Intrinsic Quantum Noise in Faraday Rotation Measurements of a Single Electron Spin

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(Dated: 2 September 2008)

Faraday rotation is one way to realize quantum non-demolition measurement of electron spin in quantum dots. To describe Faraday rotation, semiclassical models are typically used, based on quantized electron spin states and classical electromagnetic fields. Such treatments neglect the entanglement between electronic and photonic degrees of freedom that produce intrinsic quantum noise, limiting the ultimate sensitivity of this technique. We present a fully quantum-mechanical description of Faraday rotation, and quantify this intrinsic noise. A method for measuring the purity of a given spin state is suggested based on this analysis.

INTRODUCTION

Because of the discovery of long-lived spin coherence in semiconductors such as GaAs[1], the essential requirement of manipulating spins for spintronics and quantum information is now possible. The first quantum computing proposal of Loss and DiVincenzo used electron spin qubits in semiconductor quantum dots (QD)[2] and forecast the importance of measuring single electrons and their spins.

The first step to realizing coherent manipulation of a single electron spin is to orient the spin. Such orientation can be achieved optically (by exciting with circularly polarized light)[3], electrically (by driving the electrons toward a ferromagnetic surface)[4] or thermodynamically (by application of a uniform magnetic field at low temperatures). Photoluminescence (PL) allows for measurement of electron spin polarization through the relation between the circular polarization of light and electron spin orientation. However, PL is destructive in that it involves recombination of the electron with a hole. PL measurements are intrinsically limited by the lifetime of the state, and it is not possible to monitor electron spin continuously. Furthermore, unless one uses a technique such as time-resolved upconversion[5, 6, 7] or streak camera measurements, dynamical information is lost.

Time-resolved Faraday and Kerr rotation methods (hereafter referred to as Faraday rotation) have been extensively developed[8, 9, 10], and allow one to probe the spin dynamics of a single electron in a quantum dot. Faraday rotation results from a fundamental interaction between electronic and photonic degrees of freedom. Seigneur et al.[11] have proposed a scheme to implement quantum computation by using the single photon Faraday effect. However, in most semiconductors the Faraday effect is usually quite weak, corresponding to rotation angles \( \theta_F \sim 10^{-3} \text{rad} \) for single electrons. Dynamic information is usually obtained using pump-probe optical techniques: a circularly polarized pump beam creates an initially spin-polarized electron population, and a probe beam subsequently interrogates the spin state at a later time. The experiment is performed repeatedly as a function of the delay to obtain a time-resolved signal with an acceptably high signal-to-noise ratio. In the case of a single electron in a quantum dot, spin coherence can be achieved in the following manner: the quantum dot is configured (either through biasing or doping) to begin in a state that contains a single electron in the conduction band and no holes in the valence band. The quantum dot is excited, promoting a second electron into the conduction band and leaving behind a hole in the valence band. This state is often referred to as a “trion”. After one of the electrons recombines with the hole, the remaining electron spin is partially polarized. A linearly polarized probe pulse measures the spin of this electron via the Faraday effect. In most cases, the electron neither begins in a pure state nor remains in one. Hyperfine interactions with nuclear spins quickly produce a mixed state on time scales \( \sim 1-10 \text{ns} \)[12, 13, 14, 15, 16, 17, 18]. In this paper, we study the noise introduced by the mixed quantum state of the electron spin analytically and numerically. In this paper, our previous analysis[19] about the noise is extended to a more formal quantum mechanical frame. Since it’s from the spin state itself, we call it intrinsic noise.

![FIG. 1: light induced interband transition](image)
THEORETICAL MODEL

Here we discuss in detail the quantum-mechanical source of this noise using a theoretical model that treats both the electron and light field quantum mechanically. We model the interaction between a single electron in a QD and a linearly polarized monochromatic probe laser field. The Hamiltonian for the photon field can be written as

\[ H_P = \hbar \omega_P (a_L^\dagger a_L + a_R^\dagger a_R), \]

where \( \omega_P \) is the optical frequency of the probe laser, \( a_L^\dagger \) and \( a_L \) are creation and annihilation operators for left circularly polarized (LCP) photons; \( a_R^\dagger \) and \( a_R \) are creation and annihilation operators for right circularly polarized (RCP) photons. Due to optical selection rules [3] spin-up (spin-down) electrons interact only with LCP (RCP) photons. (See FIG. [1]). The raising and lowering operators satisfy boson commutation relations

\[ [a_m, a_n^\dagger] = \delta_{mn}, (m, n = L, R); \]
\[ [a_m, a_n] = 0, [a_m^\dagger, a_n^\dagger] = 0. \]

The electron state is quantized as well. We assume that the electron resides in the conduction band quantum-confined ground state in an s orbital, which means it has total angular momentum \( J = \frac{3}{2} \). In the valence band, the electronic ground states are constructed from p-orbitals, and hence the total angular momentum is \( J = \frac{5}{2} \). The Hamiltonian for the electron is given by [20]

\[ H_e = \hbar \omega_e (\sigma_{uz} + \sigma_{dz}), \]

where

\[ \sigma_{uz} = b^\dagger_{cu} b_{cu} - b^\dagger_{vu} b_{vu}, \]
\[ \sigma_{dz} = b^\dagger_{cd} b_{cd} - b^\dagger_{vd} b_{vd}; \]

subscript "c" and "v" indicate conduction band and valence band respectively; subscripts "u" and "d" refer to spin-up or spin-down states of the electron. The fermion operators satisfy anticommutation relations:

\[ \{b_{ij}, b^\dagger_{ji}\} = \delta_{ij}\delta_{\mu\nu}, \]
\[ \{b_{ij}, b_{kj}\} = 0, \{b^\dagger_{ij}, b^\dagger_{kj}\} = 0, \]

where \( i \) and \( j \) indicate conduction band or valence band, and \( \mu \) and \( \nu \) indicate spin-up or spin-down. Heavy-hole and light-hole intermixing is neglected for simplicity and because it is not expected to affect qualitatively our results. Only the heavy-hole subband is accounted for in our calculation. A LCP photon couples to a transition between \( \left| + \frac{3}{2} \right> \) and \( \left| + \frac{5}{2} \right> \), while a RCP photons couples to a transition between \( \left| - \frac{1}{2} \right> \) and \( \left| - \frac{3}{2} \right> \). The interaction Hamiltonian is given by

\[ H_I = \lambda_{Lu} (a_L^\dagger \sigma_{uz} + a_L^\dagger \sigma_{dz}) + \lambda_{Rd} (a_R^\dagger \sigma_{uz} + a_R^\dagger \sigma_{dz}), \]

where

\[ \lambda_{Lu} \propto <\frac{1}{2}|x + iy| - \frac{3}{2}>, \]
\[ \lambda_{Rd} \propto <\frac{1}{2}|x - iy| - \frac{3}{2}>, \]
\[ \sigma_{uz} = b^\dagger_{cu} b_{cu}, \sigma_{uz} = \sigma_{uz}, \]
\[ \sigma_{dz} = b^\dagger_{cd} b_{cd}, \sigma_{dz} = \sigma_{dz}. \]

The full Hamiltonian of the entire system is given by

\[ H = H_P + H_e + H_I \]

By applying the Wigner-Eckart theorem, it can be shown that the two coupling strengths \( \lambda_{Lu} \) and \( \lambda_{Rd} \) must be equal \( (\lambda_{Lu}=\lambda_{Rd}=\lambda) \). Based on the defining anticommutation relations, it can be explicitly shown that \( \sigma_{uz}, \sigma_{uz} \) and \( \sigma_{uz} \) have the following commutation relations

\[ [\sigma_{uz}, \sigma_{uz}] = \delta_{\mu\nu}\sigma_{uz}, \]
\[ [\sigma_{uz}, \sigma_{uz}] = 2\delta_{\mu\nu}\sigma_{uz}, \]
\[ [\sigma_{uz}, \sigma_{uz}] = 2\delta_{\mu\nu}\sigma_{uz}. \]

These commutation relations for \( \sigma_{uz}, \sigma_{uz} \) and \( \sigma_{uz} \) are formally identical to those for the Pauli operators, even though they are actually products of fermionic creation and annihilation operators. This feature makes it possible to find an analytic solution within the Heisenberg picture [21]. In the limit where the coupling strength is much smaller than the incident photon frequency or the characteristic frequency of the electron, the approximate solution for photon operators is as follows:

\[ a_L^\dagger (t) = e^{-i\omega t}a_L^\dagger + g(t)(\sigma_{uz} + \alpha \sigma_{uz} a_L^\dagger), \]

where

\[ \alpha = \frac{\lambda}{\omega_L - \omega_e}, \]

\[ a_R^\dagger (t) = e^{-i\omega t}a_R^\dagger + g(t)(\sigma_{uz} + \alpha \sigma_{uz} a_R^\dagger), \]

and

\[ g(t) = \frac{1}{\hbar} \int_{-\infty}^{t} e^{i\omega \tau} H_I(\tau) d\tau. \]
FIG. 2: The left figure shows the polarization ellipse in real space. The right figure is the Stokes representation of the same polarization.

\[ \Omega = \lambda \alpha, \]

\[ g(t) = \alpha(1 - e^{-i(\omega_p - \omega_e) t}). \]

This approximate solution is correct only when the coupling strength \( \lambda \) is much smaller than \( \omega_e \) and \( \omega_P \). In the section of RESULTS, one can see this criteria is satisfied in the sense that the coupling strength of our sample is in the order of \( \sim 10^9 \) Hz, but the frequency of laser and characteristic frequency of electron is in the order of \( \sim 10^{15} \) Hz. This solution, therefore, is a very good approximation and based on this, one can derive Faraday rotation angle.

**Faraday rotation operator**

Quantum Stokes operators can be used to describe Faraday rotation. They are the quantum-mechanical analogue of classical Stokes parameters. Classical Stokes parameters are defined as the following [22]

\[
\begin{align*}
S_0 &= E_x^* E_x + E_y^* E_y \\
S_1 &= E_x^* E_y - E_y^* E_x \\
S_2 &= i(E_x^* E_y + E_y^* E_x) \\
S_3 &= E_x^* E_x - E_y^* E_y
\end{align*}
\]

In electrodynamics, the polarization of light can be parameterized by two angles \( \varphi \) and \( \chi \) in the polarization ellipse. There is a one-to-one correspondence between the polarization-ellipse representation and the Stokes representation (See FIG. 2).

Once the light field is known, the Stokes parameters can be computed. The physical interpretation of \( S_0 \) is the light intensity; hence, all parameters can be normalized to \( S_0 \) (See FIG. 3).

Quantum Stokes operators are defined in the following way [23, 24]

\[
\begin{align*}
S_0 &= a_L^\dagger a_L + a_R^\dagger a_R \\
S_1 &= a_L^\dagger a_R + a_R^\dagger a_L \\
S_2 &= i(a_L^\dagger a_R - a_R^\dagger a_L) \\
S_3 &= a_L^\dagger a_R - a_R^\dagger a_L
\end{align*}
\]

Information about polarization is obtained by calculating the expectation values of these operators. In a typical Faraday experiment, the probe light is linearly polarized at a 45° angle with respect to a final polarizing beamsplitter. After the interaction between the probe light and the electron, the polarization of the transmitted light will be rotated from its initial position by an angle \( \theta_F \), known as the Faraday rotation angle. In the Stokes representation, the initial polarization vector lies along the positive \( S_2 \) axis. Faraday rotation will result in a rotation of the vector within the \( S_1 - S_2 \) plane (See FIG. 4). This vector \( P \) is confined to the plane as long as there is no circular dichroism that can lead to a non-zero expectation value for \( S_3 \).

In our calculations, we aim to reproduce the overall magnitude of the rotation angle that has been reported in experimental work [8, 9, 10]. The experimentally observed rotation angle is small: \( \theta_F \sim 10^{-5} \) rad. Hence, it can be expressed as

\[
\theta_F = \frac{1}{2} \tan^{-1} \left( \frac{<S_1>}{<S_2>} \right) \approx \frac{<S_1>}{2 <S_2>}. \quad (6)
\]
In order to start with following condition must be satisfied

\[ |\text{light field}, \quad \text{where } |\nu_L^2| \text{ and } |\nu_R^2| \text{ are the average number of left and right circularly polarized photons. These states satisfy the canonical eigenvalue equations for the (non-Hermitian) photon annihilation operators:} \]

\[ a_L|\nu_L, \nu_R > = \nu_L|\nu_L, \nu_R >, \]

\[ a_R|\nu_L, \nu_R > = \nu_R|\nu_L, \nu_R >, \]

Using the form \( \nu_L = N_L e^{i\theta_L} \) and \( \nu_R = N_R e^{i\theta_R} \), the expectation value of Stokes operators in this coherent state can be found

\[
\begin{bmatrix}
< S_0 > \\
< S_1 > \\
< S_2 > \\
< S_3 > \\
\end{bmatrix} = \begin{bmatrix}
N_L^2 + N_R^2 \\
2N_L N_R \cos(\theta_L - \theta_R) \\
2N_L N_R \sin(\theta_L - \theta_R) \\
N_R^2 - N_L^2 \\
\end{bmatrix}.
\]  

(7)

In order to start with +45° linearly polarized light, the following condition must be satisfied

\[
\begin{cases}
N_L^2 = N_R^2 \\
\theta_L - \theta_R = \pi / 2
\end{cases}
\]

To describe the mixed state of electron, a density matrix formula is employed.

\[
\rho_e = \tau |+\frac{1}{2}> + (1 - \tau) |1> - \frac{1}{2} |\frac{1}{2}>
\]

(8)

Here, \( \tau \) is a parameter that varies between 0 and 1. For \( \tau = 0 \) and \( \tau = 1 \), one has a pure state, while \( \tau = 1/2 \) corresponds to a fully mixed (unpolarized) state. Because the electron Hamiltonian is expressed in terms of creation and annihilation operators, caution must be taken when applying those operators onto electron state. When operators for spin-up electron are applied to the spin-up state, one obtains

\[
\sigma_{u+}| + \frac{1}{2} > = | + \frac{1}{2} >, \]

\[
\sigma_{u+}| + \frac{3}{2} > = | + \frac{3}{2} >, \]

and

\[
\sigma_{u-}| + \frac{1}{2} > = 0, \]

\[
\sigma_{u-}| + \frac{3}{2} > = | + \frac{3}{2} >, \]

\[
\sigma_{u-}| + \frac{3}{2} > = 0.
\]

Spin-down operators have the same rules when applied to the spin-down state. If a spin-up operator operates on a spin-down state, however, one gets zero. For example,

\[
\sigma_{d+}| - \frac{1}{2} > = (b^\dagger \nu_u b \nu_u - b^\dagger \nu_u b \nu_u)| - \frac{1}{2} > = 0.
\]

The initial state of the whole system is then

\[
\rho_0(\tau) = |\nu_L, \nu_R > < \nu_L, \nu_R| \otimes (\tau |\uparrow> + (1 - \tau) |\downarrow>)
\]

The solution of the whole system is then

\[
\rho_0(\tau) = |\nu_L, \nu_R > < \nu_L, \nu_R| \otimes (\tau |\uparrow> + (1 - \tau) |\downarrow>)
\]

The initial state of the whole system is then

\[
\rho_0(\tau) = |\nu_L, \nu_R > < \nu_L, \nu_R| \otimes (\tau |\uparrow> + (1 - \tau) |\downarrow>)
\]

According to the solution (4) and (5), the analytical expression for \( S_1 \) and \( S_2 \) can be obtained and the expectation values calculated

\[
< S_1 > = Tr(S_1(t)\rho_0(\tau)), \]

\[
< S_2 > = Tr(S_2(t)\rho_0(\tau)).
\]

The rotation angle is given by

\[
\theta_F(t, \tau) = \frac{Tr(S_1(t)\rho_0(\tau))}{2Tr(S_2(t)\rho_0(\tau))},
\]

(10)

After some algebra, one finds the following expression for the Faraday rotation:

\[
\theta_F(t, \tau) = (2\tau - 1)\left(\frac{\lambda^2}{\delta^2} \sin(\delta t) - \sin\left(\frac{\lambda^2}{\delta} t\right)\right),
\]

(11)

where \( \delta \equiv \omega_p - \omega_e \). For initial pure spin-up state \( \tau = 1 \), the rotation angle is

\[
\theta_+ \equiv \theta_F(t, 1) = \left(\frac{\lambda^2}{\delta^2} \sin(\delta t) - \sin\left(\frac{\lambda^2}{\delta} t\right)\right).
\]

(12)
For initial pure spin-down state $\tau = 0$, the rotation angle is
\[ \theta_- = \theta_F(t, 0) = \left( \frac{\lambda^2}{\delta^2} \sin(\delta t) - \sin \left( \frac{\lambda^2}{\delta} t \right) \right). \quad (13) \]

The fluctuation is given by
\[ \Delta \theta_F(t, \tau) = \frac{\sqrt{Tr(S_1^2(t)\rho_0(\tau)) - Tr(S_1(t)\rho_0(\tau))^2}}{2Tr(S_2(t)\rho_0(\tau))}. \quad (14) \]

From (11), (12) and (13), the following intuitive result can be proven very easily:
\[ \theta_F(t, \tau) = \tau \theta_+ + (1 - \tau)\theta_- , \quad (15) \]

where $\theta_+$ ($\theta_-$) is the Faraday rotation angle for an initial state which is a pure spin-up (spin-down) state.

An analytical derivation shows that the fluctuation is a function of both photon number and the initial electron state.
\[ \Delta \theta_F(t, \tau) = \sqrt{\frac{1}{4N} + \tau(1 - \tau)(\theta_+ - \theta_-)^2}, \quad (16) \]

The second term under the square root is the so-called intrinsic noise term.

Numerical simulation is done so that we compare our analytical calculation to recent Kerr rotation experimental results on single electrons. From Berezosky et al\[3\], one finds from a PL plot that the energy for a neutral exciton is about 1.633meV. That corresponds to the band gap between the top of the valence band and the bottom of the conduction band. From this number, the frequency $\omega_e = \frac{E}{\hbar} = 2.48 \times 10^{15} Hz$. Choosing probe light of wavelength 760nm, which means the frequency is $\omega_P = 2.47 \times 10^{15} Hz$. From the paper\[20\], one can take the value of the coupling strength to be $\lambda = 98$GHz. In our experiment, the probe power is about 1.57$\mu$W, the corresponding photon number is about $5 \times 10^5$. In the simulation, the interaction time between the spin and the photon is set to be 20ps. Notice that as expected, if the initial electron state is a pure state, the rotation angle has opposite values for the spin-up state and spin-down state, respectively (See FIG. 5). For pure spin-up states or spin-down states, the fluctuation (quantum noise) scales with photon number $N$ as $N^{-1/2}$, as expected for shot noise. For mixed states or superposition states, the fluctuation saturates even when the photon number approaches infinity (See FIG. 6).

One scheme to measure $\tau$ is proposed here. Suppose the photon number is so large that the shot noise term in Equation (16) could be neglected. Notice that $\theta_+ = \theta_-$ and when the rotation angle is zero, according to Equation (11) it means $\tau = \frac{1}{2}$. If the value $\tau = \frac{1}{2}$ is used in Equation (16), one obtains $\Delta \theta_{F0} = \theta_+$. This result implies that one can use the measured values of Faraday angle fluctuation at an extreme value ($\Delta \theta_F$) and at a zero crossing ($\Delta \theta_{F0}$) to measure the purity of the spin state as quantified by $\tau$:
\[ \tau = \frac{1}{2} \left( 1 + \sqrt{\frac{\Delta \theta_{F0}^2 - \Delta \theta_F^2}{\Delta \theta_{F0}^2}} \right). \quad (17) \]

In the limit of a large number of photons (i.e., where shot noise can be neglected), the above expression simplifies further
\[ \tau = \frac{1}{2} \left( 1 + \sqrt{1 - \frac{\Delta \theta_F^2}{\Delta \theta_{F0}^2}} \right). \quad (18) \]

The above analysis is based on the assumption that every device in the experiment is perfect, and the noise is only introduced by quantum state of the electron spin.
FIG. 6: (a) Fluctuation of Faraday rotation angle as a function of time and parameter $\tau$. (b) Shot noise and intrinsic noise as a function of photon number $N$. Shot noise (black) is from a pure spin-up (spin-down) state, while intrinsic noise (green) is from a maximally mixed state. In the simulation, the interaction time is chosen to be 20ps in order to make the splitting more obvious, in which case the intrinsic noise saturates at about $\Delta \theta_{F0} = 20\text{mrad}$.

fluctuation due to quantum states

$$\Delta \theta^2_F = \Delta \theta^2_M - \Delta \theta^2_B,$$

and

$$\Delta \theta^2_{F0} = \Delta \theta^2_{M0} - \Delta \theta^2_B.$$

The above equation therefore becomes

$$\tau = \frac{1}{2} \left( 1 \pm \sqrt{1 - \frac{\Delta \theta^2_{M} - \Delta \theta^2_B}{\Delta \theta^2_{M0} - \Delta \theta^2_B}} \right).$$ (19)

Notice that in the above equation, $\Delta \theta_M$ is the "real" noise we see at the extreme points of the rotation angle in the actual experiment. This noise has two sources: external noise, which is $\Delta \theta_B$ and intrinsic noise, which is $\Delta \theta_F$. In the actual experiment, $\Delta \theta_M$ should be larger than $\Delta \theta_B$ due to the fact that the pump for the electron spin is not perfect, therefore the spin that interacts with photon is in a mixed state. However, if $\Delta \theta_M = \Delta \theta_B$, that means the intrinsic noise contribution is zero. From Eq. (16), one can see in the limit of large photon number, $\Delta \theta_F = 0$ indicates that $\tau$ is either 0 or 1, which is consistent with the result if one plugs $\Delta \theta_M = \Delta \theta_B$ into Eq. (19). In other words, if in the real experiment, one observes $\Delta \theta_M = \Delta \theta_B$, then the pumped spin is in either pure spin-up or spin-down state and one can also pin down the orientation of spin by looking at the sign of measured rotation angle.

CONCLUSION

Using a quantum-mechanical model of Faraday rotation, we find that both the Faraday rotation angle and the fluctuation are functions of the initial electron spin state. If the electron spin is initially in a mixed state, intrinsic noise fluctuations will contain not only shot noise but also intrinsic noise due to weak measurement of the electron's spin state. The reason that this intrinsic noise appears in this scheme is that the measurement done here is non-destructive, and differs from a projective measurement, which causes the collapse of the electron spin wave function to a certain spin direction. Analysis of the noise spectrum should enable quantification of the purity of a given spin state.

[1] J. M. Kikkawa and D. D. Awschalom, Phys. Rev. Lett. 80, 4313 (1998).
[2] D. Loss and D. P. DiVincenzo, Phys. Rev. A 57, 120 (1998).
[3] F. Meier and B. P. Zakharchenya, Optical Orientation (Elsevier Science Publishers B. V., Amsterdam, 1984).
[4] R. J. Epstein, I. Malajovich, R. K. Kawakami, Y. Chye, M. Hanson, P. M. Petroff, A. C. Gossard, and D. D. Awschalom, Phys. Rev. B, 65, 121202 (2002).
[5] J. F. Smyth, D. A. Tulchinsky, D. D. Awschalom, N. Samarth, H. Luo, and J. K. Furdyna, Phys. Rev. Lett. 71, 601 (1993).
[6] M. P. Hehlen, G. Frei, and H. U. Güdel, Phys. Rev. B 50, 16264 (1994).
[7] R. Kumar, A. S. Vengurlekar, S. S. Prabh, J. Shah, and L.N.Pfeiffer, Phys. Rev. B 54, 4891 (1996).
[8] J. Berezovsky, M. H. Mikkelsen, O. Gywat, N. G. Stoltz, L. A. Coldren, and D.D.Awshalom, Science 314, 1916 (2006).
[9] M. Atature, J. Dreiser, A. Badolato, and A. Imamoglu, Nature Physics 521, 101 (2007).
[10] M. H. Mikkelsen, J. Berezovsky, N. G. Stoltz, L. A. Coldren, and D.D.Awshalom, Nature Physics 736, 1 (2007).
[11] H. P. Seigneur, M. N. Leuenberger and W. V. Schoenfeld, J. Appl. Ph ys. 104, 014307 (2008).
[12] I. A. Merkulov, Al. L. Efros, and M.Rosen, Phys. Rev. B 65, 205309 (2002).
[13] A. V. Khaetskii, D. Loss, and L.Glazman, Phys. Rev. Lett. 88, 186802 (2002).
[14] A. C. Johnson, J. R. Petta, J. M. Taylor, A. Ya- coby, M. D. Lukin, C. M. Marcus, M. P. Hanson, and A.C.Gossard, Nature 435, 925-928 (2005).
[15] A. S. Bracker, E. A. Stina, 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).
[16] P. F. Braun, X. Marie, L. Lombez, B. Urbaszek, T. Amand, P. Remucci, V. K. Kalevich, K. V. Kavokin, O. Krebs, P. Voisin, and Y. Masumoto, Phys. Rev. Lett. 94, 116601 (2005).
[17] F. H. L. Koppens, J. A. Folk, J. M. Elzerman, R. Hanson, L. H. Willems van Beveren, I. T. Vink, H. P. Tranitz, W. Wegscheider, L. P. Kouwenhoven, and L. M. K. Vandersypen, Science 309, 1346-1350 (2005).
[18] J. R. Petta, A. C. Johnson, J. M. Taylor, E. A. Laird, A. Yaocoby, M. D. Lukin, C. M. Marcus, M. P. Hanson, and A. C. Gossard, Science 309, 2180-2184 (2005).
[19] P. Irvin, P. S. Fodor, and J. Levy, Optics Express 15, 11756 (2007).
[20] M. Sugita, S. Machida, and Y. Yamanoto, arxiv:quant-ph/0301064v1 (2003).
[21] J. R. Ackerhalt and K. Bzazewski, Phys. Rev. A 12, 2549 (1975).
[22] L. Mandel and E. Wolf, Optical Coherence and Quantum Optics (Cambridge University Press, New York, 1995).
[23] N. Korolkova, G. Leuchs, R. Loudon, T. C. Ralph, and C. Silberhorn, Phys. Rev. A 65, 052306 (2002).
[24] Y. Takahashi, K. Honda, N. Tanaka, K. Toyoda, K. Ishikawa, and T. Yabuzaki, Phys. Rev. A 60, 4974 (1999).
[25] A. V. Khaetskii, D. Loss, and L. Glazman, Phys. Rev. Lett. 88, 186802 (2002).
[26] J. Schliemann, A. V. Khaetskii, and D. Loss, Phys. Rev. Lett. 66, 245303 (2002).
[27] A. Khaetskii, D. Loss, and L. Glazman, Phys. Rev. B 67, 195329 (2003).
[28] W. A. Coish and D. Loss, Phys. Rev. B 70, 195340 (2004).
[29] W. Zhang, V. V. Dobrovitski, K. A. Al-Hassanieh, E. Dagotto, and B. N. Harmon, Phys. Rev. B 74, 205313 (2006).
[30] G. Khitrova, H. M. Gibbs, M. Kira, S. W. Koch, and A. Scherer, Nature Physics 2, 81 (2006).