Coherent transfer of light polarization to electron spins in a semiconductor

| 著者 | この研究に貢献した皆様の名前を列挙します。
<table>
<thead>
<tr>
<th></th>
<th></th>
</tr>
</thead>
<tbody>
<tr>
<td>物理ジャーナルのタイトル</td>
<td>Physical Review Letters</td>
</tr>
<tr>
<td>卷</td>
<td>100</td>
</tr>
<tr>
<td>号</td>
<td>9</td>
</tr>
<tr>
<td>頁</td>
<td>096602</td>
</tr>
<tr>
<td>年</td>
<td>2008</td>
</tr>
<tr>
<td>URL</td>
<td><a href="http://hdl.handle.net/10097/53524">http://hdl.handle.net/10097/53524</a></td>
</tr>
</tbody>
</table>

doi: 10.1103/PhysRevLett.100.096602
We demonstrate that the superposition of light polarization states is coherently transferred to electron spins in a semiconductor quantum well. By using time-resolved Kerr rotation, we observe the initial phase of Larmor precession of electron spins whose coherence is transferred from light. To break the electron-hole spin entanglement, we utilized the big discrepancy between the transverse g factors of electrons and light-holes. The result encourages us to make a quantum media converter between flying photon qubits and stationary electron-spin qubits in semiconductors.

None of the physical implementations of quantum bits or qubits is suitable for all aspects of quantum-information technology. Quantum-information interfaces [1,2] connecting those qubits are thus needed to place the right qubit at the right position. Photons are the natural messenger qubits for fast and reliable communication [3], whereas electron spins in a semiconductor are suitable processor qubits for stable and scalable computation [4–6]. Quantum media conversion (QMC) between the photon qubit and the electron-spin qubit will drastically extend the potential of quantum-information and -communication technologies.

Quantum coherence is one of the essential features of quantum phenomena and is a central issue in quantum-information technology. Researchers exploring the feasibility of semiconductor-based quantum computation have demonstrated that not only the optical dipole coherence of excitons [7,8], but also electron-spin coherence can be controlled optically [9–13] and electrically [5,6]. The electron-spin coherence time, or transverse spin-relaxation time $T_2^*$, is generally much longer than the optical dipole coherence time and could be further prolonged by storing the electron-spin coherence in a nuclear spin [14]. On the other hand, researchers exploring the feasibility of semiconductor-based quantum communication have shown that triggered single photons [15,16] and polarization-entangled photon pairs [17–20] can also be generated in semiconductors. One of the remaining challenges in developing all-semiconductor quantum-information and -communication devices is to transfer the quantum coherence of the superposition state of a photon (based on polarization degree of freedom) to that of an electron in a semiconductor (based on spin degree of freedom). Although the physics of optical orientation has a long history [21], little attention has been paid to the coherent transfer of the polarization state of light to electron spins.

In this Letter, we demonstrate for the first time coherent transfer of light polarization to electron spins. We not only show the projective up and down spin states or spin population as reported before [22–25] but also show that in a V-shaped three-level system with Voigt geometry an arbitrary coherent superposition state can be transferred from light polarization to electron spins. Although we used classical light and a spin ensemble, this coherent transfer is the first demonstration of the conditions needed for transferring the quantum state [26–28] of a photon to an electron.

We begin by summarizing the operating principle of our QMC scheme (Fig. 1). Any light polarization state is represented by a Bloch vector in the Bloch sphere (BS) and stationary electron-spin qubits in semiconductors.
the resonant transition between the \([-\chi]_{\text{LH}}\) state and the degenerated electron states, the same selection rules lead to the transition from \(\alpha|\sigma^+\rangle + \beta|\sigma^-\rangle\) to \((\alpha|\downarrow\rangle + \beta|\uparrow\rangle) \otimes |\chi]_{\text{LH}}\). Since the transferred electron-spin state is not entangled with the LH state, its coherence time is not limited by the short LH spin coherence time or the fast relaxation to the ground heavy-hole (HH) state. This is the key point of our QMC scheme. It is essential to use the LH states since we cannot sufficiently lift the Kramer’s degeneracy of the HH state by applying an in-plane magnetic field because of the very small transverse \(g\) factor [29]. The electron coherence time is limited by the electron-hole exchange interaction, which can be greatly extended by extracting a hole while leaving an electron in a quantum dot [9, 12, 22, 23], which is usually visualized by time-resolved Kerr rotation measurements [9, 10]. Although the Kerr rotation provides information only about the spin projection or population along the growth axis, we can infer the spin orientation of the superposition state from the initial phase of the spin precession. Here, we show how the inferred initial phase of the generated electron spin is related to the pump-light polarization. We first selected six light polarization states \(|\sigma^+\rangle, |\sigma^-\rangle, |D^+\rangle, |D^-\rangle\) (±45° degree polarization), and \(|H\rangle, |V\rangle\) (horizontal and vertical polarization, respectively) and measured the temporal evolution of the Kerr rotation angle \(\theta_K\) under \(B_x = 7\) T [Fig. 2(c)]. The lower-energy side of the LH exciton peak \((\Delta E_L = 796.2\) nm) was excited with one of the above-mentioned polarizations and probed at \(\lambda_p = 795.7\) nm with the \(H\) polarization [Fig. 2(a)]. Four of the states perpendicular to the \(x\) axis \((|\sigma^+\rangle, |D^+\rangle, |\sigma^-\rangle,\) and \(|D^-\rangle\) provided approximately the same amplitude, while the phases were sequentially shifted by the angle of \(\pi/2\). These results suggest that the \(|D^-\rangle\) states, which are coherent superpositions of \(|\sigma^+\rangle\) and \(|\sigma^-\rangle\) states \(|\sigma^+\rangle \pm i|\sigma^-\rangle\), create equivalent superpositions of the \(|\uparrow\rangle\) and \(|\downarrow\rangle\) states \(|\uparrow\rangle \pm i|\downarrow\rangle\) (normalization omitted for brevity).
It is natural that the states along the $x$ axis (|$H$⟩ and |$V$⟩) provided negligible amplitude because those state vectors are parallel to the field. From the decay of precession amplitude seen in Fig. 2(d), the $T_2^*$ is estimated to be 160 ps. In contrast to the LH excitation, the HH excitation ($\lambda_K = \lambda_P = 800.8$ nm) gave the conventional result [Fig. 2(e)], where the $\sigma^\pm$ pump corresponding to the projected states |$\uparrow\rangle_e$ provided significant amplitude but the others did not. The lack of the amplitude seen with $D^+$ and $D^-$ pumping indicates that superposition of the |$\downarrow\rangle_e$ and |$\rdownarrow\rangle_e$ states could not be created by the HH excitation. We verified the one-to-one correspondence between the light polarization phase $\varphi_{iK}^e$ and the electron-spin phase $\varphi_{iK}^e$ as shown in Fig. 3. The phases $\varphi_{iK}^e$ and $\varphi_{iK}^e$ are, respectively, the angles of the state vectors in the PS and the BS, such that |$D^+$⟩ = $\cos(\varphi_{iK}^e/2)|\sigma^+⟩ + i\sin(\varphi_{iK}^e/2)|\sigma^-⟩$ as shown in Fig. 1(a). The electron phases $\varphi_{iK}^e$ inferred from the peaks well coincide with the pump-light phases $\varphi_{iK}^e$. This coincidence is evidence of coherent polarization transfer.

To show the role of the lift of the hole state Kramer’s degeneracy induced by the in-plane magnetic field, we measured $B_x$-field dependence of $\theta_K$ with the $D^+$ polarization [Fig. 4(a)]. The times to reach maximum and minimum signals are well reproduced by the theoretical calculations shown by the dashed curves. The first maximum points are delayed by $\pi/2$ for all $B_x$ as in Fig. 2(c). Note that as the magnetic field decreases, the precession amplitude decreases much faster than $T_2^*$, indicating that the electron-spin coherence is gradually lost through the entanglement with the degenerated hole spin as in the HH case. Although the exchange interaction between the electron and the LH in $|-x⟩_LH$ state might act as an additional effective magnetic field along the $x$ axis, as in the case of HH [22], it is not remarkable because of the fast LH relaxation to the HH. The distinguishability of the LH states is demonstrated in Fig. 4(b), where $\theta_K$ is measured at $\Delta t = 12$ ps ($\pi/2$ phase) with the $D^+$ polarization at $B_x = 7$ T is plotted against both the pump and probe wavelengths. The splitting of the peaks along the pump wavelength originates from the LH Zeeman splitting. The

![FIG. 2](color). (a) PL (blue) and PLE (red) spectra detected at HH exciton emission under $B_x = 7$ T. The broken lines show the spectra of the pump (red) and probe (blue) lights. (b) Pump power dependence of the Kerr rotation angle ($\theta_K$) with $D^+$ under $B_x = 7$ T at $\Delta t = 12$ ps. Error bars show standard deviation. (c) Temporal evolution of $\theta_K$ with six basis polarizations under $B_x = 7$ T. Each of the consecutive phase shifts in the cyclic order of $|\sigma^+⟩$ (red) → $|D^+⟩$ (green) → $|\sigma^-⟩$ (blue) → $|D^-⟩$ (orange) → $|\sigma^+⟩$ in $\pi/2$ steps coincides well with that of electron-spin states expected from the pump-light polarizations. In contrast, excitation with the $|H⟩$ (pink) and $|V⟩$ (sky blue) polarization, which correspond to the $B_0$ direction, generate negligible signals. (d) Large-time-scale view of Fig. 2(c). (e) HH-excitation case for comparison. Only $|\sigma^\pm⟩$ excitations generate relatively large signals, indicating that spins are always projected along the $z$ axis.

![FIG. 3](color). Time evolution of $\theta_K$ as a function of pump-light polarization under $B_x = 7$ T. The linear streak pattern indicates coincidence of the phase angle of the electron spins $\varphi_{iK}^e$ to that of the light polarization $\varphi_{iK}^e$. Definitions of the phases are shown in Fig. 1(a).
FIG. 4 (color). (a) Time evolution of $\theta_K$ as a function of $B_z$ with $D^+$ pumping. The cut of $\theta_K$ at $B_z = 4$ is shown by the dots fitted to a damped sine function. The dashed curves are the calculated tracks of the $B_z - \Delta t$ values needed to detect electrons aligned to $|\pm z\rangle_e$, when an electron transverse $g$ factor of $-0.21$ is assumed. (b) $\theta_K$ as a function of pump/probe wavelengths with $D^+$ at $\Delta t = 12$ ps under $B_z = 7$ T. Dots show the pump spectrum of $\theta_K$ at the probe wavelength used for all the LH measurements (795.7 nm) with a curve fitted by a two-component Gaussian.

low (high)-energy peak showing the positive (negative) sign corresponds to the LH spin state $|\pm x\rangle_{LH} = |\pm i\rangle_{LH}$, where the pump-light polarization state $|D^+\rangle$ is transferred to the electron-spin state $|\pm y\rangle_e = |\pm i\rangle_e$, respectively. The pump and probe wavelengths in the above experiments were chosen to have the peak signal. We can fit the pump spectrum of $\theta_K$ (dots) at a fixed probe wavelength by using two Gaussian functions with a separation given by the calculated transverse LH $g$ factor of $-3.5$ (solid curve). This well-separated spectrum is not observed in the HH-excitation case, where because of the Kramer’s degeneracy, only one peak is observed with any polarization of pump light.

The coherent polarization transfer demonstrated here is essentially different from the previously demonstrated coherent electron-spin controls based on an optical Stark effect [10], spin-flip Raman scattering [11], or Rabi oscillation [12], each of which needs intense light as a driving force defined and is therefore not suitable for transferring the phase state of a single photon. Our QMC scheme is not limited to the transfer of a single-particle state but can be generalized to the transfer of a two-particle state or an entangled state, which is the kind of transfer needed for quantum repeaters [33,34] or routers for scalable quantum computation. Although the wavelength used here is about half of the telecom wavelength, QMC between photons of different wavelengths by a nonlinear up-conversion process has been shown to be possible [2]. The transfer of light polarization only in the $y$-$z$ plane in the PS as shown here is sufficient for implementation of a quantum repeater based on the Bennett-Brassard 1984 cryptographic scheme [2,3] by converting the time-bin phase coding to the polarization coding.

In conclusion, we have demonstrated coherent polarization transfer from light to electron spins in a semiconductor. From the phase of spin precession, we inferred the initial phase that should correspond to the phase of light polarization. The electron-hole spin entanglement was broken by the big discrepancy between the $g$ factor of the electron and the $g$ factor of the light hole. Although we have demonstrated only a necessary condition of QMC, this is a step toward the realization of all-semiconductor quantum interfaces between a photon and an electron spin.

We are grateful to Professors Hideo Ohno, Toshihide Takagahara, Keiji Ono, and Jaw-Shen Tsai for fruitful discussions. This work was supported in part by the SCOPE program (No. 41402001) of MIC Japan, a grant from MEXT Japan (No. 16710061), and a grant from NEDO Japan.