Coherent optical manipulation of triplet-singlet states in coupled quantum dots

Hakan E. Türeci,1 J. M. Taylor,2 and A. Imamoglu1

1Institute of Quantum Electronics, ETH-Zürich, CH-8093, Zürich, Switzerland
2Department of Physics, MIT, MA 02139, USA

(Dated: March 23, 2022)

We show that spin-orbit coupling in a quantum dot molecule allows for coherent manipulation of two electron spin states using Raman transitions. Such two-electron spin states defined by the singlet and triplet states of two exchange coupled quantum dots can have favorable coherence properties. In addition, two of the four metastable ground states in this system can be used as auxiliary states that could facilitate implementation of tasks such as mapping of spin states to that of a single propagating photon. We find that even weak spin-orbit effects — manifesting themselves as slightly different g-factors for the electron and the hole — would allow for the coherent Raman coupling of the singlet-triplet states. We also discuss the possibilities for implementing quantum optical techniques for spin preparation and manipulation.

I. INTRODUCTION

Over the past decade semiconductor quantum optical systems, implemented in quantum wells and particularly quantum dots, have been paradigmatic for the exploration of novel quantum mechanical effects in the solid state. Potential applications as single photon sources or as quantum bits for quantum information storage and processing have driven investigations of single quantum dots with ground states containing zero or a single electron charge. These systems bear marked resemblance to noble gas and alkali atoms. However, systems similar to alkaline-earth atoms (with two valence electrons) or homo-nuclear alkali molecules remain largely unexplored. Such systems typically have a more complex fine structure, leading to metastable spin states with useful decoherence properties and have well-established semiconductor realizations.

In this article we examine approaches for optically coupling the four metastable ground (spin) states of a two-electron double quantum dot system via optical Raman transitions. These states, split into a singlet and triplet manifold, have demonstrated useful properties with respect to spin-related dephasing, as seen in recent experiments in electrically-controlled double quantum dots. Our approach for coupling singlet and triplet states via optical fields can lead to the integration of optical manipulation, measurement and entanglement techniques with demonstrated approaches to controlling fine structure states in electrical quantum dots.

While we focus on the case of Zinc-Blende (III-V) semiconductor quantum dots, we find that optical coupling of ground state spins can be realized even with weak spin-orbit interaction, where a sufficient condition is that the electron and hole states have differing g-factors; this is believed to be the case for small radius carbon nanotubes. Thus our approach for working with fine structure states can find wide application in a variety of quantum dot systems.

II. COHERENT OPTICAL MANIPULATION OF COUPLED QUANTUM DOTS WITH STRONG-SPIN ORBIT EFFECTS

In this section we develop an approach to coupling singlet and triplet fine structure states of a double quantum dot system via optical Raman transitions. We rely upon the dramatic difference between exchange energies for two electrons on the same quantum dot and two electrons in separated quantum dots to provide a reliable means of using single spin selection rules to develop controlled, two-spin selection rules. We find that for doubly-charged double quantum dots with spin-orbit coupling, such as Zinc-Blende semiconductor quantum dots, polarization- and energy-selective transitions between all fine structure states are possible. This allows techniques, such as STImulated Raman Adiabatic Passage (STIRAP) to be used to initialize arbitrary superpositions of singlet and triplet spin states.

We consider here the optical transitions in a doubly-charged coupled quantum dot (CQD), a situation that can for instance be realized with an asymmetric pair of stacked InAs quantum dots embedded in a Schottky diode structure in an appropriate gate voltage regime. We assume that the left dot (L) is blue-shifted with respect to the right dot (R) and that the lowest conduction level of the right dot is detuned by Δ with respect to the left one (see Fig. 1 inset). Higher orbital states can be neglected as the associated transitions are well separated in energy in the gate voltage regime of interest. For a range of gate voltages where the two dots contain a total of two electrons, by fine tuning the voltage it is possible to convert between atomic and molecular orbital states, i.e., between charge states (0, 2), (2, 0) and (1, 1). Labels (m, n) here refer to the number of electrons confined in the (left, right) dot.
presence of finite interdot tunneling and in the regime where (1, 1) is the lowest energy charge configuration, the two resident electrons hybridize resulting in an energetically isolated singlet-triplet subspace where the effective degree of freedom is the total spin of the two electrons. This regime is denoted by (II) in Fig. 1 and in what follows we will work in this regime.

In the (1,1) regime (II), the ground state manifold is given by the states $| (1,1)S \rangle = \frac{1}{\sqrt{2}}(c_{L1}^\dagger c_{R1}^\dagger + c_{L1}^\dagger c_{R1}^\dagger)|0\rangle$, $| (1,1)T_0 \rangle = \frac{1}{\sqrt{2}}(c_{L1}^\dagger c_{R1}^\dagger + c_{L1}^\dagger c_{R1}^\dagger)|0\rangle$, $| (1,1)T_+ \rangle = c_{L1}^\dagger c_{R1}^\dagger|0\rangle$ and $| (1,1)T_- \rangle = c_{L1}^\dagger c_{R1}^\dagger|0\rangle$ with energies $E_{S_0} = \Delta + e_{c}^{LR} - J$, $E_{T_0}, T_\pm = \Delta + E_{c}^{LR}$, where $J \approx T_e^2/(E_{C}^{RR} - E_{C}^{LR})$. The typical energies for an InAs self-assembled quantum dot are given by $E_{C}^{LL} \approx E_{C}^{RR} = V_{c}^{ee,LL,LL} \approx 20$ meV, $E_{C}^{LR} = V_{c}^{ee,LL,RR} \approx 10$ meV, where $V_{ab} = \int d\mathbf{r} \epsilon_{\sigma}^{a}(\mathbf{r})\epsilon_{\sigma}^{b}(\mathbf{r})V_{c}(\mathbf{r} - \mathbf{r}')\phi_{\sigma}^{a}(\mathbf{r}')\phi_{\sigma}^{b}(\mathbf{r})$ and $\phi_{\sigma}^{a}(\mathbf{r})$ are the single-particle envelope wavefunctions on dot $j = L, R$ for conduction band electrons ($a = e$) or holes ($a = h$). The tunneling matrix element $T_e$ appearing in the exchange splitting $J$ is given by $T_e = t_e + V_{c}^{ee,LL,RR} \approx t_e + V_{c}^{ee,RR,RR} \lesssim 1$ meV and connects $| (1,1)S \rangle$ to $| (2,0)S \rangle$ and $| (0,2)S \rangle$. Here, $t_e$ is the bare interdot electron tunneling matrix element.

Consider now a right-hand circularly polarized ($\sigma_+$) optical excitation with its optical axis along the heterostructure growth direction ($z$-axis). This axis is defined with respect to the crystal axes of the quantum well structure on which the dots are grown, which typically defines the axis of shape asymmetry of the quantum dots (see Fig. 2). The light-matter interaction Hamiltonian in the dipole, rotating wave and envelope function approximations is given by

$$V_+ = \frac{i\hbar}{2} \Omega_+ \left( M_{RR} c_{R1}^\dagger h_{R\uparrow}^\dagger + M_{LR} c_{L1}^\dagger h_{L\uparrow}^\dagger \right) e^{-i\omega_+ t} + h.c. \quad (1)$$

Here, $\Omega_+ = \frac{1}{\sqrt{2}}(L^x = 0, L^z = 0 \mid x + iy \mid L^h = 1, L^z = +1)$, where $| L^x = 0, L^z = 0 \rangle$ is the periodic part of the conduction band Bloch wavefunction which has $s$-character, and $| L^h = 1, L^z = +1 \rangle$ is that of the valence band electrons (or holes) which has $p$-character, while $M_{RR}$ and $M_{LR}$ are the overlaps of the electron and hole envelope wavefunctions $\zeta(L|R)_h$ and $\zeta(L|R)_h$ respectively. The implicit assumption here is that the light-hole levels and the spin-orbit split-off hole-band is energetically well-separated from the heavy-hole band denoted by $| \uparrow \rangle = | J^h = 3/2, J_z^h = 3/2, \rangle$, $| \downarrow \rangle = | J^h = 3/2, J_z^h = -3/2, \rangle (J^h = L^h + S^h)$ so that their optical coupling can be neglected, a well-justified approximation in Zinc-Blende (III-V) semiconductor quantum dots. Note that it is the correlation between the spin and spatial parts of the heavy-hole states ($| \uparrow \rangle = | L^h_{\uparrow} = 1; \uparrow \rangle$, $| \downarrow \rangle = | L^h_{\downarrow} = -1; \downarrow \rangle$) which enables us to use selection rules based on pseudospin conservation. The excited state manifold of doubly charged excitons ($X^{2-}$) is
FIG. 2: Schematics showing the optical excitation of the CQD system embedded in a Schottky-diode structure. The light pulse \( \Omega_+ (t) \) is incident along the crystal growth direction (\( z \)-axis).

The action of \( \sigma_+ \) optical excitation on the \((1,1)\) ground state manifold

\[
\begin{align*}
V_+ |S\rangle &= -\mathcal{M}_{RR} |(1,2)\uparrow\rangle - \mathcal{M}_{LR} |(2,1)\downarrow\rangle \\
V_+ |T_0\rangle &= -\mathcal{M}_{RR} |(1,2)\uparrow\rangle + \mathcal{M}_{LR} |(2,1)\downarrow\rangle \\
V_+ |T_+\rangle &= \mathcal{M}_{RR} |(1,2)\uparrow\rangle - \mathcal{M}_{LR} |(2,1)\downarrow\rangle \\
V_+ |T_-\rangle &= 0
\end{align*}
\]

where \( \sigma_+ = \uparrow, \downarrow \). We assume that because of the particular structure of the dots and the bias of choice, the optically generated hole always resides on the right dot within its lifetime. Here, \( E_{\text{LLR}} = E_{\text{LL},LL} + 2E_{\text{LL},RR} - E_{\text{LR},LR} - 2V_{\text{ee}}^{LR,RR} - V_{\text{eh}}^{LR,RR} \) for instance and \( E_{\text{LLR}} > E_{\text{LLRR}} \) because of the e-h attraction. An additional electron tunneling matrix element connects states \(|(2,1)\uparrow\rangle\) and \(|(1,2)\uparrow\rangle\) as well as \(|(2,1)\downarrow\rangle\) and \(|(1,2)\downarrow\rangle\), which gives rise to an anti-crossing at around the bias \( \Delta \approx E_{\text{LLR}} - E_{\text{LR}} \approx 15 - 20 \) meV. If we operate in the \((1,1)\) regime (II) close to \( V = V_0 \), we can safely neglect any mixing of these \( X^{2-} \) states.

Consider the action of \( \sigma_- \) optical excitation on the \((1,1)\) ground state manifold

\[
\begin{align*}
E_1 &= E_{\text{LLR}} \\
E_2 &= E_{\text{LLR}} \\
E_3 &= E_{\text{LLRR}} + \Delta \\
E_4 &= E_{\text{LLRR}} + \Delta
\end{align*}
\]

The e-h exchange interaction in the relevant optically excited states would be negligibly small since the unpaired electron is exchange coupled to the hole that resides in a different dot. Starting with \(|(1,1)S, T_0\rangle\), the states \(|(1,2)\uparrow\rangle\) and \(|(1,2)\downarrow\rangle\) will form the "bright" excitons \( X^{2-} \), whereas the states \(|(1,2)\uparrow\rangle\) and \(|(1,2)\downarrow\rangle\) will form the bright excitons for the \(|(1,1)T_\pm\rangle\) subspace.

Since \( M_{RR} \gg M_{LR} \), the upper row transitions are strongest, while the lower row may be neglected.

The above table illustrates the optical selection rules for transitions from \(|(1,1)S\rangle\) and \(|(1,1)T_0\rangle\).

\[
\begin{array}{c|c|c}
\sigma_+ & |\uparrow\rangle & |\downarrow\rangle & E_1 \\
\hline
M_{RR} & (1,2) & (1,2) & E_1 \\
M_{LR} & (2,1) & (2,1) & E_3
\end{array}
\]

Since \( M_{RR} \gg M_{LR} \), the upper row transitions are strongest, while the lower row may be neglected.

The \( \sigma_- \) optical excitation on the \((1,1)\) ground state manifold via optical transitions for circular polarization. One can envision manipulating the system with a combination of a static external magnetic field and optical fields to overcome this problem.

Consider applying an in-plane magnetic field \( B = B_0 \mathbf{z} \). This will mix both the ground state and the excited state manifolds. Noting that the in plane hole \( g \)-factor is negligible, instead of rewriting the new states in the
FIG. 3: Diagram showing the optical selection rules from the singlet-triplet ground state of the CQD to the intermediate excited states in a Raman spin-flip scheme. While \(|(1,1)S\rangle | (1,1)T_0\rangle\) couple to the \(|M| = 1\) excited state manifold, \(|(1,1)T_+\rangle | (1,1)T_-\rangle\) couple to the \(|M| = 2\) subspace. Here \(M\) refers to the spin projection of the total pseudo-spin of the left electron and the right hole.

\[
\begin{align*}
(1,1)S & \quad \Omega_+^+ \quad (1,1)T_0 \quad \Omega_+^-
(1,1)T_+ & \quad \Omega_+^+ \quad (1,1)T_- \quad \Omega_-
\end{align*}
\]

FIG. 4: Diagram showing the selection rules in the presence of an in-plane magnetic field. The previously derived selection rules for \(B = 0\) are no longer valid in this case, allowing for either \(|T_+\rangle\) or \(|T_-\rangle\) state to be coupled to the same intermediate state as that of \(|T_0\rangle\) and \(|S\rangle\). The two sets of ground state manifolds can be accessed by using either left or right hand circularly polarized fields.

Note that, \(\sigma_+\) polarized light couples \(|(1,1)S\rangle\) and \(|(1,1)T_0\rangle\) to \(|\downarrow_L\rangle\) as well. In this regard, it is important to have a Zeeman splitting \(\epsilon_Z \leq J\) (\(\epsilon_Z = g^e \mu_B B\) is the single electron Zeeman splitting) that is sufficiently large so that coupling of a single spin state to a single intermediate excited optical state may be possible.

Having an auxiliary state at disposal for optical manipulation is important in the implementation of robust qubit rotations based on stimulated Raman adiabatic passage (STIRAP)\(^{14}\). The qubits in our CQD scheme above are formed by \((1,1)S\) and \((1,1)T_0\), which is a subspace that can be protected from hyper-fine induced dephasing by spin-echo techniques\(^{11}\). Combined with the immunity to dephasing of the intermediate state that can be achieved through classical optical control, it allows for the realization of fault-tolerant quantum computation.
by STIRAP, the above described scheme seems to be well suited for generation and manipulation of qubit states for quantum information protocols in Zinc-Blende (III-V) semiconductor quantum dots.

III. OPTICAL MAPPING OF SPIN STATES

In this section we discuss a scheme to efficiently prepare, manipulate and map spin states of a doubly charged CQD into photon polarization for long distance quantum state transfer.

Consider the CQD system placed inside a high-Q cavity and the gate voltage tuned such that the system is in regime (II), but close to the anti-crossing at \( V = V_1^* \). We furthermore apply an in-plane magnetic field. This will ensure that the states \( S, T_0 \) and \( T_\pm \) are well-separated in energy. Let us assume that the system is initialized to the superposition state

\[
|\Psi\rangle = \alpha |S\rangle + \beta |T_0\rangle
\]

in the qubit space composed of \((|S\rangle, |T_0\rangle)\). We first map this state into the state

\[
|\Psi\rangle = \alpha |T_-\rangle + \beta |T_0\rangle
\]

by a STIRAP sequence using \( \sigma_- \) polarized light as described for example in Ref.\(^\text{16}\). The Raman nature of the scattering will ensure a robust state mapping.

We next turn on lasers with frequencies \( \omega_a, \omega_b \), polarizations \( \sigma_- , \sigma_+ \) and time dependent Rabi frequencies \( \Omega_a(t) \) and \( \Omega_b(t) \), which are detuned by \( \delta_1 \) and \( \delta_2 \) from \(|T_-\rangle - |\uparrow_L\psi\rangle\) and \(|T_0\rangle - |\uparrow_L\phi\rangle\) transitions respectively. The detunings \( \delta_1 \) and \( \delta_2 \) are carefully chosen so that the transitions \(|\uparrow_L\psi\rangle - |S\rangle\) and \(|\uparrow_L\phi\rangle - |S\rangle\) are resonant with a pair of degenerate cavity modes \( \omega_- \) and \( \omega_+ \) with orthogonal polarizations (see Fig. 5). Such cavities can be engineered using photonic band gap structures where the cavity-mode splitting can be precisely tuned by AFM oxidation so as to enforce the degeneracy condition required\(^\text{17}\). The two-photon resonance condition for both polarizations will result in cavity assisted Raman transition down to \(|S\rangle\) and a transfer of the original spin state \((\alpha |T_-\rangle + \beta |T_0\rangle)\) to the polarization of the emitted single photon via

\[
(\alpha |T_-\rangle + \beta |T_0\rangle)_{ph} \rightarrow |S\rangle (\alpha |1_{\sigma_-}\rangle_{ph} + \beta |1_{\sigma_+}\rangle_{ph})
\]

The resulting photons can then be used for long-range quantum communication in a distributed network of nodes containing CQD-cavity systems. For this to work reliably, we require the Purcell-enhanced decay rate to be faster than spin dephasing rates. Experimental parameters\(^\text{18}\) suggest this can be achieved by a factor of 100.
IV. COHERENT OPTICAL MANIPULATION OF COUPLED QUANTUM DOTS WITH QUASI-DEGENERATE VALENCE BANDS

We now show that coherent optical coupling between fine structure states of doubly-charged double quantum dots could be achieved even when the intrinsic spin-orbit coupling of the semiconductor material is weak. We will rely upon an external magnetic field and weak electron-hole symmetry breaking (differing electron and hole $g$-factors) to show that all relevant Raman transitions are accessible.

Consider the possibility of optical manipulation of quantum dots where the spin-orbit effects are weak. The weakened selection rules result in the following optical interaction Hamiltonian

$$V_+ = \frac{i\hbar}{2} \Omega_{\text{RR}} \left( h_{R,3/2,3/2}^+ e_{R,1/2,-1/2}^+ + \frac{1}{\sqrt{3}} h_{R,3/2,1/2}^+ e_{R,1/2,1/2}^+ - \sqrt{\frac{2}{3}} h_{R,1/2,1/2}^+ e_{R,1/2,1/2}^+ \right) e^{-i\omega_+ \tau} + \text{h.c.}$$  \hspace{1cm} (6)

Here, $h_{R,J,J}^\dagger$ creates a hole on the right dot with total angular momentum $J^h = L^h + S^h$ and magnetic quantum number $J_z$, while $e_{R,1/2,\sigma}^\dagger$ creates an electron in the conduction band with spin projection $\sigma = \pm 1/2$. We only keep the strong, direct transition terms of the optical Hamiltonian. Considering now the action of $\sigma_+$ optical excitation on the $(1,1)$ ground state manifold, we obtain

$$V_+ |S\rangle = \mathcal{M}_{\text{RR}} \left( (1 \downarrow, 2 \downarrow) 3/2, 3/2 ) + \frac{1}{\sqrt{3}} (1 \uparrow, 2 \uparrow) 3/2, 1/2 ) - \sqrt{2/3} (1 \uparrow, 2 \uparrow) 1/2, 1/2 ) \right)$$

$$V_+ |T_0\rangle = \mathcal{M}_{\text{RR}} \left( -(1 \downarrow, 2 \downarrow) 3/2, 3/2 ) + \frac{1}{\sqrt{3}} (1 \uparrow, 2 \uparrow) 3/2, 1/2 ) - \sqrt{2/3} (1 \uparrow, 2 \uparrow) 1/2, 1/2 ) \right)$$

$$V_+ |T_+\rangle = -\mathcal{M}_{\text{RR}} \left( (1 \uparrow, 2 \uparrow) 3/2, 3/2 ) \right)$$

$$V_+ |T_-\rangle = \mathcal{M}_{\text{RR}} \left( \frac{1}{\sqrt{3}} (1 \downarrow, 2 \downarrow) 3/2, 1/2 ) - \sqrt{2/3} (1 \downarrow, 2 \downarrow) 1/2, 1/2 ) \right)$$

FIG. 6: The energy level structure of the excited state manifold. On the left, the energies are given in units of $\mu_B B_x$. The additional smaller splittings are due to the possibly different $g$-factors of electrons and holes, i.e. nonzero $\epsilon^{h,+}$. The states are written out explicitly in the format $|L^h_x, S^h_x\rangle |S^e_x\rangle$ where the electron spin refers to the left electron and the hole is on the right.

The doubly occupied singlet electrons in the right quantum dot are not shown.

In principle, the degeneracy in the excited state manifold presents a problem for coherent optical protocols due to potential destructive quantum interference between different pathways. To overcome this difficulty, we consider the effect of a magnetic field. For simplicity we assume a Voigt configuration ($B = B_x \hat{z}$, say). Writing the Zeeman Hamiltonian as

$$H_Z = \mu_B B_x \cdot ( L^h_x + g^h S^h_x + g^e S^e_x )$$  \hspace{1cm} (7)

where $L^h_x$ and $S^h_x$ are the $x$-components of the hole orbital and spin angular momentum respectively. Let us take $g^{h,e} = 2 + \epsilon^{h,e}$ with the hindsight that the spin-orbit interaction is weak. Then, one can easily see that the total
angular momentum is not conserved. Thus we have to use the \((L^h, L^g, S^h, S^g, S^e)\) basis. The energy level structure is given in Fig. 6. In this basis, the optical interaction Hamiltonian is given by

\[
V_+ = \frac{i\hbar}{2} \Omega_+ M_{RR} \left( H_{\uparrow}^\dagger e_{R\uparrow}^\dagger + H_{\downarrow}^\dagger e_{R\downarrow}^\dagger \right) + \text{h.c.}
\]

where

\[
H_{\sigma} = \frac{1}{2} h_{R,-1,\sigma} - \frac{1}{\sqrt{2}} h_{R,0,\sigma} - \frac{1}{2} h_{R,+1,\sigma} \quad \sigma = \uparrow, \downarrow
\]

The optical ground state is still given as above with the magnetic quantum numbers measured along \(x\) and with additional Zeeman energy contributions. The Zeeman contributions also split up the excited state manifold as shown in Fig. 6. Consider for instance the transitions to the final state \(|f\rangle = |-1, \downarrow\rangle_h|\uparrow\rangle_e = e_{R\uparrow}^\dagger e_{R\downarrow}^\dagger h_{R,-1,\downarrow}^\dagger\). We get

\[
\langle f|V_+|S\rangle = 1/2
\]

\[
\langle f|V_+|T_0\rangle = 1/2
\]

\[
\langle f|V_+|T_+\rangle = 0
\]

\[
\langle f|V_+|T_-\rangle = 0
\]

However, if one is to consider a spin-flip Raman scattering from \(S\) to \(T_0\) there are two available paths via intermediate states \(|f_1\rangle = |-1, \downarrow\rangle_h|\uparrow\rangle_e\) and \(|f_2\rangle = |-1, \uparrow\rangle_h|\downarrow\rangle_e\). Only when the electron and hole \(g\)-factors are different would these contributions allow a non-zero transition amplitude, otherwise the two available paths will interfere destructively. Hence it should in principle be possible to couple the states \(S, T_0\) to each other via two-photon processes if the electron and hole \(g\)-factors are different. In practice, the \(g\)-factors should be sufficiently different to ensure \(|\mu_B B_x(g^h - g^g)| > \gamma\), where \(\gamma\) is the broadening of the optically-excited states.

V. CONCLUSION

We have presented a scheme for optical manipulation of the metastable singlet-triplet subspace of a doubly-charged CQD. We find that with strong as well as weak spin-orbit interactions it is possible to implement Raman spin-flip transitions between a singlet and a triplet fine-structure-split ground state. Such Raman transitions enable the implementation of arbitrary coherent rotations robustly via quantum optical techniques previously proposed for single atoms or ions.

A key issue is the identification of the conditions that need to be satisfied in order to generate a spin-flip Raman scattering within a singlet-triplet space. In the absence of spin-orbit coupling and a degenerate p-like valence band, the optical Hamiltonian is a spin-0 operator. Thus, unless the spin-conservation law is broken in the intermediate state, it will not be possible to flip \(S\) to \(T_0\). In the strong spin orbit case, the optical interaction Hamiltonian however is a reducible spin operator with non-zero projections on spin-1 and spin-0 subspaces (where spin is actually a pseudo-spin due to the restriction to the heavy-hole band). This particular property makes it possible to optically connect \(S\) and \(T_0\).

We expect that our findings will stimulate experimental research aimed at combining electrical and optical manipulation of confined spin states. The fact that optical manipulation is possible even for quantum dot structures with weak spin-orbit interaction enhances the prospects for pursing experimental realization of quantum communication protocols in such systems.

Acknowledgments

HET would like to thank Mete Atatüre, Jan Dreiser and Alex Högele for useful discussions.

---

1 D. D. Awschalom, N. Samarth, and D. Loss, eds., Semiconductor Spintronics and Quantum Computation (Springer-Verlag, Berlin, 2002).
2 P. Michler, A. Kiraz, C. Becher, W. V. Schoenfeld, P. M. Petroff, L. D. Zhang, E. Hu, and A. Imamoglu, Science 290, 2282 (2000).
3 M. Pelton, C. Santori, J. Vuckovic, B. Y. Zhang, G. S. Solomon, J. Plant, and Y. Yamamoto, Physical Review Letters 89 (2002).
4 A. Imamoglu, Fortschr. Phys. 48, 987 (2000).
5 P. C. Haljan, P. J. Lee, K.-A. Brickman, M. Acton, L. Deslauriers, and C. Monroe, Physical Review A 72, 062316 (2005).
6 H. J. Krenner, M. Sabathil, E. C. Clark, A. Kress, D. Schuh, M. Bichler, G. Abstreiter, and J. J. Finley, Physical Review Letters 94, 057402 (2005).
7 E. A. Stinaff, M. Scheibner, A. S. Bracker, I. V. Ponomarev, V. L. Korenev, M. E. Ware, M. F. Doty, T. L. Reinecke, and D. Gammon, Science 311, 636 (2006).
8 P. Zanardi and M. Rasetti, Phys. Rev. Lett. 79, 3306 (1997).
9 D. A. Lidar, I. L. Chuang, and K. B. Whaley, Phys. Rev. Lett. 81, 2594 (1998).
10 L.-M. Duan and G.-C. Guo, Phys. Rev. Lett. 79, 1953 (1997).
11 J. R. Petta, A. C. Johnson, J. M. Taylor, E. A. Laird, A. Yacoby, M. D. Lukin, C. M. Marcus, M. P. Hanson, and A. C. Gossard, Science 309, 2180 (2005).
12 P. Jarillo-Herrero, S. Sapmaz, C. Dekker, L. P. Kouwenhoven, and H. S. J. van der Zant, Nature 429, 389 (2004).
13 T. B. Bahder, Physical Review B 41, 11992 (1990).
14 Z. Kis and F. Renzoni, Physical Review A 65 (2002).
15 J. I. Cirac, P. Zoller, H. J. Kimble, and H. Mabuchi, Physical Review Letters 78, 3221 (1997).
16 F. Renzoni and S. Stenholm, Optics Communications 189, 69 (2001).
17 K. Hennessy, C. Hogele, E. Hu, A. Badolato, and A. Imamoglu, Applied Physics Letters 89 (2006).
18 A. Badolato, K. Hennessy, M. Atature, J. Dreiser, E. Hu, P. M. Petroff, and A. Imamoglu, Science 308, 1158 (2005).