Entanglement and quantum state engineering
in the optically driven two-electron double-dot structure

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We study theoretically the quantum dynamics of two interacting electrons in the symmetric double-dot structure under the influence of the bichromatic resonant pulse. The state vector evolution is studied for two different pulse designs. It is shown that the laser pulse can generate the effective exchange coupling between the electron spins localized in different dots. Possible applications of this effect to the quantum information processing (entanglement generation, quantum state engineering) are discussed.

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I. INTRODUCTION

In recent years, the low-dimensional semiconductor structures containing a small number of electrons in the size-quantized conduction band have attracted much attention. The main reason for that interest is the progress in up-to-date technology that allows one to fabricate the nanostructures with high precision [1]. Along with the use of advanced techniques of population control [2], it makes the creation of macroatoms with desired properties a standard experimental tool at hand. There is a lot of possible applications of this field of solid state physics [3]. One of them is concerned with rapidly developing quantum computing and quantum information processing [4].

The key point of implementation of quantum computer’s hardware is to find an appropriate physical system characterized by the well-defined Hilbert space that would allow for efficient external (classical) control. Among the systems proposed for operation with the quantum information, a two-level system (qubit) is the most studied one. For practical applications, one has to look for a physical system that could serve as a base for a scalable quantum computer. Only such scalable quantum computers would outperform their classical counterparts for several classes of computational problems. It is commonly believed that the problem of scalability can be effectively solved with solid-state systems. Many proposals for solid-state qubit realization have been made. Here we mention the superconducting devices using a Cooper-pair box [5], phosphorous donors embedded in a silicon host [6, 7, 8],
and a wide class of the systems based on the quantum dots (QDs) (see, e. g., 9, 11, 12, 13, 17, 19, 20, 21, 22, 23).

Loss and DiVincenzo have proposed to encode the quantum information into electron spins contained in laterally coupled electrostatically-formed QDs 9. In their model, the local alternating magnetic field is used for single-qubit rotations, while the electrostatic gate voltage brings about the interqubit coupling. Besides, the static magnetic field can be applied to the double-dot structure to adjust the exchange coupling between electrons. Upon the control of the exchange during the voltage pulse action, the logic operations on the neighboring electron spins become possible. To organize a non-trivial two-qubit gate (e. g., XOR) one needs to perform the ”square-root of swap” operation combined with appropriate single-qubit rotations 10. Furthermore, the exchange control allows one to transform the product state of the qubits into the non-separable superposition of the qubit states which is the highly-entangled state 17, 18, 49, 50.

Here we present a novel scheme for manipulation with the quantum states of two interacting electrons in the double-dot structure. It is based on the interaction of electrons with the coherent electromagnetic pulse. We are interested in the design of the external driving field that allows to achieve the reliable control on the quantum dynamics of the system. As we shall show, this may be done by the generalization of the driving schemes developed for a single electron confined in the double-dot structure 13, 14, 15, 16, 18. The sequence of laser pulses, instead of electrostatic
ones, is applied to the double-dot structure, and the all-optical quantum state engineering is realized via the optically induced transitions between the size-quantized two-electron levels. The resonant character of the electron-pulse interaction provides the high selectivity of the corresponding transitions, thus making the scheme robust against the unwanted excitations far from the resonance. We consider the resonant transitions between two degenerate ground states and the auxiliary excited states. We use the ground (localized) states of the system as qubit states. This allows us to isolate them from each other after the pulse is off. Besides, the scheme permits one to operate with the strongly-detuned pulses in the Raman-like regime, where one (or several) auxiliary state remains unpopulated. This seems to be important for the structures with the strong decoherence. We shall see that two resonant pulses are sufficient for implementation of such a rotation in the subspace spanned by the ground states that corresponds to generation of the entanglement of electron spins. Thus, the two-electron double-dot structure may be exploited as the spin entangler.

The paper is organized as follows. In Sec. II we study the two-electron double-dot structure interacting with the external pulse. The energy spectrum and the stationary eigenstates of the structure are obtained from the analysis of the extended Hubbard model. The probability amplitudes of the eigenstates relevant for the quantum dynamics are found from the solution of the non-stationary Schrödinger equation. In Sec. III we describe the quantum operations that may be realized in this system under the influence of the laser pulses and propose the scheme for the
entanglement generation. The results are summarized in Sec. IV.

II. THE MODEL

A. The eigenstates and the eigenenergies of the stationary Hamiltonian

We consider the double-dot structure (see Fig.1) containing two interacting electrons 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 $\varphi_{A(B)\,0}(r) = \langle r \mid A(B)\,0 \rangle$ and $\varphi_{A(B)\,1}(r) = \langle r \mid A(B)\,1 \rangle$, 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 tunneling $[13]$.

Following the procedure used in Ref. [25], we transform the one-electron orbitals of the isolated QDs into the orthonormal one-electron orbitals, accounting for the hybridization of the excited states: $\tilde{\varphi}_{A0}(r) = \varphi_{A0}(r), \tilde{\varphi}_{B0}(r) = \varphi_{B0}(r), \tilde{\varphi}_{A1}(r) = (1/2\sqrt{1+s} + 1/2\sqrt{1-s})\varphi_{A1}(r) + (1/2\sqrt{1+s} - 1/2\sqrt{1-s})\varphi_{B1}(r)$, $\tilde{\varphi}_{B1}(r) = (1/2\sqrt{1+s} + 1/2\sqrt{1-s})\varphi_{B1}(r) + (1/2\sqrt{1+s} - 1/2\sqrt{1-s})\varphi_{A1}(r)$. Here $s = \langle A1 \mid B1 \rangle$ is the overlap between the excited states of the QDs.
The Hamiltonian of two interacting electrons confined in the symmetric double-dot structure is $H_0 = h(\mathbf{r}_1) + h(\mathbf{r}_2) + w(|\mathbf{r}_1 - \mathbf{r}_2|)$ \([10, 25]\), where $h(\mathbf{r}_i)$ is the one-particle Hamiltonian that includes the kinetic and potential terms, and $w(|\mathbf{r}_1 - \mathbf{r}_2|)$ is the Coulomb interaction between the electrons. We consider the case that the electrons have the opposite spins. The Hamiltonian can be expressed in terms of the extended Hubbard model \([25]\):

$$H_0 = \sum_{\sigma} \left[ \varepsilon_0 (n_{A0,\sigma} + n_{B0,\sigma}) + \varepsilon_1 (n_{A1,\sigma} + n_{B1,\sigma}) \right] - \sum_{\sigma} \left( a_{A1,\sigma}^+ a_{B1,\sigma} + h.c. \right) +$$

$$+ \sum_{\sigma \neq \sigma'} \left[ V_{00} n_{A0,\sigma} n_{B0,\sigma'} + V_{01} (n_{A0,\sigma} n_{B1,\sigma'} + n_{B0,\sigma} n_{A1,\sigma'}) + V_{11} n_{A1,\sigma} n_{B1,\sigma'} \right] +$$

$$+ U_{00} (n_{A0,\uparrow} n_{A0,\downarrow} + n_{B0,\uparrow} n_{B0,\downarrow}) + U_{11} (n_{A1,\uparrow} n_{A1,\downarrow} + n_{B1,\uparrow} n_{B1,\downarrow}) +$$

$$+ U_{01} \sum_{\sigma \neq \sigma'} (n_{A0,\sigma} n_{A1,\sigma'} + n_{B0,\sigma} n_{B1,\sigma'}) ,$$

where $a_{k,\sigma}^+$ creates an electron in the state with the wave function $\tilde{\varphi}_k(\mathbf{r})$ ($k = A0, B0, A1, \text{and} B1$) and the spin $\sigma = +1/2, -1/2$; $n_{k,\sigma} = a_{k,\sigma}^+ a_{k,\sigma}$ is the particle number operator acting on the state vectors in the occupation number representation; $\varepsilon_0 = \int \tilde{\varphi}_{A(B)0}^* (\mathbf{r}) h(\mathbf{r}) \tilde{\varphi}_{A(B)0}(\mathbf{r}) d\mathbf{r}$ and $\varepsilon_1 = \int \tilde{\varphi}_{A(B)1}^* (\mathbf{r}) h(\mathbf{r}) \tilde{\varphi}_{A(B)1}(\mathbf{r}) d\mathbf{r}$ are the one-electron energies of the ground and excited states, respectively (the same for both QDs); $\tilde{t} = \int \tilde{\varphi}_{A1}^* (\mathbf{r}) h(\mathbf{r}) \tilde{\varphi}_{B1}(\mathbf{r}) d\mathbf{r}$ is the matrix element for the electron hopping between the excited states of the QDs; $V_{00} = V_{A0,B0}$, $V_{01} = V_{A0,B1} = V_{A1,B0}$, $V_{11} = V_{A1,B1}$, $U_{01} = U_{A0,A1} = U_{B0,B1}$ are the Coulomb interaction energies for electrons occupying different states, and $U_{00} = U_{A0,A0} = U_{B0,B0}$, $U_{11} = U_{A1,A1} = U_{B1,B1}$ are the Coulomb interaction energies for electrons occupying the same orbital state.

For the sake of simplicity we have omitted in Eq. (1) the higher-order terms such as the cotunneling, the direct exchange, e.t.c. In what follows, we consider the
system in the strong-confinement regime where the Coulomb correlations are small compared to the level spacing, i.e., $V_{00}, V_{01}, V_{11} \ll U_{00}, U_{01}, U_{11} \ll \varepsilon_1 - \varepsilon_0$.

The straightforward diagonalization of the Hamiltonian results in the two-electron eigenstates $|n\rangle$ ($n = 1 - 16$) that can be expressed in terms of the four-site basis vectors $|n_{A0,\sigma}\rangle \otimes |n_{B0,\sigma'}\rangle \otimes |n_{A1,\sigma''}\rangle \otimes |n_{B1,\sigma'''\rangle} = |n_{A0,\sigma}, n_{B0,\sigma'}, n_{A1,\sigma''}, n_{B1,\sigma'''}\rangle$ as follows:

\begin{align*}
|1\rangle &= |1_\uparrow, 1_\downarrow, 0, 0\rangle, \\
|2\rangle &= |1_\downarrow, 1_\uparrow, 0, 0\rangle, \\
|3\rangle &= |2, 0, 0, 0\rangle, \\
|4\rangle &= |0, 2, 0, 0\rangle, \\
|5\rangle &= C_+ \left( v + \sqrt{\nu^2 + t^2} \right) |1_\uparrow, 0, 0, 1_\downarrow\rangle + \tilde{t} |1_\uparrow, 0, 1_\downarrow, 0\rangle, \\
|6\rangle &= C_+ \left( v + \sqrt{\nu^2 + t^2} \right) |0, 1_\downarrow, 1_\uparrow, 0\rangle + \tilde{t} |0, 1_\downarrow, 0, 1_\uparrow\rangle, \\
|7\rangle &= C_+ \left( v + \sqrt{\nu^2 + t^2} \right) |0, 1_\downarrow, 1_\downarrow, 0\rangle + \tilde{t} |0, 1_\downarrow, 0, 1_\downarrow\rangle, \\
|8\rangle &= C_+ \left( v + \sqrt{\nu^2 + t^2} \right) |1_\downarrow, 0, 0, 1_\uparrow\rangle + \tilde{t} |1_\downarrow, 0, 1_\uparrow, 0\rangle, \\
|9\rangle &= C_- \left( v - \sqrt{\nu^2 + t^2} \right) |1_\uparrow, 0, 0, 1_\downarrow\rangle + \tilde{t} |1_\uparrow, 0, 1_\downarrow, 0\rangle, \\
|10\rangle &= C_- \left( v - \sqrt{\nu^2 + t^2} \right) |0, 1_\downarrow, 1_\uparrow, 0\rangle + \tilde{t} |0, 1_\downarrow, 0, 1_\uparrow\rangle, \\
|11\rangle &= C_- \left( v - \sqrt{\nu^2 + t^2} \right) |0, 1_\downarrow, 1_\downarrow, 0\rangle + \tilde{t} |0, 1_\downarrow, 0, 1_\downarrow\rangle, \\
|12\rangle &= C_- \left( v - \sqrt{\nu^2 + t^2} \right) |1_\downarrow, 0, 0, 1_\uparrow\rangle + \tilde{t} |1_\downarrow, 0, 1_\uparrow, 0\rangle, \\
|13\rangle &= \tilde{C}_+ \left[ \tilde{v} + \sqrt{\tilde{\nu}^2 + 4t^2} \right] \left[ \frac{[0,0,1_\uparrow,1_\downarrow]}{\sqrt{2}} + \frac{[0,0,1_\downarrow,1_\uparrow]}{\sqrt{2}} \right] + \tilde{t} \left[ \frac{[0,0,2,0]}{\sqrt{2}} + \frac{[0,0,0,2]}{\sqrt{2}} \right], \\
|14\rangle &= \frac{[0,0,1_\uparrow,1_\downarrow]}{\sqrt{2}} - \frac{[0,0,1_\downarrow,1_\uparrow]}{\sqrt{2}}, \\
|15\rangle &= \frac{[0,0,2,0]}{\sqrt{2}} - \frac{[0,0,0,2]}{\sqrt{2}}, \\
|16\rangle &= \tilde{C}_- \left[ \tilde{v} - \sqrt{\tilde{\nu}^2 + 4t^2} \right] \left[ \frac{[0,0,1_\uparrow,1_\downarrow]}{\sqrt{2}} + \frac{[0,0,1_\downarrow,1_\uparrow]}{\sqrt{2}} \right] + \tilde{t} \left[ \frac{[0,0,2,0]}{\sqrt{2}} + \frac{[0,0,0,2]}{\sqrt{2}} \right],
\end{align*}

where

\begin{align*}
C_\pm &= \frac{1}{\sqrt{t^2 + (v \pm \sqrt{\nu^2 + t^2})^2}}, \\
\tilde{C}_\pm &= \frac{1}{\sqrt{4\tilde{t}^2 + (\tilde{v} \pm \sqrt{\tilde{\nu}^2 + 4\tilde{t}^2})^2}},
\end{align*}

and

\begin{align*}
u &= \frac{U_{01} + V_{01}}{2}, \quad v = \frac{U_{01} - V_{01}}{2}, \\
\tilde{\nu} &= \frac{U_{11} + V_{11}}{2}, \quad \tilde{v} = \frac{U_{11} - V_{11}}{2}.
\end{align*}
The corresponding eigenenergies $E_n \ (n = 1 - 16)$ are $E_1 = E_2 = 2\varepsilon_0 + V_{00}$, $E_3 = E_4 = 2\varepsilon_0 + U_{00}$, $E_5 = E_6 = E_7 = E_8 = \varepsilon_0 + \varepsilon_1 + u - \sqrt{v^2 + t^2}$, $E_9 = E_{10} = E_{11} = E_{12} = \varepsilon_0 + \varepsilon_1 + u + \sqrt{v^2 + t^2}$, $E_{13} = 2\varepsilon_1 + \tilde{u} - \sqrt{\tilde{v}^2 + 4\tilde{t}^2}$, $E_{14} = 2\varepsilon_1 + V_{00}$, $E_{15} = 2\varepsilon_1 + U_{01}$, and $E_{16} = 2\varepsilon_1 + \tilde{u} + \sqrt{\tilde{v}^2 + 4\tilde{t}^2}$.

Since we suppose $V_{00}$, $V_{01}$, $V_{11} << U_{00}$, $U_{01}$, $U_{11} << \varepsilon_1 - \varepsilon_0$, the coupling between the electrons gives rise to three energy manifolds around the values of the energy of two non-interacting electrons, $2\varepsilon_0$, $\varepsilon_0 + \varepsilon_1$, and $2\varepsilon_1$. The eigenstates $|1\rangle$ and $|2\rangle$ with the energy $\varepsilon_{00} = E_1 = E_2$ are the ground states of the two-electron system, where both electrons are localized in the lowest orbital levels of different QDs. The states $|3\rangle$ and $|4\rangle$, where the electrons occupy the lowest orbital level in the same QD, have the energy $\tilde{\varepsilon}_{00} = E_3 = E_4$. Next eight eigenstates correspond to the situation where one electron is localized in one of the QDs and another electron is excited and delocalized over the double-dot structure due to tunneling. Four of those states, namely $|5\rangle$, $|6\rangle$, $|7\rangle$, and $|8\rangle$, constitute the subspace with the energy $\varepsilon_{01}^- = E_5 = E_6 = E_7 = E_8$, whereas other four states, i.e. $|9\rangle$, $|10\rangle$, $|11\rangle$, and $|12\rangle$, form the subspace with the energy $\varepsilon_{01}^+ = E_9 = E_{10} = E_{11} = E_{12}$. The states $|13\rangle$, $|14\rangle$, $|15\rangle$, and $|16\rangle$ are the states where both electrons are excited. Their properties have been studied in Refs. [10, 24, 25, 27, 39]. The eigenstate $|13\rangle$ is the singlet state with the energy $\varepsilon_{11}^s = E_{13}$. The triplet state $|14\rangle$ has the energy $\varepsilon_{11}^t = E_{14}$. Two remaining eigenstates, $|15\rangle$ and $|16\rangle$, are the singlet states with the energies $\varepsilon_{11}^a = E_{15}$ and $\varepsilon_{11}^{as} = E_{16}$, respectively. The superscript $s(a)$ denotes the
symmetry (antisymmetry) of the corresponding state under the spatial reflection operation.

We see that the states $|1\rangle$ and $|2\rangle$ are degenerate. It follows from the fact that both tunneling and exchange processes are prohibited due to the strong confinement. Thus if the system was initialized, say, in the state $|1\rangle$, we assert that an electron in the state $|A0(B0)\rangle$ has the spin $\sigma = 1/2(-1/2)$ \[24\]. Instead, the states $|13\rangle - |16\rangle$ are hybridized, and electron spins in those states are correlated. In the next section, we shall work with the state $|13\rangle$. We shall show that together with the doubly-occupied states $|3\rangle$ and $|4\rangle$, it can be used to entangle the electrons, initially decoupled from each other. To achieve this goal, one needs an effective coupling mechanism for the controlled interaction between the electrons.

B. Laser-induced electron dynamics

The standard method usually considered for electron charge and spin manipulations in a double-dot structure is based on the electrostatic and magnetic field control. As was proposed in Ref. \[10\], by applying an adiabatically switched voltage pulse one can rise or lower the potential barrier between QDs. This allows one to organize the controlled electron-electron interaction and, as a consequence, the quantum state engineering. The main difficulty inherent to this method lies in the trade-off between the requirement of the small operation times, as compared to the decoherence time, and the adiabatic character of the electron tunneling process. In particular, the violation of adiabatic conditions will lead to the double-occupancy
of one of the QDs, resulting in the quantum information leakage from the computational subspace. To solve this problem, several approaches have been developed, see Refs. [26, 27].

Here we propose an alternative scheme of the controlled quantum dynamics of the two-electron system. It is based on the resonant laser pulses that induce the transitions between the electron states in the double-dot structure. In recent works it was shown that similar techniques may be utilized to drive an excess electron between two QDs [13, 14, 15, 16]. For the exciton-based quantum computers the interqubit coupling schemes based on the coherent dynamics of localized [23, 28, 29, 30, 31] and delocalized [32, 33] excitons under the influence of appropriately tuned laser pulses were developed. The superexchange coupling between two deep donor impurities in silicon mediated by the optically-excited electron of the control atom was discussed in Ref. [34]. In our model we consider the electron dynamics involving the two-electron states (Eq. (2)) and show how to choose the pulse parameters to obtain the final quantum state with the desired properties.

Let our system be initially in the state \( |1 \rangle \). To drive it into the superposition of the states \( |1 \rangle \) and \( |2 \rangle \), we make use of two resonant laser pulses that act simultaneously and induce the optical dipole transitions between each of those states and the auxiliary excited states. As we shall show, the proper choice of the pulse parameters, such as the intensities, the detunings of pulse frequencies from the resonance, and the pulse durations, allows one to realize the inversion operation and create the
maximally entangled states, e. g., the Bell states. There are several ways to achieve this goal. We study two of them, the most transparent in our opinion.

We consider the quantum evolution of the system under the influence of two laser pulses. The model Hamiltonian has the form

$$H = H_0 + [V_1 \cos (\omega_1 t + \varphi_1) + V_2 \cos (\omega_2 t + \varphi_2)] [\theta (t) - \theta (t - T)] ,$$

where $V_k = -e E_k (r_1 + r_2) , k = 1, 2; E_k, \omega_k, \text{ and } \varphi_k$ are the amplitude, the frequency and the phase of the $k$-th pulse, respectively, $e$ is the electron charge, $\theta (t)$ is the step function and $T$ is the pulse duration.

The state vector of the system may be represented in terms of the eigenstates $|1\rangle - |16\rangle$, Eqs. (2), of the stationary Hamiltonian $H_0$ as

$$|\Psi (t)\rangle = \sum_{n=1}^{16} c_n (t) \exp (-iE_n t) |n\rangle$$

(hereafter $\hbar = 1$).

The quantum evolution of the state vector under the influence of two laser pulses is governed by the non-stationary Schrödinger equation

$$i \frac{\partial |\Psi (t)\rangle}{\partial t} = H |\Psi (t)\rangle .$$

Since the electron-pulse dipole interaction gives rise to the one-particle excitations only, we cannot generate an entanglement in a two-electron system using the one-photon process only. However, it is seen from the energy level structure that the states $|1\rangle$ and $|2\rangle$ can be coupled by the four-stage two-photon transition scheme
involving the ground states, the states from the subspace with the energy \( \varepsilon_{01}^\pm \), and
the hybridized \( |13\rangle \) - \( |16\rangle \) or doubly-occupied \( |3\rangle, |4\rangle \) states.

First, we consider the transition scheme that involves the states \( |1\rangle, |2\rangle, |5\rangle \) - \( |8\rangle \),
and \( |13\rangle \) (Fig. 2). Because of the degeneracy of the states \( |5\rangle \) - \( |8\rangle \), two pulses are
sufficient to induce the transitions between the states \( |1\rangle \) and \( |2\rangle \). We set \( \omega_1 = \varepsilon_{01} - \varepsilon_{00} + \Delta_1 \) and \( \omega_2 = s^{11} - \varepsilon_{01} + \Delta_2 \), where \( \Delta_1 \) and \( \Delta_2 \) are the detunings. The
transformation of the coefficients \( c_n \) in Eq. (6) according to \( c_1 = \tilde{c}_1 \), \( c_2 = \tilde{c}_2 \), \( c_k = \tilde{c}_k \exp (-i \Delta_1 t) \) for \( k = 5 - 8 \), \( c_{13} = \tilde{c}_{13} \exp [-i (\Delta_1 + \Delta_2) t] \) and their substitution
into Eqs. (6) and (7) result in the set of equations:

\[
\begin{align*}
\dot{\tilde{c}}_1 &= \lambda_1 (\tilde{c}_5 + \tilde{c}_6), \\
\dot{\tilde{c}}_2 &= \lambda_1 (\tilde{c}_7 + \tilde{c}_8), \\
\dot{\tilde{c}}_k &= -\Delta_1 \tilde{c}_k + \lambda_1^* (\delta_{k5} \tilde{c}_1 + \delta_{k6} \tilde{c}_1 + \delta_{k7} \tilde{c}_2 + \delta_{k8} \tilde{c}_2) + \lambda_2 \tilde{c}_{13}, \quad k = 5 - 8 \\
\dot{\tilde{c}}_{13} &= - (\Delta_1 + \Delta_2) \tilde{c}_{13} + \lambda_2^* (\tilde{c}_5 + \tilde{c}_6 + \tilde{c}_7 + \tilde{c}_8),
\end{align*}
\]

(8)

where \( \lambda_1 = \frac{1}{2} \exp (i \varphi_1) \langle 1 \mid V_1 \mid 5 \rangle \) and \( \lambda_2 = \frac{1}{2} \exp (i \varphi_2) \langle 5 \mid V_2 \mid 13 \rangle \) are the matrix
elements of the electron-pulse interaction (here we take into account the identities
\( \langle 1 \mid V_1 \mid 5 \rangle = \langle 1 \mid V_1 \mid 6 \rangle = \langle 2 \mid V_1 \mid 7 \rangle = \langle 2 \mid V_1 \mid 8 \rangle \) and \( \langle 5 \mid V_2 \mid 13 \rangle = \langle 6 \mid V_2 \mid 13 \rangle = \langle 7 \mid V_2 \mid 13 \rangle = \langle 8 \mid V_2 \mid 13 \rangle \)).

The eigenfrequencies of the set of Eqs. (8) can be found from the following
equation:

\[
x^3 - i (2 \Delta_1 + \Delta_2) x^2 + \left[ 4 |\lambda_2|^2 + 2 |\lambda_1|^2 - \Delta_1 (\Delta_1 + \Delta_2) \right] x - 2 i |\lambda_1|^2 (\Delta_1 + \Delta_2) = 0.
\]

(9)

The solution of Eq. (9) is straightforward. However, only in the case that \( \Delta_1 = -\Delta_2 = \Delta \) it may be presented in a rather simple form. This case corresponds to
the two-photon resonance.

If the system is initially in the superposition of the states $|1\rangle$ and $|2\rangle$, i.e., $|\Psi(0)\rangle = \alpha |1\rangle + \beta |2\rangle$, the initial conditions are $\tilde{c}_1(0) = \alpha$, $\tilde{c}_2(0) = \beta$, $\tilde{c}_{k\neq 1,2}(0) = 0$ and we obtain from the set of Eqs. (8) the following expressions for the coefficients $c_n$ in the laboratory frame:

$$c_{1,2} = \pm \frac{\alpha - \beta}{2} \exp\left(i \frac{\Delta t}{2}\right) \left[ \cos(\Omega_1 t) - i \frac{\Delta}{2|\lambda_1|^2} \sin(\Omega_1 t) \right] + \frac{\alpha + \beta}{2} \left\{ \frac{2|\lambda_2|^2}{|\lambda_1|^2 + 2|\lambda_2|^2} \right\} \exp\left(i \frac{\Delta t}{2}\right) \left[ \cos(\Omega_2 t) - i \frac{\Delta}{2|\lambda_2|^2} \sin(\Omega_2 t) \right], \quad (10)$$

$$c_5 = c_6 = -i \exp\left(i \frac{\Delta t}{2}\right) \left[ \frac{\alpha - \beta}{2} \frac{\lambda_1^*}{\Omega_1} \sin(\Omega_1 t) + \frac{\alpha + \beta}{2} \frac{\lambda_2^*}{\Omega_2} \sin(\Omega_2 t) \right], \quad (11)$$

$$c_7 = c_8 = -i \exp\left(i \frac{\Delta t}{2}\right) \left[ -\frac{\alpha - \beta}{2} \frac{\lambda_1^*}{\Omega_1} \sin(\Omega_1 t) + \frac{\alpha + \beta}{2} \frac{\lambda_2^*}{\Omega_2} \sin(\Omega_2 t) \right], \quad (12)$$

$$c_{13} = (\alpha + \beta) \frac{|\lambda_1|^2|\lambda_2|^2}{|\lambda_1|^2 + 2|\lambda_2|^2} \left\{ -1 + \exp\left(i \frac{\Delta t}{2}\right) \left[ \cos(\Omega_2 t) - \frac{\Delta}{2|\lambda_2|^2} \sin(\Omega_2 t) \right] \right\}. \quad (13)$$

Here $\Omega_1 = \sqrt{\frac{\Delta^2}{4} + 2|\lambda_1|^2}$ and $\Omega_2 = \sqrt{\frac{\Delta^2}{4} + 2\left(|\lambda_1|^2 + 2|\lambda_2|^2\right)}$ are the Rabi frequencies.

If $\Delta = 0$, the system is in the exact resonance with both pulses. To localize the system completely in the ground-state subspace $\{|1\rangle, |2\rangle\}$, the matrix elements of the electron-pulse interaction $\lambda_1$ and $\lambda_2$ must satisfy the condition

$$\frac{|\lambda_2|}{|\lambda_1|} = \sqrt{\frac{1}{2} \left( \frac{4m^2}{k^2} - 1 \right)}, \quad (14)$$

where $k, m$ are integers chosen so that the right-hand side of Eq. (14) be real. The Eq. (14) ensures that $c_{n\neq 1,2}(T_k) = 0$. We see that in this case for the pulse durations $T_k = \pi k/|\lambda_1| \sqrt{2}$ the inversion operation $\sigma_x$ is realized in the ground-state subspace if $k$ is odd, while the initial state is not changed if $k$ is even.
There is an interesting particular case of $\Delta = \left(\varepsilon_{11} + \varepsilon_{00}\right)/2 - \varepsilon_{01}$ similar to that discussed in Ref. [28]. In this case $\omega_1 = \omega_2$ and only one pulse is sufficient for the state inversion. If the value of $\Delta$ is much larger than the matrix elements of the electron-pulse interaction, the populations $|c_k|^2$ of the auxiliary states $|5\rangle$ - $|8\rangle$ are of the order of $(|\lambda_1|/\Delta)^2 << 1$, i.e., much smaller than the populations of states $|1\rangle$, $|2\rangle$ and $|13\rangle$. It corresponds to the well-known adiabatic elimination procedure, widely used in the quantum optics (see, e.g., Ref. [45]). The auxiliary states $|k\rangle$ with $k = 5 - 8$ are populated only virtually, and the effective three-level scheme is realized, giving rise to the simultaneous excitation of two electrons from the ground-state subspace to the state $|13\rangle$. We may view this process as the generation of the effective exchange coupling between the electron spins driven by one resonant pulse. The corresponding dynamics can be obtained through the expansion Eqs. (10) - (13) in terms of the small parameters $|\lambda_{1,2}|/|\Delta| << 1$. This process is slow compared to that with $\Delta = 0$, since its Rabi frequency is $\Omega_{eff} = |\lambda_1|^2/|\Delta| << |\lambda_1|$. The stroboscopical evolution of the state vector in the ground state subspace $|\Psi (T_k)\rangle = \exp (-i\varepsilon_{00}T_k) U_k |\Psi (0)\rangle$, where $|\Psi (0)\rangle = (\alpha, \beta)^T$ and $T_k = \pi k \Delta / \left(|\lambda_1|^2 + 2 |\lambda_2|^2\right)$, is given by the matrix $U_k$:

$$U_k = \exp (-i\psi_k) \begin{pmatrix}
    \cos (\psi_k) & \exp (i\pi/2) \sin (\psi_k) \\
    -\exp (-i\pi/2) \sin (\psi_k) & \cos (\psi_k)
\end{pmatrix},$$

(15)

where $\psi_k = \Omega_{eff} T_k = \pi k |\lambda_1|^2 / \left(|\lambda_1|^2 + 2 |\lambda_2|^2\right)$. Here the pulse duration $T_k$ is determined by the condition $c_{13} (T_k) = 0$. We see that Eq. (15) corresponds to the in-plane rotation through the angle $\psi_k$. This operation creates the maximally
entangled state for the single-spin qubit encoding scheme of Ref. [9] but, of course, is not sufficient for an arbitrary rotation of the qubit state vector on the Bloch sphere for the scheme where the states $|1\rangle$ and $|2\rangle$ are used as qubit states (see below).

Next we consider the situation where $\Delta_1 \neq \Delta_2$ and one (or both) of the transitions is off-resonant. Assuming $\Delta_1 = 0$, $|\Delta_2| >> |\lambda_2|$, we exclude adiabatically the singlet state from the set of Eqs. (8) as follows:

$$\tilde{c}_{13} = 0, \quad \tilde{c}_{13} = \frac{\lambda_2^*}{\Delta_2} (\tilde{c}_5 + \tilde{c}_6 + \tilde{c}_7 + \tilde{c}_8),$$

(16)

and substitute Eq. (16) into Eq. (8), thus arriving at the following expressions:

$$c_{1,2} = \pm \frac{\alpha - \beta}{2} \cos \left( \sqrt{2} |\lambda_1| t \right) + \frac{\alpha + \beta}{2} \exp \left( -i \frac{2|\lambda_2|^2 t}{\Delta_2} \right) \left[ \cos \left( \tilde{\Omega}_2 t \right) + i \frac{2|\lambda_2|^2}{\Delta_2 \tilde{\Omega}_2} \sin \left( \tilde{\Omega}_2 t \right) \right],$$

(17)

$$c_5 = c_6 = -i \sqrt{2} \left[ \frac{\alpha - \beta}{2} \sin \left( \sqrt{2} |\lambda_1| t \right) + \frac{\alpha + \beta \sqrt{2} \lambda_1^*}{\tilde{\Omega}_2} \exp \left( -i \frac{2|\lambda_2|^2 t}{\Delta_2} \right) \sin \left( \tilde{\Omega}_2 t \right) \right],$$

(18)

$$c_7 = c_8 = -i \sqrt{2} \left[ -\frac{\alpha - \beta}{2} \sin \left( \sqrt{2} |\lambda_1| t \right) + \frac{\alpha + \beta \sqrt{2} \lambda_1^*}{\tilde{\Omega}_2} \exp \left( -i \frac{2|\lambda_2|^2 t}{\Delta_2} \right) \sin \left( \tilde{\Omega}_2 t \right) \right],$$

(19)

where $\tilde{\Omega}_2 = \sqrt{\frac{4|\lambda_2|^4}{\Delta_2^2} + 2 |\lambda_1|^2}$. As for the cases discussed above, there is a relation between the matrix elements of the electron-pulse interaction necessary to achieve a complete localization of the system in the subspace spanned by the states $|1\rangle$ and $|2\rangle$:

$$\frac{|\lambda_2|}{|\lambda_1|} = \frac{|\Delta_2|}{\sqrt{2} |\lambda_2|} \sqrt{\frac{m^2}{k^2} - 1},$$

(20)

where $m, k$ are integers. If this condition is fulfilled, the ground state evolution
matrix at $T_k = \pi k / |\lambda_1| \sqrt{2}$ is given by the expression:

$$U_k = \exp (i \pi k) \exp (-i \psi_k) \begin{pmatrix} \cos (\psi_k) & \exp (-i \pi/2) \sin (\psi_k) \\ -\exp (i \pi/2) \sin (\psi_k) & \cos (\psi_k) \end{pmatrix},$$

(21)

with $\Omega_{e\text{ff}} = |\lambda_2|^2 / |\Delta_2|$, $\psi_k = \Omega_{e\text{ff}} T_k = \pi \sqrt{m^2 - k^2} / 2$ if $m - k$ is even and $\bar{U}_k = \exp (i \pi) \sigma_x U_k = \exp (i \pi) U_k \sigma_x$ if $m - k$ is odd.

The cases $\Delta_2 = 0$, $|\Delta_1| >> |\lambda_1|$, $|\Delta_2| >> |\lambda_1|$, $|\lambda_2| > |\lambda_1|$ enable a complete localization of the system in the ground state subspace at any time without restrictions on the matrix elements of the electron-pulse interaction. However, the former case does not reveal the non-trivial evolution up to the fourth order in the ratio $|\lambda_1| / \Delta_1$, and the effective Rabi frequency in the latter case is $\Omega_{e\text{ff}} \sim |\lambda_1|^2 |\lambda_2|^2 / \Delta_1^2 |\Delta_2|$ that makes the rotations defined by Eq. (15) too slow and unviable in view of spin decoherence processes. Here we don’t consider these cases in details.

Next we study another way to achieve a rotation of the quantum state of the two-electron system. As it was mentioned before, the doubly-occupied states $|3\rangle$ and $|4\rangle$ may also be exploited for the creation of the effective exchange coupling between the electrons localized in $|1\rangle$ or $|2\rangle$ states. To see how it may be realized we present the scheme for the electron transitions that provides such coupling (Fig. 3). We see that only the states $|1\rangle - |8\rangle$ are involved in the process, while the doubly-excited states $|13\rangle - |16\rangle$ do not participate the dynamics. In complete analogy with the procedure described above, we represent the state vector of the system in terms of the stationary eigenstates of Eq. (1) with the time-dependent coefficients $c_k$, where $k = 1 - 8$. Next we substitute it into Eq. (7), where the frequencies of two laser
pulses now are given by the expressions \( \omega_1 = \varepsilon_{01} - \varepsilon_{00} + \Delta_1 \) and \( \omega_2 = \varepsilon_{01} - \varepsilon_{00} - \Delta_2 \).

The coefficients may be obtained from the set of equations:

\[
\begin{align*}
    i \dot{\tilde{c}}_1 &= \lambda_1 (\tilde{c}_5 + \tilde{c}_6), \\
    i \dot{\tilde{c}}_2 &= \lambda_1 (\tilde{c}_7 + \tilde{c}_8), \\
    i \dot{\tilde{c}}_3 &= - (\Delta_1 + \Delta_2) \tilde{c}_3 + \tilde{\lambda}_2^* (\tilde{c}_5 + \tilde{c}_8), \\
    i \dot{\tilde{c}}_4 &= - (\Delta_1 + \Delta_2) \tilde{c}_4 + \tilde{\lambda}_2^* (\tilde{c}_6 + \tilde{c}_7), \\
    i \dot{\tilde{c}}_k &= - \Delta_k \tilde{c}_k + \lambda_k^* (\delta_{k5} \tilde{c}_1 + \delta_{k6} \tilde{c}_1 + \delta_{k7} \tilde{c}_2 + \delta_{k8} \tilde{c}_4) \\
    &+ \tilde{\lambda}_2 (\delta_{k5} \tilde{c}_3 + \delta_{k6} \tilde{c}_3 + \delta_{k7} \tilde{c}_4 + \delta_{k8} \tilde{c}_4), \quad k = 5 - 8,
\end{align*}
\]

where now \( \tilde{\lambda}_2 = \frac{1}{2} \exp (i \varphi_2) \langle 5 | V_1 | 3 \rangle \) and the identities \( \langle 5 | V_2 | 3 \rangle = \langle 8 | V_2 | 3 \rangle = \langle 6 | V_2 | 4 \rangle = \langle 7 | V_2 | 4 \rangle \) are assumed.

The set of Eqs. (22) is equivalent to the set of Eqs. (8) if we set in Eqs. (22) \( \tilde{c}_3 = \tilde{c}_4 = \tilde{c}_{13} / \sqrt{2} \) and \( \tilde{\lambda}_2 = \lambda_2 \sqrt{2} \). This substitution allowed in the case \( c_{13}(0) = 0 \), \( c_3(0) = 0 \), \( c_4(0) = 0 \) provides the formal analogy between the first and the second schemes. Thus one can derive all time-dependent coefficients \( c_k \) relevant for the quantum dynamics from Eqs. (10) - (21).

Note that apart from the restrictions imposed by the state localization requirement, the matrix elements of the electron-pulse interaction, as well as the detunings from resonance, have to satisfy another condition, justifying the resonant approximation used in Eqs. (8) and Eqs. (22). If we consider the transitions involving the states \( |5\rangle - |8\rangle \) or the states \( |9\rangle - |12\rangle \), the following inequalities must be fulfilled:

\[ |\Delta_1|, |\Delta_2| < \lambda_1, |\lambda_2| < \varepsilon_{01}^+ - \varepsilon_{01}^- \]. For the off-resonant transition we should keep \( |\Delta| < \varepsilon_{01}^+ - \varepsilon_{01}^- \). (Of course, the energy differences of the states \( |13\rangle - |16\rangle \) must satisfy similar inequalities as well.) If we take the matrix elements of the electron-pulse interaction \( |\lambda_{1,2}| \sim 10^{-5} \) eV and assume that \( \tilde{t} \sim 10^{-3} \) eV and \( \tilde{t} \sim \varepsilon_{01}^+ - \varepsilon_{01}^- \),
it is possible to meet the condition above in both resonant and off-resonant cases. In the resonant case, this choice of the structure parameters implies the operation times \( \tau \sim 1/\min(|\lambda_1|,|\lambda_2|) \) to be of the order of tens of picoseconds. The setting of the matrix elements of the electron-pulse interaction \( \lambda_{1,2} \) can be readily achieved via adjusting the strengths of the pulses.

In the next section, we consider several important potential applications of the results obtained to the quantum information processing.

III. ENTANGLEMENT AND QUANTUM STATE ENGINEERING

Electron spin localized in the QD is now extensively studied as the promising candidate for the solid-state qubit implementation. Let us investigate the possibility of exploiting the electron spins in the double QD structure for quantum computations in view of the realization of quantum operations by optical means. We discuss the situation where the states \(|1\rangle = |1\uparrow, 1\downarrow, 0, 0\rangle\) and \(|2\rangle = |1\downarrow, 1\uparrow, 0, 0\rangle\) are used as the qubit logical states \([46]\). Thus, a double-dot structure is now viewed as a single qubit. What kind of quantum operations may be performed on such a qubit? We have found from Eqs. (15) and (21) that the rotation along the meridian \( \varphi = \pi/2 \) through the polar angle \( \psi_k \) may be realized. It is easy to see, however, that the states \(|1\rangle \) and \(|2\rangle \) belong to the same charge configuration. To generate, say, \( \sigma_z \) operation one should distinguish between them, making use of the Pauli exclusion principle only. The charge and spin degrees of freedom are decoupled from each other, and
the laser pulse cannot change the electron spin in a given QD. To construct a quantum gate that acts directly on the electron charge only but is able to operate with the spin-encoded qubit, some additional tools are required. For example, one can exploit the exciton-based techniques [32], where the circularly-polarized laser pulse generates the electron transition to the conduction band depending on the spin of the electron occupying the lowest size-quantized level.

With the help of the optics it is possible to transport an individual spin through the sequence of the QDs by swapping electron spins in the neighboring QDs, induced by the laser pulse that generates \( \sigma_x \) operation. This effect may be helpful for the quantum state transfer in the models based on the Loss - DiVincenzo’s proposal.

There exist several more spin qubit encoding schemes [47] - [50]. We can arrive at these schemes by organizing the entanglement between the QD spins (see below), that is equivalent to the superposition of the states \( |1\rangle \) and \( |2\rangle \), and using the entangled states as logical ones. This offers an opportunity to handle with quantum information within the decoherence-free subspaces that, in its turn, provides robust quantum information processing. Note that in the four-spin encoding scheme of Ref. [48] both \( \sigma_x \) and \( \sigma_z \) operations can be performed via the sequence of the corresponding two-spin inversion operations.

Next we discuss the possible application of the optically-induced quantum evolution of the two-electron system in the entanglement generation. A lot of proposals based on a few-electron QD, using the static and/or alternating electric field con-
trolling techniques, have been made to achieve this purpose [35] - [44]. Some of them seem to be very promising, especially those handling with two interacting electrons in the double QD under the action of an oscillatory electric field. As it was demonstrated, both static [40] - [44] and alternating [35] - [39] electric fields satisfying some conditions may be used to entangle two electrons and localize them in the entangled state. Besides, by an appropriate switching procedure one can drive the system between the delocalized and fully localized states and coherently destroy the electron tunneling process. In those schemes, the orbital electron states localized in the left (right) QD serve as the logical states. The Coulomb repulsion plays a significant role in such systems; depending on the Hubbard ratio $t/U$, several regimes may be realized. All of the dynamical properties can be derived through the analysis of the Floquet spectrum of the two-electron system in the electric field.

We present one another way to produce an entanglement in the two-electron double-dot structures. If we treat the states $|1\rangle = |1_{\uparrow}, 1_{\downarrow}, 0, 0\rangle$ and $|2\rangle = |1_{\downarrow}, 1_{\uparrow}, 0, 0\rangle$ as the two-qubit states in the Loss - DiVincenzo’s scheme, it is easy to create the maximally entangled Bell state of two spins $(|1\rangle + i |2\rangle)/\sqrt{2}$ as follows. One may use the single pulse with the detuning $\Delta = (\epsilon_{11} + \epsilon_{00})/2 - \epsilon_{01}$ and $|\lambda_2| = \sqrt{3/2} |\lambda_1|$. The evolution described by Eq. (15) is then realized, and if the system has started from the state $|1\rangle$, the Bell state (up to the common phase) is obtained at $\tau_{\text{Bell}} = \pi |\Delta|/4|\lambda_1|^2$. Note that this state may be realized in two ways considered in Sec. II B. Besides of the formation of the spin entangled states,
there is also a possibility to generate the entanglement between the charge states, as it was introduced in Refs. \[35\], \[39\]. To do this one should start from the superposition of the states like \((|1⟩ + |2⟩)/\sqrt{2}\) or from one of the doubly-occupied states. The entangled states of two different forms, \((|2,0,0,0⟩ ± |0,2,0,0⟩)/\sqrt{2}\) and \((|0,0,2,0⟩ ± |0,0,0,2⟩)/\sqrt{2}\), can be obtained by the use of the corresponding transition scheme. If one needs to obtain the state similar to that discussed in Ref. \[39\], i.e. \((|0,0,2,0⟩ + |0,0,0,2⟩)/\sqrt{2}\), one should replace the state \(|13⟩\) by the state \(|16⟩\) in the transition scheme of Eqs. (10) - (13) and tune the system parameters so that the regime with \(\tilde{t} << \tilde{v}\) be realized. Then the state \(|16⟩\) turns out to be the equally-weighted superposition of the doubly-occupied excited states. Of course, the quantum state engineering procedure retains all of the operation tools, e.g., Coulomb repulsion control proposed earlier \[35\]. On the other hand, the entanglement of the doubly-occupied states \(|3⟩\) and \(|4⟩\) may be achieved in the transition scheme presented by Eqs. (22) without any modification.

Here we draw attention to the main distinguishing features of our model from those mentioned above. First, we consider the two-level QDs instead of the single-level QDs studied in Refs. \[35\] - \[44\]. This enables us to localize an electron in the well-isolated ground-state subspace without use of any additional resource. Second, we operate with the pulse intensities much smaller as compared to those used in the cited works. It would help one to carefully isolate the logical and auxiliary states from other ones during the pulse action. The resonant character of the structure-
field interaction serves for the same purpose. What we would like to mention, our scheme turns out to be useful in the entanglement generation for both spin and charge encoding schemes.

**IV. CONCLUSIONS**

We have considered the quantum dynamics of two interacting electrons confined in the double-dot structure. The quantum transitions occur between the eigenstates of the stationary two-electron Hamiltonian under influence of the resonant laser pulse. By an appropriate choice of the pulse parameters we may generate the superposition of some states. There are several ways for the quantum state manipulations by the use of the electromagnetic field. All of them involve the hybridized or doubly-occupied states as the auxiliary states. Those states play an important role in the non-trivial quantum state evolution and serve as "the entanglers" in the schemes presented here. The maximally-entangled Bell states may be generated for both electron spin states and electron charge states. We can say that the laser pulse generates (and destroys) the effective exchange coupling between the electron spins. This makes the two-electron double-dot system very interesting for the quantum information processing. Several important quantum operations may be constructed using this system. If we consider the doubly-degenerate ground states (or their combinations) as the logical states, the optically-driven single qubit operation $\sigma_x$ may be realized in the logical subspace. In this work we have studied a simplified model to catch the principal features of the behavior of two confined interacting
electrons in the resonant field. The next steps of investigations are the study of the influence of structure asymmetry and the pulse imperfections on the quantum dynamics as well as the decoherence effects. However, one should expect that further analysis will retain all of qualitative results obtained here. We hope that our study will be helpful for the solid-state quantum computer design.

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Fig. 1. Schematics of the states for a single electron confined in the double-dot structure. These one-electron states are used for construction of the two-electron states, see text for details.
Fig. 2. Transition scheme connecting the states $|1\rangle$ and $|2\rangle$ through the use of the auxiliary states $|5\rangle$, $|6\rangle$, $|7\rangle$, $|8\rangle$, and $|13\rangle$ in the case of the two-photon resonance $\Delta_1 = -\Delta_2 = \Delta$. Here $\Delta_1$ and $\Delta_2$ are the detunings of the pulse frequencies $\omega_1$ and $\omega_2$, respectively, from the exact resonance.
Fig. 3. Transition scheme connecting the states $|1\rangle$ and $|2\rangle$ through the use of the auxiliary states $|3\rangle$, $|4\rangle$, $|5\rangle$, $|6\rangle$, $|7\rangle$, and $|8\rangle$ in the two-photon resonant case, see Fig. 2.