PHYSICAL REVIEW B **88**, 155321 (2013)

Creation of entangled spin qubits between distant quantum dots

Nobuhiko Yokoshi,[1,*] Hiroshi Imamura,[2] and Hideo Kosaka[3]
[1]*Department of Physics and Electronics, Osaka Prefecture University, Sakai, Osaka 599-8531, Japan*
[2]*Spintronics Research Center, AIST, Tsukuba, Ibaraki 305-8568, Japan*
[3]*Laboratory for Nanoelectronics and Spintronics, RIEC, Tohoku University, Sendai 980-8577, Japan*
(Received 24 May 2013; published 25 October 2013)

We theoretically show that an entangled state can be prepared between distant heavy-hole spin qubits by using III–V semiconductor quantum dots (QDs) in cascaded cavities. In this scheme, using quantum state transfer between photons and carrier spins, the local spin-singlet state in a virtually excited trion is stretched into the two separated QDs via a single-photon transmission. We calculate the effective Liouville equation and determine the optimal conditions for entanglement creation. We also discuss the possibility of distributing entangled electron spins in the same manner.

DOI: 10.1103/PhysRevB.88.155321 PACS number(s): 73.21.La, 03.67.Bg, 03.67.Mn, 71.18.+y

The distribution of entanglements in solid-based qubits is a key step toward building a quantum information network. Such a network consists of several nodes that act as quantum repeaters and transfer informations via quantum teleportation.[1,2] The photon is the strongest contender for an information carrier to connect the distant nodes because it can travel with little decoherence. Thus, the workability of such a network should be determined by how efficiently and reliably we store the quantum state carried by a flying photon and establish quantum entanglement between distant network nodes.

There are many candidates for a memory qubit which stores the quantum state of a photon. Particularly, in order to distribute entangled states to trapped atoms, a lot of photon-mediated schemes have been proposed,[3–9] and several experiments have actually demonstrated entanglement in distant nodes.[10–13] Also, in separated nitrogen-vacancy defects in diamonds, an entanglement state was recently achieved.[14] On the other hand, electron spins confined in quantum dots (QDs) are plausible candidates for solid-state processing qubits. In such qubits, we can electrically conduct various operations and measurements[15–19] and may achieve Bell state measurements.[20,21] Thus, quantum state transfer from photons to electron spins in a semiconductor has opened the possibility of realizing solid-based quantum repeaters.[22–24] Recently this technique is being extended to the case of a single photon and single spin in a QD.[25,26] However, electron spins are inevitably exposed to the fluctuating spins of the host nuclei, which shorten the spin relaxation time T_2^*.[27,28]

The qubit we focus on here is a heavy-hole spin in a QD. In III–V semiconductors such as the GaAs compound, a valence hole has an atomic p-like orbital which vanishes at the location of each nucleus. Actually, the long spin relaxation time[29] and coherent spin operation[30] of a single heavy hole have been observed. The aim of this work is to propose a scheme for creating entangled heavy-hole spins in distant QDs embedded in optical cavities. The essence of our scheme is the quantum state transfer from photons to heavy-hole spins, by which a spin-singlet state in a virtually excited trion is stretched by a photon propagating between the distant QDs.

First, we investigate quantum state transfer between photons and heavy-hole spins. The mechanism of the transfer is intrinsically the same as that for a photon-electron spin interface.[22–24] A photon state on a Poincaré sphere can be spanned by the right and left circular polarization states, i.e., $|\sigma_+\rangle$ and $|\sigma_-\rangle$. In III–V semiconductors, a heavy-hole exciton state is connected by the photon polarization through optical selection rules, i.e., $|\sigma_+\rangle \to -|\frac{3}{2},\frac{3}{2}\rangle_{i\mathrm{h}} \otimes |\frac{1}{2},\frac{1}{2}\rangle_{i\mathrm{e}}$ and $|\sigma_-\rangle \to |\frac{3}{2},-\frac{3}{2}\rangle_{i\mathrm{h}} \otimes |\frac{1}{2},-\frac{1}{2}\rangle_{i\mathrm{e}}$.[31] Here the orbital wave function $|J,m\rangle$ is labeled by the total and magnetic angular momentum $\hbar J$ and $\hbar m$. The subscript $\{i\mathrm{e(h)}\}$ denotes the electron (hole) state on the orbital in the ith QD. Fortunately, the in-plane g factor of heavy holes is approximately 0 ($g_\parallel \sim 0$) in a magnetic field up to 6 T.[32] Then, applying an in-plane magnetic field, we can lift the degeneracy of x-directed electron spins, while the Zeeman shift for heavy-hole spins remains negligible. Thus, it is possible to selectively access one of the electron states $|\pm\rangle_{ie} = (1/\sqrt{2})(|\frac{1}{2},\frac{1}{2}\rangle_{ie} \pm |\frac{1}{2},-\frac{1}{2}\rangle_{ie})$ by a single photon. When one selects the state $|-\rangle_{ie}$ and neglects the terms with $|+\rangle_{ie}$, the following optical transition is achieved:

$$\alpha|\sigma_+\rangle + \beta|\sigma_-\rangle \to \left(\alpha\left|\tfrac{3}{2},-\tfrac{3}{2}\right\rangle_{i\mathrm{h}} + \beta\left|\tfrac{3}{2},\tfrac{3}{2}\right\rangle_{i\mathrm{h}}\right) \otimes |-\rangle_{i\mathrm{e}}. \tag{1}$$

Note that the exciton state in Eq. (1) is the direct product of the electron and heavy-hole states. Thus, the electron spin relaxation time does not set up bounds to the coherence of the spin qubit.[22–24] By extracting the electron to the continuum,[33] the desired quantum state transfer is achieved.[22] Hereafter, we employ the notations for heavy-hole spin $|\frac{3}{2},-\frac{3}{2}\rangle_{i\mathrm{h}} \equiv |\Uparrow\rangle_i$ and $|\frac{3}{2},\frac{3}{2}\rangle_{i\mathrm{h}} \equiv |\Downarrow\rangle_i$.

An overview of the proposed scheme is shown in Fig. 1. We initially prepare a heavy hole in the first QD; i.e., the initial state is some pure or mixed state composed of $|\Uparrow\rangle_1$ and $|\Downarrow\rangle_1$, whereas the second QD is in a vacuum state, $|\mathrm{vac}\rangle_2$. When an off-resonant light is irradiated onto the first QD, it virtually creates an additional pair of an electron with x-directed spin and a heavy hole. On that occasion, the trion states with triplet hole spins lie far above the one with singlet spins because of the strong exchange interaction. Thus, we can consider that only the trion state in $|X^+\rangle_1 = (1/\sqrt{2})(|\Uparrow\rangle_1|\Downarrow\rangle_1 - |\Downarrow\rangle_1|\Uparrow\rangle_1) \otimes |-\rangle_{1\mathrm{e}}$ contributes to the excitation.

Consider a case where the trion radiates a photon, leaving a hole in some state $|\xi\rangle_1 = u|\Uparrow\rangle_1 + v|\Downarrow\rangle_1$ with $|u|^2 + |v|^2 = 1$. According to Eq. (1), the polarization of the radiated photon should be $|\mathrm{ph}\eta\rangle = -v^*|\sigma_+\rangle + u^*|\sigma_-\rangle$. The photon propagates unidirectionally through the waveguide with a polarization-independent isolator. If the second QD successfully absorbs

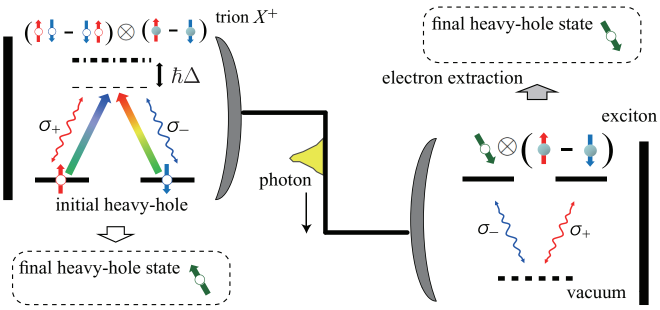

FIG. 1. (Color online) Schematic of the proposed method. In the initial state, we prepare a heavy hole in the first QD. Introducing an off-resonant light to the QD, we virtually excite the trion state, which consists of an electron with x-directed spin and heavy holes forming a spin singlet. The trion radiates a photon and relaxes into a heavy-hole state. The photon propagates unidirectionally and creates the exciton according to the selection rule in the second QD. After extracting the electron, the heavy-hole spins in the QDs should be entangled, reflecting the spin singlet in the trion.

it, the exciton in $|\text{ex}\eta\rangle_2 = -v^*|\text{ex}\uparrow\rangle_2 + u^*|\text{ex}\downarrow\rangle_2$ is created, where $|\text{ex}\uparrow(\downarrow)\rangle_2 = |\uparrow(\downarrow)\rangle_2 \otimes |-\rangle_{2\text{e}}$. On the other hand, when the spin state $|\eta\rangle_1 = -v^*|\uparrow\rangle_1 + u^*|\downarrow\rangle_1$ is left in the first QD, the radiated photon creates an exciton state $|\text{ex}\xi\rangle_2 = u|\text{ex}\uparrow\rangle_2 + v|\text{ex}\downarrow\rangle_2$ in the second QD. The resulting state is

$$\frac{1}{\sqrt{2}}(|\xi\rangle_1|\text{ex}\eta\rangle_2 - |\eta\rangle_1|\text{ex}\xi\rangle_2) = |\sigma\rangle \otimes |-\rangle_{2\text{e}},$$

where $|\sigma\rangle \equiv (1/\sqrt{2})(|\uparrow\rangle_1|\downarrow\rangle_2 - |\downarrow\rangle_1|\uparrow\rangle_2)$. Therefore, extracting the electron into the continuum, the heavy-hole spins in both QDs are entangled.

Henceforth we investigate these processes using a density matrix. In our calculation, we do not take the light holes $|\frac{3}{2}, \pm\frac{1}{2}\rangle_{lh}$ into account. This can be justified when the quantum well containing the QD is sufficiently thin. For a quantum well of $\sim$10-nm thickness, the energy splitting between the heavy hole and the light hole is estimated to be $E_{\text{hh}-\text{lh}} \sim 10$ meV,[34] which is much larger than the light bandwidth. The Liouville equation for the total density matrix ρ_{T} is

$$\frac{d}{dt}\rho_{\text{T}} = -\frac{i}{\hbar}[H,\rho_{\text{T}}] + \mathcal{L}_{\text{spon}}\rho_{\text{T}} + \mathcal{L}_{\text{cav}}\rho_{\text{T}} + \mathcal{L}_{\text{ex}}\rho_{\text{T}}. \quad (2)$$

The Hamiltonian $H = H_{\text{cav}} + H_{\text{qd}} + H_{\text{int}}$ consists of cavity photons, QDs and their interaction. Both the cavities are assumed to have the same frequency: $H_{\text{cav}} = \Sigma_{i=1,2}\hbar\omega_{\text{cav}}(\hat{a}^\dagger_{i+}\hat{a}_{i+} + \hat{a}^\dagger_{i-}\hat{a}_{i-})$. The operator $\hat{a}_{i,p}$ annihilates the cavity photon, with $p = \pm$ denoting the circular polarization. Owing to the zero g factor of the heavy holes and the characteristic selection rule in Eq. (1), the Hamiltonian is described as

$$H_{\text{qd}} = E_c|X^+\rangle_{11}\langle X^+| - E_{\text{ex}}|\text{vac}\rangle_{22}\langle\text{vac}| + \left[\frac{\hbar\Omega}{2}\text{e}^{-i\omega_{\text{L}}t}(|X^+\rangle_{11}\langle\downarrow| - |X^+\rangle_{11}\langle\uparrow|) + \text{H.c.}\right], \quad (3)$$

where $\hbar\Omega$ is the dipole coupling energy to the H-polarized pump light with frequency ω_{L}. The excitation energy of the trion E_c differs from that of the neutral exciton E_{ex} by $\hbar\Delta$, which originates from the extra Coulomb potential by the prepared heavy hole, and is of the order of milli–electron volts.[35] The interaction part is described by

$$H_{\text{int}} = i\hbar g_1(|X^+\rangle_{11}\langle\downarrow|\hat{a}_{1+} - |X^+\rangle_{11}\langle\uparrow|\hat{a}_{1-}) + i\hbar g_2(|\text{ex}\uparrow\rangle_{22}\langle\text{vac}|\hat{a}_{2+} + |\text{ex}\downarrow\rangle_{22}\langle\text{vac}|\hat{a}_{2-}) + \text{H.c.}, \quad (4)$$

with $\hbar g_i$ being the coupling constant.

The term $\mathcal{L}_{\text{spon}}\rho_{\text{T}} = (\mathcal{L}_1 + \mathcal{L}_2)\rho_{\text{T}}$ describes the decays of the excited states $|X^+\rangle_1$ and $|\text{ex}\uparrow(\downarrow)\rangle_2$ with the rates being $\gamma_{1,2}$. The damping of the cavity modes through their mirrors and the unidirectional coupling[36,37] are

$$\mathcal{L}_{\text{cav}}\rho_{\text{T}} = \sum_{p=\pm}\Bigg[\sum_{i=1,2}\frac{\kappa}{2}(2\hat{a}_{i,p}\rho_{\text{T}}\hat{a}^\dagger_{i,p} - \{\hat{a}^\dagger_{i,p}\hat{a}_{i,p},\rho_{\text{T}}\}) - \sqrt{\epsilon\kappa^2}([\hat{a}^\dagger_{2,p},\hat{a}_{1,p}\rho_{\text{T}}] + [\rho_{\text{T}}\hat{a}^\dagger_{1,p},\hat{a}_{2,p}])\Bigg], \quad (5)$$

in which the decay rate κ is set to be common in both cavities. The parameter ϵ $(0 \leqslant \epsilon \leqslant 1)$ models the coupling inefficiency at the second QD and photon losses in the waveguide.[36,37]

The electron continuum, into which one extracts the electron $|-\rangle_{\text{e}}$, is assumed to be positioned just outside the second QD.[38] The extraction is characterized by the spectral density function $\hbar\gamma_{\text{e}}(\varepsilon) = 2\pi\,\Sigma_l|w_l|^2\delta(\varepsilon - \varepsilon_l)$. Here w_l is the tunneling matrix element between the exciton state and the continuum level ε_l. We apply the Markov approximation also to the process and then treat $\gamma_{\text{e}}(\varepsilon)$ as a constant,[38] i.e.,

$$\mathcal{L}_{\text{ex}}\rho_{\text{T}} = \frac{\gamma_{\text{e}}}{2}(2\hat{c}\rho_{\text{T}}\hat{c}^\dagger - \{\hat{c}^\dagger\hat{c},\rho_{\text{T}}\}), \quad (6)$$

where the operator $\hat{c}$ annihilates the electron from the exciton, i.e., $\hat{c}|\text{ex}\uparrow(\downarrow)\rangle_2 = |\uparrow(\downarrow)\rangle_2$.

Because the input light is a continuous one, it can excite the first QD at the same time that the transferred photon excites the second QD. However, we set the pump light sufficiently weak so that it can include only a few photons in the mean. This condition allows us to disregard the states where both QDs are excited, e.g., $|X^+\rangle_1|\text{ex}\uparrow\rangle_2$. Moreover, we consider here the use of a pump light that is in resonance with the cavity modes and the exciton states, i.e., $\hbar\omega_{\text{L}} = \hbar\omega_{\text{cav}} = E_{\text{ex}}$.

In order to simplify the equation further, we treat the above model with some approximations. First, it is assumed that the optical cavities are strongly coupled with the waveguide. Then, the system is in the so-called bad cavity limit ($\kappa \gg \Omega,\gamma_1,\gamma_2,g_1,g_2$). Besides, the frequency detuning in the first QD is typically higher than the pump rate and the decay rates $\kappa \ll \Delta$. These conditions make it possible to apply the adiabatic elimination to the cavity modes.[39] Then the decay rate in $\mathcal{L}_i\bar{\rho}$ is replaced with the cavity-enhanced rate $\Gamma_i = \gamma_i + 4g_i^2/\kappa$. The coupling between the QDs is described by the term

$$\mathcal{L}_{\text{coup}}\bar{\rho} = \frac{\Gamma_{12}}{2}([|\downarrow\rangle_{11}\langle X^+|,|\text{ex}\uparrow\rangle_{22}\langle\text{vac}|\bar{\rho}] - [|\uparrow\rangle_{11}\langle X^+|,|\text{ex}\downarrow\rangle_{22}\langle\text{vac}|\bar{\rho}] + \text{H.c.}), \quad (7)$$

where $\Gamma_{12} = 4g_1g_2\sqrt{\epsilon}/\kappa$ and $\bar{\rho}$ is the reduced density matrix with respect to the cavity modes. Subsequently, we eliminate the trion state $|X^+\rangle_1$ adiabatically and disregard the spontaneous emissions in the first QD. Thus the Rabi frequency in the

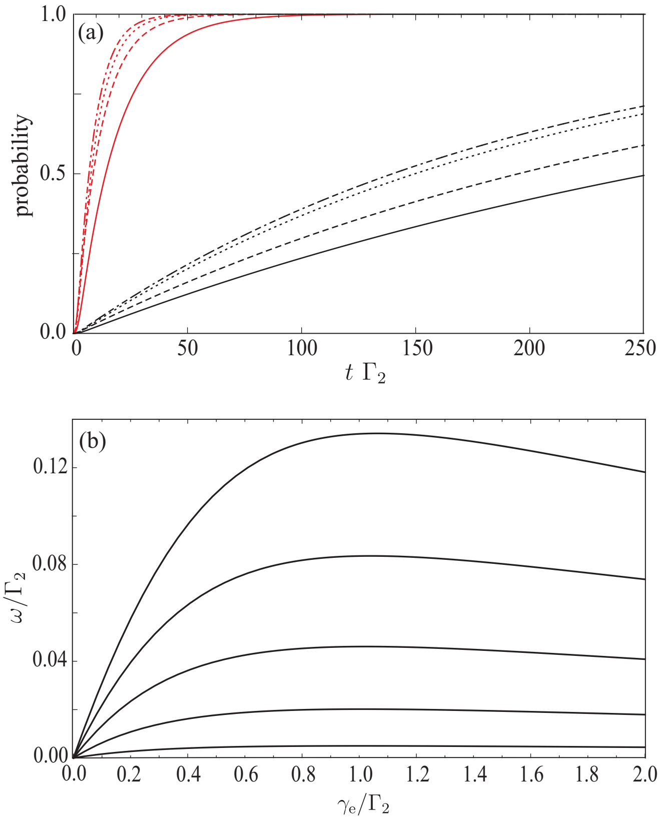

FIG. 2. (Color online) (a) Time dependence of the probability of the entanglement creation is plotted. The lower group of black lines corresponds to the probabilities when the extraction rate γ_e is $0.2\Gamma_2$ (solid line), $0.6\Gamma_2$ (dashed line), $1.0\Gamma_2$ (dashed-dotted line), and $1.4\Gamma_2$ (dotted line). Here we have set the effective inter-QD coupling $\Omega\Gamma_{12}/\sqrt{2}\Delta = 0.1\Gamma_2$. The upper group of (red) lines corresponds to the same plots for $\Omega\Gamma_{12}/\sqrt{2}\Delta = 0.5\Gamma_2$. (b) The inverse of the time constant of entanglement creation is plotted against the extraction rate γ_e. The coupling is set to be $\Omega\Gamma_{12}/\sqrt{2}\Delta = 0.1\Gamma_2$, $0.2\Gamma_2$, $0.3\Gamma_2$, $0.4\Gamma_2$, and $0.5\Gamma_2$ in increasing order of gradient at the origin. One can see that the fastest creation is achieved at $\gamma_e = \Gamma_2$.

first QD becomes $\Omega^2/2\Delta$, and the effective coupling constant is $\Omega\Gamma_{12}/\Delta$[6].

Because only the first QD holds a heavy hole in the initial state, we consider the mixed state of $|\uparrow\rangle_1|\text{vac}\rangle_2$ and $|\downarrow\rangle_1|\text{vac}\rangle_2$ for $\bar{\rho}(0)$. For example, we employ the condition where the probabilities of both states are 1/2 and no coherence exists between them. For the initial state, we calculate the probability of being transferred into the entangled state $|\sigma\rangle$. In Fig. 2(a), we show the time evolution of the probability for different parameters. Here we consider the usage of GaAs-AlGaAs QDs, in which the decay rate is, e.g., $\gamma_2/2\pi = 3.4$ GHz(=14 μeV).[33] One can see that the entanglement is certainly achieved. When the effective inter-QD coupling $\Omega\Gamma_{12}/\sqrt{2}\Delta = 0.5\Gamma_2$, the time it takes is about $30/\Gamma_2 \sim 100/\Gamma_2$. In particular, for $\kappa/2\pi = 40$ GHz (=160 μeV) and $g/2\pi = 20$GHz(=80 μeV), which can be achieved in photonic-crystal cavities,[40] the required time is estimated to be $0.12 \sim 0.4$ ns. Even for smaller coupling ($\Omega\Gamma_{12}/\sqrt{2}\Delta = 0.1\Gamma_2$), the probability approaches 1 within time $\leqslant 500/\Gamma_2 \sim 2.0$ ns (not shown in the figure). Therefore, even when we take into account the time delay of slow light in passing through the waveguide,[41] we can create an entangled pair of heavy-hole spins within the dephasing time $T_2^* \simeq 1$ ms.[29]

Moreover, in increasing the extraction rate on the second QD γ_e, the alteration in the time evolution is not monotonic. This can be confirmed by the analytical expression of the solution of the approximated Liouville equation. Here we investigate the inverse of the time constant for the entanglement creation, which is

$$\omega = g - R\cos\left(\tfrac{1}{3}\arccos\Lambda\right), \tag{8}$$

where $g = (\Gamma_2 + \gamma_e)/2$ and

$$R = \sqrt{\frac{2}{3}}\sqrt{2g^2 - \left(\frac{\Omega\Gamma_{12}}{\Delta}\right)^2}, \tag{9}$$

and

$$\Lambda = \frac{(\Gamma_2 - \gamma_e)}{R^3}\left(\frac{\Omega\Gamma_{12}}{\Delta}\right)^2. \tag{10}$$

We take a plot of Eq. (8) against the extraction rate [Fig. 2(b)]. The required time for the whole process reaches a minimum at $\gamma_e = \Gamma_2$. This is because a large γ_e not only effectively leaves a heavy hole in the second QD, but also prevents the creation of the exciton there. However, the influence of the deviation from the optimal point $\gamma_e = \Gamma_2$ is appreciably small for the large-γ_e regime. Therefore, in verifying the scheme in the experiment, it may be desirable to target the regime $\gamma_e \geqslant \Gamma_2$.

For III–V (or II–VI) semiconductor QDs in a cavity, a controlled-NOT operation between electron spins has been proposed, which can create entanglement.[42] This protocol is based on temporal switching of effective long-range spin-spin coupling between the QDs mediated by the cavity modes. Our proposal directly utilizes the local spin-singlet state of the pure trion state in one QD, and such operations are not necessary. Besides, compared with previous works,[3–7,42] neither preparing the initial state in some pure state nor utilizing more than one laser fields is required. As long as the pump laser is radiated to the first QD, the event in Fig. 1 is repeated without any initializations, which is the most important advantage of our scheme over other similar schemes already put forward. Furthermore, in our proposal, the entanglement creation is accomplished only when the extracted electron is detected. Thus, the fidelity at $t \to \infty$ does not depend on the efficiency of the cavity-photon coupling ϵ. One may be afraid that phonons in the QD disturb the trion state and inhibit the entanglement creation. However, the laser field is detuned from $|X^+\rangle_1$ by $\hbar\Delta \sim 1$ meV and is excited only virtually. In such a situation, the effects of the phonons are negligible at low temperature.[43]

In summary, we have shown that an entangled state can be prepared between distant heavy holes by using characteristic optical transitions in III–V semiconductor QDs. Our proposal utilizes the spin-singlet state in the trion due to local spin-spin interaction: we convert the singlet state in one QD into an entangled spin state in distant QDs through a spontaneously radiated photon. We determine the optimal conditions for smooth entanglement creation. Related technologies have grown dramatically in the past few years.[23–26,29,30] Thus, although the proposed setup is still a challenging one, we believe the proposed scheme can be achieved by combining these technologies.

The proposed scheme can be applied as well to the creation of entangled electron spins. Quantum state transfer from photons to electron spins can be achieved by preparing semiconductor QDs in which the electron g factor is controlled to be near zero.[23–25] Then, in the above calculations, we just have to redefine $|\uparrow(\downarrow)\rangle_i$ by $|\frac{1}{2},\pm\frac{1}{2}\rangle_{ie}$ and replace x-directed electron spin $|-\rangle_{ie}$ with light-hole spin $|-\rangle_{ih} = (1/\sqrt{2})(|\frac{3}{2},\frac{1}{2}\rangle_{ih} - |\frac{3}{2},-\frac{1}{2}\rangle_{ih})$[22]. In practice, in order to experimentally verify the scheme for the electron spin case, the problem of a short dephasing time T_2^* is necessarily encountered. Then, in addition to g-factor engineering, it requires manipulations, e.g., dynamic nuclear polarization.[44,45] However, if we become able to distribute an entanglement efficiently into processor electron spin qubits, this should take us a step closer to the establishment of solid-based quantum repeaters and quantum information networks.

The authors would like to thank T. Kutsuwa, T. Inagaki, and H. Ishihara for fruitful discussions. This work was partly supported by CREST of the Japan Science and Technology Agency. N.Y. was supported by Grant-in-Aid for Young Scientists (B) No. 24760046 from MEXT, Japan. H.K. was supported by Grants-in-Aid for Scientific Research Nos. 24244044 and 23104701 from MEXT and the Funding Program for World-Leading Innovative R&D on Science and Technology.

---

*Corresponding author: yokoshi@pe.osakafu-u.ac.jp

[1] C. H. Bennett, G. Brassard, C. Crepeau, R. Jozsa, A. Peres, and W. K. Wootters, Phys. Rev. Lett. **70**, 1895 (1993).
[2] H.-J. Briegel, W. Dür, J. I. Cirac, and P. Zoller, Phys. Rev. Lett. **81**, 5932 (1998).
[3] J. I. Cirac, P. Zoller, H. J. Kimble, and H. Mabuchi, Phys. Rev. Lett. **78**, 3221 (1997).
[4] A. S. Parkins and H. J. Kimble, J. Opt. B: Quantum Semiclass. Opt. **1**, 496 (1999).
[5] S. Bose, P. L. Knight, M. B. Plenio, and V. Vedral, Phys. Rev. Lett. **83**, 5158 (1999).
[6] S. Clark, A. Peng, M. Gu, and S. Parkins, Phys. Rev. Lett. **91**, 177901 (2003).
[7] D. Pinotsi and A. Imamoglu, Phys. Rev. Lett. **100**, 093603 (2008).
[8] X.-L. Feng, Z.-M. Zhang, X.-D. Li, S.-Q. Gong, and Z.-Z. Xu, Phys. Rev. Lett. **90**, 217902 (2003).
[9] L.-M. Duan, M. D. Lukin, J. I. Cirac, and P. Zoller, Nature **414**, 413 (2001).
[10] D. L. Moehring, P. Maunz, S. Olmschenk, K. C. Younge, D. N. Matsukevich, L.-M. Duan, and C. Monroe, Nature **449**, 68 (2007).
[11] D. N. Matsukevich, P. Maunz, D. L. Moehring, S. Olmschenk, and C. Monroe, Phys. Rev. Lett. **100**, 150404 (2008).
[12] P. Maunz, S. Olmschenk, D. Hayes, D. N. Matsukevich, L.-M. Duan, and C. Monroe, Phys. Rev. Lett. **102**, 250502 (2009).
[13] S. Olmschenk, D. N. Matsukevich, P. Maunz, D. Hayes, L.-M. Duan, and C. Monroe, Science **323**, 486 (2009).
[14] H. Bernien, B. Hensen, W. Pfaff, G. Koolstra, M. S. Blok, L. Robledo, T. H. Taminiau, M. Markham, D. J. Twitchen, L. Childress, and R. Hanson, Nature **497**, 86 (2013).
[15] J. 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).
[16] K. C. Nowack, F. H. L. Koppens, Yu. V. Nazarov, and L. M. K. Vandersypen, Science **318**, 1430 (2007).
[17] J. M. Taylor, H.-A. Engel, W. Dür, A. Yacoby, C. M. Marcus, P. Zoller, and M. D. Lukin, Nat. Phys. **1**, 177 (2005).
[18] R. Hanson and G. Burkard, Phys. Rev. Lett. **98**, 050502 (2007).
[19] N. Yokoshi, H. Imamura, and H. Kosaka, Phys. Rev. Lett. **103**, 046806 (2009).
[20] H.-A. Engel and D. Loss, Science **309**, 586 (2005).
[21] N. Yokoshi, H. Imamura, and H. Kosaka, Phys. Rev. B **81**, 161305(R) (2010).
[22] R. Vrijen and E. Yablonovitch, Physica E **10**, 569 (2001).
[23] H. Kosaka, H. Shigyou, Y. Mitsumori, Y. Rikitake, H. Imamura, T. Kutsuwa, K. Arai, and K. Edamatsu, Phys. Rev. Lett. **100**, 096602 (2008).
[24] H. Kosaka, T. Inagaki, Y. Rikitake, H. Imamura, Y. Mitsumori, and K. Edamatsu, Nature **457**, 702 (2009).
[25] M. Kuwahara, T. Kutsuwa, K. Ono, and H. Kosaka, Appl. Phys. Lett. **96**, 163107 (2010).
[26] A. Pioda, E. Totoki, H. Kiyama, T. Fujita, G. Allison, T. Asayama, A. Oiwa, and S. Tarucha, Phys. Rev. Lett. **106**, 146804 (2011).
[27] W. A. Coish and D. Loss, Phys. Rev. B **72**, 125337 (2005).
[28] F. H. L. Koppens, C. Buizert, I. T. Vink, K. C. Nowack, T. Meunier, L. P. Kouwenhoven, and L. M. K. Vandersypen, J. Appl. Phys. **101**, 081706 (2007).
[29] B. D. Gerardot, D. Brunner, P. A. Dalgarno, P. Öhberg1, S. Seidl, M. Kroner, K. Karrai, N. G. Stoltz, P. M. Petroff, and R. J. Warburton, Nature **451**, 441 (2008).
[30] K. D. Greve, P. L. McMahon, D. Press, T. D. Ladd, D. Bisping, C. Schneider, M. Kamp, L. Worschech, S. Höfling, A. Forchel, and Y. Yamamoto, Nat. Phys. **7**, 827 (2011).
[31] F. Meier and B. P. Zakharchenya (eds.), *Optical Orientation* (Elsevier, Amsterdam, 1984).
[32] J. G. Tischler, A. S. Bracker, D. Gammon, and D. Park, Phys. Rev. B **66**, 081310(R) (2002).
[33] B. Zhang, D. W. Snoke, and A. P. Heberle, Solid State Commun. **152**, 296 (2012).
[34] P. Brick, C. Ell, G. Khitrova, H. M. Gibbs, T. Meier, C. Sieh, and S. W. Koch, Phys. Rev. B **64**, 075323 (2001).
[35] A. S. Bracker, E. A. Stinaff, D. Gammon, M. E. Ware, J. G. Tischler, A. Shabaev, Al. L. Efros, D. Park, D. Gershoni, V. L. Korenev, and I. A. Merkulov, Phys. Rev. Lett. **94**, 047402 (2005).
[36] C. W. Gardiner, Phys. Rev. Lett. **70**, 2269 (1993).
[37] H. J. Carmichael, Phys. Rev. Lett. **70**, 2273 (1993).
[38] Y. Rikitake, H. Imamura, and H. Kosaka, J. Phys. Soc. Jpn. **76**, 114004 (2007).

[39] J. I. Cirac, Phys. Rev. A **46**, 4354 (1992).
[40] A. Faraon, A. Majumdar, H. Kim, P. Petroff, and J. Vučković, Phys. Rev. Lett. **104**, 047402 (2010).
[41] T. Baba, Nat. Photon. **2**, 465 (2008).
[42] A. Imamoḡlu, D. D. Awschalom, G. Burkard, D. P. DiVincenzo, D. Loss. M. Sherwin, and A. Small, Phys. Rev. Lett. **83**, 4204 (1999).
[43] A. Grodecka, C. Weber, P. Machnikowski, and A. Knorr, Phys. Rev. B **76**, 205305 (2007).
[44] D. J. Reilly, J. M. Taylor, J. P. Petta, C. M. Marcus, M. P. Hanson, and A. C. Gossard, Science **321**, 817 (2008).
[45] J. R. Petta, J. M. Taylor, A. C. Johnson, A. Yacoby, M. D. Lukin, C. M. Marcus, M. P. Hanson, and A. C. Gossard, Phys. Rev. Lett. **100**, 067601 (2008).