01 to 10.

The numerical plots showing the time dependencies of the populations \( p_n = |c_n|^2, \ n = A0, B0, A1, B1 \) are presented in the Figs. 2 (a)-(d) for \( V/|\lambda_0| = 5; 1; 0.3; 0.05 \), respectively. For large but finite values of \( V/|\lambda_0| \) the three-level Rabi oscillations picture involving the states \( |A0\rangle, |B0\rangle, (|A1\rangle + |B1\rangle)/\sqrt{2} \) becomes non-ideal due to the excitation of the state \( (|A1\rangle - |B1\rangle)/\sqrt{2} \), see Eq. (34). This effect is clearer seen in the representation of isolated QD basis since the phase and amplitude shifts between \( p_{A1} \) and \( p_{B1} \) provide us with the measure characterizing the difference between the electron populations of the excited levels in the QDs A and B. We see that for \( V/|\lambda_0| = 5 \) (Fig. 2(a)) the results of Sec. III A may be still applied while for \( V/|\lambda_0| = 1 \) (Fig. 2(b)) the regular oscillation pattern is destroyed when \( t \geq 3T \) and
the identification of excitation process is hardly possible. Figure 2(b) clearly demonstrates the strong modulations of the optically-induced Rabi oscillations caused by the electron tunnelling if \( V \approx |\lambda_0| \). Further reduction of the ratio \( V/|\lambda_0| \) amounts to the qualitative changes in the population dynamics. Figure 2(c) indicates the importance of the single-QD processes even for \( V/|\lambda_0| = 0.3 \). When \( V/|\lambda_0| = 0.05 \) (Fig. 2(d)) we observe several almost ideal two-level Rabi oscillations in the QD A slightly modulated by residual dynamics in the QD B. The total reorganization of oscillation pattern that marks the transition from one excitation scheme to another occurs for \( V/|\lambda_0| \sim 0.01 \). The reduction of the Rabi frequency by a factor of 2 and the depopulation of the states belonging to the QD B are clearly seen from the numerical plots that confirms the results obtained above.

V. CONCLUSIONS

There are a lot of proposals for the qubit design that use the basic quantum properties of low-dimensional objects to encode, to process, and to store the quantum information. The existence of purely theoretical frameworks is of great importance since they allow us to capture the principal aspects of idealized evolution of the system under consideration and then to examine it further at the more profound level. An exact solution describing the qubit dynamics is often readily achieved due to the simplified structure of the model. It is therefore desirable to look for the model that would include the main features characterizing the qubit and, at the same time, enable the analytical treatment of the dynamical problem.
In this paper we have studied in detail the one-electron double-dot structure proposed as the candidate for a qubit implementation [18]. The quantum operations in the structure may be realized by applying the resonant electromagnetic pulse driving an electron between the QDs. Within this scheme we have generalized the results recently obtained for several different pulse-structure setups [18], [21], [28] and have pointed on some delicate aspects concerning their applicability which, to our knowledge, had never been clarified before. As we have shown, the efficiency of one or another scheme is conditioned by the value of the ratio between the matrix element of electron tunnelling and the matrix element of optical dipole transition.

The mathematical model of the one-electron excitation process has permitted to study the coherent evolution of the system beyond the rotating-wave approximation. The numerical results have confirmed those obtained analytically for the parameter choices corresponding to the three- and two-level dynamics.

The results presented in this paper may be also applied to the two-electron symmetric double-dot structure [32] and to the other systems possessing the same spectral properties, say, to the superconducting devices [31, 33]. Besides we suppose that the effect of the structure asymmetry on the electron dynamics [22] may be treated in the same way.

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Fig. 1. Schematics of the states for a single electron confined in the double-dot structure.
Fig. 2. The electron populations of the ground states $|A0\rangle$ (solid line), $|B0\rangle$ (solid dashed line) and the excited states $|A1\rangle$ (thin line), $|B1\rangle$ (thin dashed line) of the symmetric double-dot structure versus the pulse duration $T$ (in units of $\lambda_0^{-1}$) for the pulse-structure parameter choice $\delta_\perp = 0$, $|\lambda_0|/\omega = 10^{-3}$ and (a) $V/|\lambda_0| = 5$, (b) $V/|\lambda_0| = 1$, (c) $V/|\lambda_0| = 0.3$, (d) $V/|\lambda_0| = 0.05$. 
Fig. 2(a)
Fig. 2(b)
Fig. 2(c)
Fig. 2(d)