Photo-induced spin filtering in a double quantum dot

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We investigate the spin-dependent electron dynamics in a double quantum dot driven by sub-picosecond asymmetric electromagnetic pulses. We show analytically that applying the appropriate pulses, specified here, allows a spin separation on a femtosecond time scale in the sense that states with a desired spin projection are localized mainly on one of the dots. It is shown how to maintain in time this photo-induced spin-dependent filtering.

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Studies of the electron dynamics in quantum well structures triggered by time-dependent electromagnetic fields revealed a wealth of interesting phenomena. The driving field can be a continuous wave (cw) laser, which can be employed to demonstrate processes such as the coherent suppression of tunneling or the low-frequency generation in symmetric double quantum well (DQW) systems. Coherent control schemes can be used for trapping of electrons in different quantum wells or driving transient current bursts due to quantum interferences.

Another possibility to control the electron dynamics on the subpicosecond time scale is to utilize ultrashort highly asymmetric electromagnetic pulses. Such pulses with the electric field consisting of an ultrashort and strong oscillation half-cycle followed by a weak and much longer oscillation tail of the opposite polarity are called half-cycle pulses (HCPs). If the amplitude of the tail is small enough it hardly influences the dynamics. The HCP can be considered then as an unipolar electromagnetic pulse with a small duration \( \tau_d \) given by the duration of its first ultrashort and strong half-cycle.

Quantum mechanically, the interaction between a HCP and a charge carrier can be approximately described as a transformation of the carrier wave function \( \Psi(p) \rightarrow \Psi(p + \Delta p) \) in the momentum representation, while in the coordinate representation we have \( \Psi(x) \rightarrow \Psi(x)e^{-i\Delta p x/\hbar} \). It is possible to localize a charge carrier on the subpicosecond time scale by applying HCP sequences. Here we deal with the question of whether it is possible to separate spin states using ultrashort light pulses, an issue which is of relevance for applications in ultrafast spintronics and quantum computing devices.

The effect of the intrinsic spin-orbit coupling on the light-driven charge and spin dynamics in a DQW system has been investigated recently. It has been shown that the charge distribution dynamics induced by applied HCPs become spin-dependent. Below we show that the spin-dependent localization of the charge carriers can be achieved by using the effect of the field-induced spin-orbit coupling. Because of this field-induced coupling an appropriately directed HCP can generate a spin-dependent coordinate transfer for a free carrier confined to one dimension so that we have \( \Psi_\uparrow(x) \rightarrow \Psi_\uparrow(x + \Delta x) \) for the spin-up state \( \uparrow \) whereas \( \Psi_\downarrow(x) \rightarrow \Psi_\downarrow(x - \Delta x) \) for the spin-down \( \downarrow \) state. In the case of an electron confined to two coupled quantum dots we will show how to apply appropriate HCPs to control coherently the electron position depending on its spin degree of freedom. The required photon fields are feasible. Nowadays, light pulses with durations in the range of femtosecond and even attoseconds can be generated.

We consider a one-dimensional symmetric DQW,
which can be realized, e.g., on the basis of semiconductor heterostructures and can be viewed as a double quantum dot (DQD) system. The geometry of the system and the applied electric field pulses are illustrated in Fig. 1. The left (right) quantum dot is located in the region \( x < 0 \) (\( x > 0 \)). Due to the quantum confinement with appropriately high energy barriers in the \( y \)- and \( z \)-directions, the carrier motion is confined only to the \( x \)-direction. The control scheme is based on applying two types of HCPs, which are polarized either along the \( x \)- or along the \( y \)-axis, respectively. Our target is to steer spatially the carrier depending on its spin state, in particular on its spin projection to the \( z \)-axis. We note that both QDs experience the same pulse field.

The system Hamiltonian is given by \( H = H_0 + H_K(t) + H_{SO}(t) \), where \( H_0 \) is the Hamiltonian of a free electron in the DQD. The interaction of the electron with the electric field \( E_x(t) \) that is polarized along the \( x \)-axis, is given by \( H_K(t) = x \cdot eE_x(t) \), where \( e \) is the elementary charge. HCPs polarized in this direction are called here \( x \)-HCPs. The electric field \( E_y(t) \) of the corresponding \( y \)-HCPs, polarized along the \( y \)-axis, induces the field-dependent Rashba-type spin-orbit interaction \( H_{SO}(t) = \frac{\hbar}{4} \alpha_{SO} E_y(t) \sigma_z \), for the electron in the DQD. Here \( \alpha_{SO} \) is the coupling coefficient, \( p_z \) is the component of the momentum operator along the \( z \)-axis and \( \sigma_z \) is the Pauli matrix. We have omitted the term proportional to \( p_x \sigma_x \) in the expression for \( H_{SO}(t) \) because the motion in the \( z \)-direction is restricted by the carrier confinement. The electric field pulse shape can be selected, e.g., as \( E(t) = E_0 \sin^2(\pi t/\tau_d) \) for \( 0 < t < \tau_d \) (the field is zero outside of this time range), where \( E_0 \) is the electric field amplitude and \( \tau_d \) is the HCP duration. The pulse is centered at \( t = t_p = \tau_d/2 \) when the electric field reaches its maximum.

In our treatment we may ignore effects of elastic scattering and electron-phonon interaction because they take place on much longer time scales compared with the times relevant for this study.\(^{27,28}\) If the duration \( \tau_d \) of the HCP is much shorter than the characteristic time \( \tau_c = 2\pi/\omega_c \) (\( \omega_c \) is the frequency corresponding to the energy difference between the ground and first excited states of the field-free system) of the undriven system, the impulsive (sudden) approximation (IA)\(^{29,30}\) can be applied when solving the general time-dependent Schrödinger equation. In the framework of the IA the electron interaction with a \( x \)-HCP can be written as

\[
H_K(t) = x\Delta p\delta(t - t_p),
\]

where \( \delta(x) \) denotes the Delta-function and \( \Delta p \) corresponds to the amount of transferred momentum, which is given by \( \Delta p = e \int_{-\infty}^{\infty} E_x(t)dt \). In the case of the sine-square pulse shape field we have \( \Delta p = eE_0\tau_d/2 \). For the Hamiltonian of the electron interaction with a \( y \)-HCP the IA leads to

\[
H_{SO} = p_x\sigma_z\Delta x\delta(t - t_p),
\]

where \( \Delta x = \frac{\hbar}{2}\alpha_{SO} \int_{-\infty}^{\infty} E_y(t)dt \) accounts for a sudden coordinate transfer to the electron. For the sine-square pulse shape we have \( \Delta x = \frac{\hbar}{2}\alpha_{SO} E_0\tau_d \).

The DQD parameters are chosen such that the two lowest energy levels are well separated from the other states. Hence, a two-level system approximation (TLSA) including the spin can be applied. The effective strengths of the \( x \)-HCP and the \( y \)-HCP can be characterized by the dimensionless parameters \( \beta = x_12\Delta p/h \) and \( \alpha = p_12\Delta x/h \), respectively. Here \( x_12 = (1|x^2|2) \) and \( p_12 = (1|p_x|2)/i \) are real coordinate and momentum matrix elements.\(^{22}\)

The time evolution depends strongly on the initial conditions. Two situations are considered: (a) the initial condition, which corresponds to an electron being completely localized in the left well (tunneling initial condition) and a zero average spin along the \( z \)-axis, i.e. \( \langle \sigma_z \rangle = 0 \), and (b) the initially delocalized state (optical initial condition) with the same spin properties, which belongs to the ground state of the system. In previous studies it was shown that a \( x \)-HCP with appropriate pulse parameters preceded by a short period of free propagation can be used to create the optical initial condition from the tunneling one on an ultrafast timescale.\(^{59}\) A spin polarization can be created if we then apply a \( y \)-HCP at the time moment just after the optical initial condition was created. To maintain the spin polarization an appropriate periodic train of \( x \)-HCPs can be used, in analogy to the maintenance of spin-unpolarized states.\(^{59}\)

The results of the corresponding numerical simulation are illustrated in Fig. 2 where the spin-resolved time-dependent probability of finding an electron in the left well \( |P_L^\uparrow(t), P_L^\downarrow(t)\rangle \) is shown. We start from the tunneling initial condition (both spin states have the same probability). After a time \( t_0 = 0.25\tau_c \) of free propagation a \( x \)-HCP with \( \beta = \pi/4 \) is applied, which transfers the system into the ground state (both spin states have the same probability \( P_L^\uparrow = P_L^\downarrow = 0.25 \)). At \( t = 0.5\tau_c \) we
apply a $y$-HCP with $\alpha = \pi/4$ and obtain immediately a nearly perfect spin separation. This means that the two spin states are localized in different wells: $P_{L\uparrow} = 0.5$ and $P_{L\downarrow} = 0$. This separation is then maintained by applying a periodic $x$-HCP train with the effective strength $\beta = \pi/2$, period $T = 0.1t_c$ and the first HCP centered at $t = 0.5t_c$. As a result, the mean values of the probabilities after the separation, averaged over the time period $2t_c$, are: $\langle P_{L\uparrow} \rangle = 0.495$ and $\langle P_{L\downarrow} \rangle = 0.046$. Thus a very good spin-separation is stabilized in time. This is also illustrated in Fig. 3b where the dynamics of the components of the spin-polarization $\vec{A}_L(t)$ in the left well are shown. Until the time moment of applying the $y$-HCP at $t = 0.5t_c$ the spin polarization in the left well is oriented along the $y$-axis. Just after this time moment the spin polarization turns abruptly into the $z$-direction. Its orientation oscillates then in a small solid angle close to the $z$-axis under the action of the maintaining $x$-HCP train. The time-averaged (over $2t_c$) spin polarization in $z$-direction is $0.982$. The corresponding polarization vector trajectory on the unit sphere is shown in Fig. 3b.

Let us discuss shortly the conditions required for an experimental realization. As a first choice for the model DQD we may consider a typical GaAs-based symmetric DQW as in Ref. 8, where $\vec{r}_a \approx 0.67$ ps. We may take $\tau_A = 40$ fs ($\tau_B$ should be varied in a way that the light can still drives effectively the relevant dynamics). Values of $\alpha_{30}$ for different $\alpha_3\beta_5$ semiconductor materials are listed in Ref. 13. For GaAs we have $\alpha_{30} \approx 4$ A$^2$. Then we can estimate the maximum required amplitude of HCPs to $E_0 \sim 10^5$ V/cm. Such high and short fields have become available recently, but one may wish to attenuate them in order to avoid undesirable effects. In particular, strong fields may alter the band structure (e.g., via the dynamic Franz-Keldysh or the AC Stark effects) that is not accounted for here. Narrow gap semiconductors, in particular InSb, are suitable as a replacement for GaAs because of the much higher value of $\alpha_{30}$.

We have shown that it is possible to separate the spin states of electrons in a double quantum dot by only two different ultrafast light pulses. We achieved a nearly perfect spin polarization in the particular direction which can be maintained for a desired period of time by applying an additional pulse train. Such light-induced spin filtering can be realized on a sub-picosecond time scale that can be of relevance for designing ultrafast spintronic and spin-qubit devices. Heterostructures based on narrow gap semiconductors with strong spin-orbit interaction are good candidates for an experimental demonstration of the predicted phenomena.

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33We select the wave functions to be real.
34These values are required for HCPs polarized in the $y$-direction, for which the motion is limited by the quantum confinement.
Required amplitudes for x-HCPs are two orders of magnitude lower.