

# Proposal for fast optical control of spin dynamics in a quantum wire

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The spin dynamics in a quantum wire with a Rashba spin-orbit interaction (SOI) is shown to be controllable via subpicosecond electromagnetic pulses shaped appropriately. If the light polarization vector is along the wire's direction, the carriers experience a momentum boost while the phase coherence in different spin channels is maintained, a fact exploitable to control the speed of a photodriven spin field effect transistor. A photon pulse with a polarization vector perpendicular to the wire results in a spin precession which is comparable to that due to the Rashba SOI and is tunable by the pulse field parameters, an effect utilizable in optically controlled spintronics devices.

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## I. INTRODUCTION

Optical semiconductor devices are indispensable components of nowadays technology<sup>1</sup> with applications ranging from optical fiber communication systems to consumer electronics. A new impetus is expected from spintronics devices,<sup>2,3</sup> i.e., from exploiting the carriers spin in addition to their charge for efficient operation or to realize new functionalities. While the optical control and manipulation of charges in conventional semiconductors<sup>4</sup> have been the key for the realization of ultrafast electronic devices, an analogous optical control of the spins is, however, not straightforward. The appropriate electromagnetic pulses are U(1) fields while spins belong to SU(2) fields. A laser pulse does not seem thus to couple to the spins directly, meaning a less efficient optical coupling to the spins than to the charges. Indeed, the key to the optical spin manipulation are inherent interactions in the system that couple the charge to the spin such as spin-orbit interaction (SOI).<sup>5,6</sup> Along this line we present here a way for an ultrafast control of the spin dynamics in a conventional spin field effect transistor (SFET) (Refs. 7 and 8) driven by shaped electromagnetic pulses. The SFET relies on the Rashba SOI to perform a controlled rotation of a carrier spin that traverses a FET-type device<sup>7,8</sup> with two magnetic leads (cf. Fig. 1). The conductance depends on the achieved rotation angle  $\Delta\theta_0$  at the drain lead. Here we propose two ways to control the time dependence of this rotation angle and hence the operation of SFET. The key ingredient is the use of asymmetrically shaped linearly polarized electromagnetic pulses.<sup>9-12</sup> The pulse consists of a very short, strong half cycle followed by a second long (compared

to the ballistic transverse time) and a much weaker half cycle of an opposite polarity.<sup>9-11</sup> Hence such pulses are often called half-cycle pulse (HCP). Experimentally the achieved asymmetry ratio of the positive and negative amplitudes can be 13:1, the peak fields can reach several hundreds of kilovolt per centimeter and have a duration  $t_p$  in the range between nanoseconds and subpicoseconds. The interaction of HCP with matter is particular in that it delivers a definite amount of momentum boost to the carriers along the optical polarization axis.<sup>6,12</sup> In case this axis is along the conductive channel we find that both the spin-precession frequency and the carrier speed increase upon irradiation but  $\Delta\theta_0$  remains unchanged. The operation speed is thus pulse controllable. If the optical polarization vector is perpendicular to the carriers propagation direction  $\Delta\theta_0$  and hence the SFET operation is determined by the pulses parameters, as shown explicitly below.

## II. MODEL SFET

We follow Refs. 7 and 8 and focus on the central conductive region of the SFET that can be considered as a one-dimensional (1D) quantum wire (QW) of length  $L$  with a SOI. The ferromagnetic leads serve as spin injector and spin detector separated from QW by an insulating barrier to achieve a higher spin injection efficiency (Fig. 1). Recently a SFET structure similar to Fig. 1 has been realized experimentally.<sup>8</sup> The experimental findings are in line with the predictions of the stationary version of the present work. In the work presented below, however, we consider 1D QW, for spin-flip transitions between the first and second subband is negligible for the system studied in this work.

The inversion asymmetry of the confining potential results in Rashba SOI (Ref. 13)  $H_R$ . For a two-dimensional electron gas (2DEG)  $H_R$  reads<sup>13</sup>

$$H_R = \alpha_R^0 (\sigma_z k_x - \sigma_x k_z), \quad (1)$$

where  $\sigma_i, i=x, y, z$  are Pauli matrixes,  $\alpha_R^0 = r_R \epsilon_0$  is the static Rashba SOI coefficient which is proportional to the perpendicular electric field  $\epsilon_0$  resulting from band bending, and  $r_R$  is a material-specific prefactor.<sup>14</sup> Under equilibrium conditions, spin transport in such a device is investigated extensively,<sup>7,14</sup> in brief, one chooses for 1D QW the  $z$  axis as

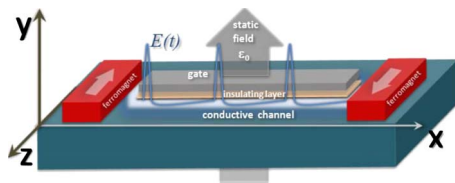


FIG. 1. (Color online) Schematics of the optically driven spin field transistor. Ferromagnetic leads are separated from the conductive channel by a tunneling barrier to enhance the spin injection efficiency. A metallic gate is used to tune the Rashba SOI via a static field  $\epsilon_0$ .  $E(t)$  is the time-dependent electric field.

the spin quantization axis (Fig. 1), meaning that

$$H_R = \alpha_R^0 \sigma_z k_x = \mu_B \sigma_z \tilde{B}_z,$$

where  $\mu_B$  is Bohr's magneton and  $\tilde{B}_z = \alpha_R^0 k_x / \mu_B$  is an effective magnetic field along  $z$ . This results in the spin splitting  $2\alpha_R^0 k_x$  between carriers injected with spin polarization parallel or antiparallel to  $z$ . The phase difference while passing through the length  $L$  is

$$\Delta\theta_0 = (k_{x\uparrow} - k_{x\downarrow})L = \frac{2m^* \alpha_R^0 L}{\hbar^2}.$$

The eigenenergies and the eigenstates are, respectively,

$$E_{k_x\mu} = \frac{\hbar^2 k_x^2}{2m^*} + \mu \alpha_R^0 k_x$$

and

$$|k_x\mu\rangle = \frac{1}{\sqrt{L}} e^{ik_x x} |\eta_\mu\rangle,$$

where  $|\eta_\mu\rangle$  is the spin states. This is the original idea of the Ref. 7.

Considering the spins to be injected aligned along the  $x$  or  $y$  directions, they precess around  $\tilde{B}_z$ . In Heisenberg picture, the spin operators vary with time as

$$\dot{\sigma}_{x(y)}(t) = \mp \omega_{k_x} \sigma_{y(x)}(t),$$

where

$$\omega_{k_x} = \frac{2\alpha_R^0 k_x}{\hbar}$$

is the precession frequency. Therefore,

$$\sigma_\pm(t) = \sigma_x \pm i\sigma_y = \sigma_\pm(0) e^{\pm i\omega_{k_x} t}.$$

Let us specify the initial orientation, say  $\sigma_y(0)=0$ , therefore

$$\sigma_x(t) = \sigma_x(0) \cos(\omega_{k_x} t) \quad \text{and} \quad \sigma_y(t) = \sigma_x(0) \sin(\omega_{k_x} t).$$

An initial spin along the  $x$  direction rotates anticlockwise with the angular frequency  $\omega_{k_x}$ . The accumulated angle through the length  $L$  is  $\Delta\theta_0 = \frac{2m^* \alpha_R^0 L}{\hbar^2}$ , which is exactly equal to the phase shift for the spin along  $z$ .

### III. FIRST DYNAMIC CASE

Having outlined the equilibrium case, we apply to the quantum wire a linearly polarized HCP with the vector potential  $\mathbf{A} = \mathbf{e}_x A(t)$ . The polarization vector  $\mathbf{e}_x$  is along the  $x$  direction. Thus,

$$\mathbf{E}(t) = -\mathbf{e}_x \frac{\partial A(t)}{\partial t} = F a(t) \mathbf{e}_x,$$

where  $F$  is the peak amplitude of the electric field and  $a(t)$  describes the pulse profile. For simplicity we use  $k$  to indicate  $k_x$  in this section. The single-particle Hamiltonian reads

$$H = \frac{\pi^2}{2m^*} + V(y,z) - e\Phi + \frac{\alpha_R^0}{\hbar} (\boldsymbol{\sigma} \times \boldsymbol{\pi})_{y,k_z=0}, \quad (2)$$

where  $\boldsymbol{\pi} = \mathbf{p} + e\mathbf{A}(\mathbf{r}, t)$  and  $\mathbf{p} = -i\hbar\Delta$  is the momentum operator. The second term in Eq. (2) is the QW confinement potential, the third term is from the scalar potential of the pulse field, and the fourth term in Eq. (2) is the Rashba SOI. We write  $H = H^0 + H'$ , with

$$H^0 = \mathbf{p}^2/2m^* + V(y,z) + (\alpha_R^0/\hbar) (\boldsymbol{\sigma} \times \mathbf{p})_{y, k_z=0}$$

and

$$H' = (e/2m^*) (\mathbf{p} \cdot \mathbf{A} + \mathbf{A} \cdot \mathbf{p}) + (e/2m^*) \mathbf{A} \cdot \mathbf{A} - e\Phi(\mathbf{r}) + (e\alpha_R^0/\hbar) \times (\boldsymbol{\sigma} \times \mathbf{A})_y.$$

$H^0$  is the pulse-free single-particle Hamiltonian, and  $H'$  is the time-dependent part. Choosing a gauge where  $\Phi=0$ ;  $\mathbf{p} \cdot \mathbf{A} + \mathbf{A} \cdot \mathbf{p} = 2A_x p_x$  and  $(\boldsymbol{\sigma} \times \mathbf{A})_y = \sigma_z A_x$  results in

$$H' = \frac{e}{m^*} A_x p_x + \frac{e^2}{2m^*} A_x^2 + \frac{e\alpha_R^0}{\hbar} \sigma_z A_x. \quad (3)$$

Defining the spinor field operator as  $\hat{\psi}(x) = \sum_{k\mu} c_{k\mu} |k\mu\rangle$ , where  $c_{k\mu}$  is the annihilation operator for the states  $|k\mu\rangle$ , in the second quantization form, we obtain

$$H' = \sum_{k\mu} \left( \frac{e\hbar}{m^*} A_x k + \frac{e^2}{2m^*} A_x^2 + \mu \frac{e\alpha_R^0}{\hbar} A_x \right) c_{k\mu}^\dagger c_{k\mu}. \quad (4)$$

The total Hamiltonian reads  $H = \sum_{k\mu} \epsilon_{k\mu}^c(t) c_{k\mu}^\dagger c_{k\mu}$ , with the time-dependent transient energy [measured with respect to the ground state of  $V(y,z)$ ],

$$\epsilon_{k\mu}^c(t) = \frac{1}{2m^*} \left\{ \left[ \hbar k + eA_x(t) + \mu \frac{p_{so}}{2} \right]^2 - \frac{p_{so}^2}{4} \right\} \quad (5)$$

and

$$p_{so} = \hbar k_{so} = \frac{2m^* \alpha_R^0}{\hbar}. \quad (6)$$

To obtain the momentum distribution upon a short pulse application (say at  $t=0$ ) we may proceed as in Refs. 6 and 12 and expand the single-particle excited state  $\Psi_{k_0\mu_0}(x,t)$  that starts from the state labeled by  $|k_0\mu_0\rangle$  in terms of the stationary eigenstates

$$\Psi_{k_0\mu_0}(x,t) = \sum_{\mu} \int dk C_{k\mu}(k_0, \mu_0, t) |k\mu\rangle e^{-iE_{k\mu} t/\hbar}. \quad (7)$$

For a sudden excitation,  $\Psi_{k_0\mu_0}(x, t=0^+)$  right after the pulse evolves from the state before the pulse  $\Psi_{k_0\mu_0}(x, t=0^-)$  as<sup>6</sup>

$$\Psi_{k_0\mu_0}(x, t=0^+) = e^{ix\bar{p}} \Psi_{k_0\mu_0}(x, t=0^-).$$

Thus

$$C_{k\mu}(k_0, \mu_0, t=0^+) = C_{k-\bar{k}, \mu}(k_0, \mu_0, t=0^-), \quad (8)$$

where

$$\hbar\bar{k} = \bar{p} = -eF \int_{-t_p/2}^{t_p/2} a(t') dt'$$

is the momentum boost delivered to the carrier by the pulse (the second weak and long half cycle of the pulse acts as a weak dc off-set field). For  $t < 0$ , Eq. (7) may stand for the injected electron state in terms of the stationary states. For  $C_{k'\mu'}(k, \mu, t < 0) = \delta(k - k') \delta_{\mu\mu'}$ , the injected electron occupies a single eigenstate. In this case the wave function after the pulse is

$$\Psi_{k_0\mu_0}(x, t > 0) = e^{i(\hbar k_0 + \bar{p})x} |\eta_{\mu_0}\rangle e^{-iE_{(k_0+\bar{k})\mu_0} t/\hbar} / \sqrt{L}.$$

The energy after the pulse is

$$\varepsilon_{k\mu}^c(t > 0) = \frac{1}{2m^*} \left[ \left( \hbar k + \bar{p} + \mu \frac{p_{so}}{2} \right)^2 - \frac{p_{so}^2}{4} \right].$$

If  $C_{k'\mu'}(k, \mu, t < 0)$  models a Gaussian wave packet centered at  $k_0$  for the injected state we infer from Eq. (8) that right after the pulse the shape of the wave packet is maintained while its central momentum is shifted by  $\bar{p}$ . We conclude: (i) the pulse field delivers a transient momentum transfer (which is proportional to the momentary vector potential) and a net momentum given by the field-amplitude time integrated over the field duration, for harmonic fields this quantity vanishes whereas for HCP it is finite and is equal to  $\bar{p}$ . (ii) This momentum boost is experienced by all the electrons speeding up the device operation. (iii) The phase difference  $\Delta\theta_0$  is maintained as for the static case, i.e., the operation speed is changed by an amount proportional to the field strength while the spin coherence is unchanged, a fact exploitable for the realization of an ultrafast SFET.

#### IV. SECOND DYNAMIC CASE

If the electric field polarization of the HCP is along the  $y$  direction, i.e., perpendicular to the 2DEG the Hamiltonian is

$$H = \frac{\pi^2}{2m^*} + V(\mathbf{r}) - e\Phi + \frac{\tilde{\alpha}(t)}{\hbar} (\boldsymbol{\sigma} \times \boldsymbol{\pi})_{y, k_z=0}, \quad (9)$$

where

$$\tilde{\alpha}(t) = \alpha_R^0 + \alpha_R^t(t)$$

and  $\alpha_R^0$  is the static Rashba SOI,

$$\alpha_R^t(t) = r_R E(t)$$

is proportional to the HCP electric field. The time-dependent part of the Hamiltonian is

$$H^t = \frac{e^2}{2m^*} A_y^2 + \frac{\alpha_R^t(t)}{\hbar} (\boldsymbol{\sigma} \times \mathbf{p})_y. \quad (10)$$

The vector potential in the canonical momentum in the second term in Eq. (10) does not couple to the electric field of the HCP. The total Hamiltonian is  $H = H^0 + H^t$ . The first term in Eq. (10) results in a phase shift for all states in all subbands, an effect which is unimportant for the following discussion and hence we ignore it and consider the total Hamiltonian:

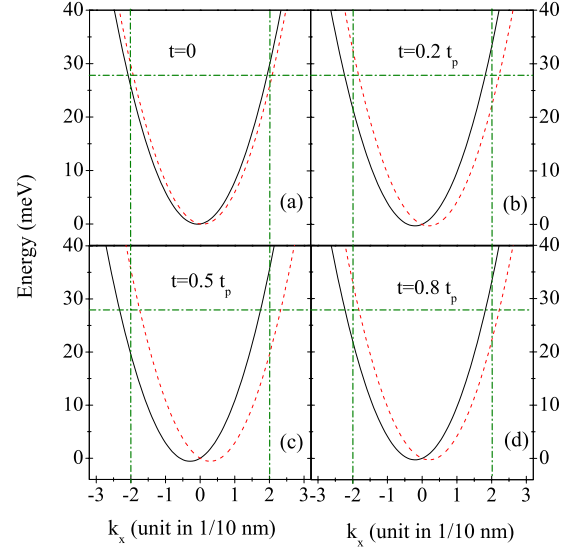


FIG. 2. (Color online) The instantaneous energy spectrum for GaSb/InAs/GaSb system (Ref. 15) at different time fraction of the pulse duration  $t_p$ . In (a) the static spectrum with SOI splitting is shown. The parameters are taken from Ref. 15:  $m^* = 0.055m_0$ ,  $\alpha_R^0 = 0.9 \times 10^{-9}$  eV cm,  $2\alpha_R^0 k_f = 4.0$  meV, and  $k_f = 2.0(10 \text{ nm})^{-1}$ . In (b)–(d) the horizontal dashed-dotted lines indicate the position of the Fermi level for equilibrium system, and the vertical dashed-dotted lines mark the positions of the Fermi wave vectors. The solid (dashed) curves correspond to the spin-up (down) state. The other parameters are selected as  $\frac{E}{\varepsilon_0} = 3$ ,  $a(t) = \sin(\pi t/t_p)$  for  $0 \leq t \leq t_p$ .

$$H = \frac{\hbar^2 k_x^2}{2m^*} + \frac{\tilde{\alpha}(t)}{\hbar} (\boldsymbol{\sigma} \times \mathbf{p})_y. \quad (11)$$

We note, the pulse field does not change the quantum number  $k$  nor the spin states. The time evolution occurs only in the parameter space specified by  $\tilde{\alpha}(t)$  which is 1D parameter space. With varying  $\alpha_R^t(t)$  the spin-dependent potential is changed and so does the energy. To be specific let us consider the GaSb/InAs/GaSb system with the parameters given in Ref. 15. For InAs QW with a width of 6.28 nm the second energy subband is separated from the ground state by  $\approx 700$  meV.<sup>15</sup> To inspect the effect of the pulse field we show in Fig. 2 the instantaneous energy spectrum. As evident from this figure the highest achieved energy level is well below the first excited subband and hence we only need to consider the intraband dynamics in the first subband. The behavior of the instantaneous energy for a state  $|k\mu\tilde{\alpha}(t)\rangle$  is such that for positive  $k$  the energy of the spin-up carrier is raised while the spin-down energy is lowered in the first quarter of the monocycle pulse; the opposite happens in the second quarter cycle. In the second half of the monocycle the spin-resolved energies evolves in an opposite way to the first half. As the shift is determined by the magnitude of the electric field peak, the very weak and long (on the transport time) tail of HCP has a minor effect. From Fig. 2 we infer that the pulse results in a time-varying potential and hence an oscillation of all electrons in the Fermi sphere in energy space. No holes are generated. The spin operators develop as  $\dot{\sigma}_{\pm}(t) = \mp i\tilde{\omega}_{k_x}(t)\sigma_{\pm}(t)$ , where  $\tilde{\omega}_{k_x}(t) = 2\tilde{\alpha}(t)k_x/\hbar$ . Hence, we find

$$\sigma_{\pm}(t) = \sigma_{\pm}(0) \exp \left\{ \pm i \int \tilde{\omega}_{k_x}(t) dt \right\}. \quad (12)$$

Defining  $\tau = -(F/\varepsilon_0)t_p\gamma$ , where  $\gamma = \int a(\xi') d\xi'$  and  $\xi' = t'/t_p$ . We obtain  $\int \tilde{\omega}_{k_x}(t) dt = \omega_{k_x}t + \theta_p$ , where

$$\theta_p = -\omega_{k_x}\tau. \quad (13)$$

For an injected carrier with a spin-polarization vector along  $z$ , the wave function right after the pulse  $\Psi_{k_{\mu_0}\mu_0}(x, t > 0)$  evolves from the stationary state before the pulse  $\Phi_{k_{\mu_0}\mu_0}(x, t < 0)$  as

$$\Psi_{k_{\mu_0}\mu_0}(x, t > 0) = e^{-i\mu_0\theta_p/2} \Phi_{k_{\mu_0}\mu_0}(x, t < 0),$$

meaning that the pulse causes a phase splitting in the spin-up and spin-down channels. The phase difference between up and down spins is equal to the induced angle rotation in Eq. (13).

## V. OPTICAL SFET

In the static case a  $0.67 \mu\text{m}$  long 1D quantum wire is needed to reach the phase shift  $\Delta\theta_0 = \pi$  in 2D  $\text{In}_x\text{Ga}_{1-x}\text{As}$ .<sup>16</sup> While a shorter 1D quantum wire is sufficient to reach  $\pi$  phase shift with length  $0.2 \mu\text{m}$  in GaSb/InAs/GaSb system,<sup>15</sup> since the Rashba SOI is larger, i.e.,  $\alpha_R^0 \approx 0.9 \times 10^{-9} \text{ eV cm}$ . According to Ref. 15 the charge density is  $n = 10^{12} \text{ cm}^{-2}$  and  $2\alpha_R^0 k_f = 4 \text{ meV}$ . We find then  $\omega_{k_x} \approx 2\pi \text{ ps}^{-1}$ . The Fermi velocity is  $0.4 \mu\text{m/ps}$ . Thus in  $0.2 \mu\text{m}$  1D QW there are 20 electrons distributed and the transport time  $t_{tr}$  for the electrons at the Fermi velocity is about 500 fs. In what follows we based our discussion on these realistic numbers. HCPs with peak field of up to several hundreds of kilovolt per centimeter and duration in the picosecond and subpicosecond regimes can be experimentally generated.<sup>10</sup> Novel principles allow the generation of unipolar pulses as short as 0.1 fs and with intensities up to  $10^{16} \text{ W/cm}^2$ .<sup>11</sup> The pulse induced precession angle is

$$\theta_p = k_{so} L_p (F\gamma/\varepsilon_0),$$

where  $L_p$  is the length traveled by a particle with a momentum  $\hbar k_x$  within the pulse duration  $t_p$ . The total precession angle accumulated while traversing the length  $L$  is

$$\Delta\theta = \Delta\theta_0 + \theta_p,$$

where the first term stems from the static Rashba SOI. To evaluate the angle  $\theta_p$  we consider the ratio

$$\lambda = \frac{\theta_p}{\Delta\theta_0} = \frac{L_p F \gamma}{L \varepsilon_0}. \quad (14)$$

(i) *Single HCP.* For  $t_p = 20 \text{ fs}$  we note  $t_p \ll t_{tr}$ , and we have  $\gamma \approx 1$ . The static fields  $\varepsilon_0$  are typically of the order of several kilovolts per centimeter, for example, in GaAs/ $\text{Al}_{0.3}\text{Ga}_{0.7}\text{As}$  quantum well.<sup>14</sup>  $F$  can be generated with several hundreds of kilovolt per centimeter. Therefore,  $F/\varepsilon_0$  can be tuned as high as 100 (without inducing intersubband transitions),

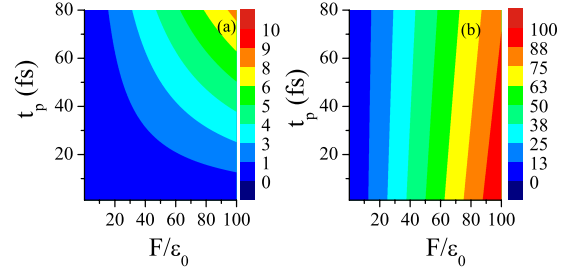


FIG. 3. (Color online) A contour plot for the ratio  $\lambda$ , as given by Eq. (14), as a function of the magnitude and of the duration of the external field (a) for a single HCP and (b) for a train of HCPs. In (a), the Fermi velocity is chosen as  $0.4 \mu\text{m/ps}$ ,  $L = 0.2 \mu\text{m}$ ,  $\gamma$  is set at 0.5. In (b),  $t_s = 0.5 \text{ ps}$ ,  $t_{tr} = 0.5 \text{ ps}$ , and  $\gamma \approx 1$ .

$$L_p = \hbar k_x t_p / m^*$$

is about  $8 \text{ nm}$  during  $t_p = 20 \text{ fs}$ . Therefore,  $\lambda \approx 4$ . The induced accumulated angle  $\theta_p$  is a swift procession angle transfer. If the initial injected spin is in  $x$  direction it suddenly rotates anticlockwise over an angle  $\theta_p$  upon applying the pulse. This conclusion is exploitable to realize a nanosize SFET. Spins with higher drift velocity experience a larger angular transfer during the period  $t_p$ . Thus, a wave packet of spins that is initially polarized in the same direction but contain different velocity components will spread over a range of angles after scattering from the pulses. This can be compensated by operating the device in the linear-response regime (small bias) where electrons at the Fermi surface dominate the transport, i.e.,  $k_x = k_f$ . A contour plot of the ratio  $\lambda$ , as introduced in Eq. (14), as function of the external field parameters is shown in Fig. 3(a). The ratio increases with increasing  $F$  and  $t_p$ .

(ii) *A train of HCPs.* We apply a pulse train with a total duration  $t_t = n(t_p + t_s)$ , where  $n$  is the number of HCP peaks and  $t_s$  is the time interval between two consecutive HCPs.  $m = t_t/t_{tr}$  is the number of electrons passing through the device during  $t_t$ . We define for a single electron an averaged  $\bar{\gamma} = n\gamma/m$  and  $L_p = L$ . For  $t_p = 1 \text{ ps}$ ,  $t_s = 0.5 \text{ ps}$ ,  $t_t = 125 \text{ ps}$ ,  $n = 50$ ,  $m = 250$ , and  $\gamma \approx 1$ , thus  $\bar{\gamma} \approx 0.2$  and  $\lambda \approx 20$ . In this case, spins with higher velocities will pass through the device faster while rotating faster. The additional precession angle induced by the pulse field is comparable to the static one or may even be larger. To show the behavior more clearly, a contour plot of  $\lambda$  with the tuning parameters of the external field is given in Fig. 3(b). Since the transport time is fixed, increasing  $t_p$  decreases the ratio  $\lambda$  which is in a contrast to the single HCP case. However, a larger  $\lambda$  is also achieved via increasing the electric field strength. As discussed for the static case, an injected spin polarized in positive or negative  $z$  direction experiences a phase shift, while passing through the length  $L$ , that is exactly equal to the precession angles for the case of a spin injected polarized along  $x$  or  $y$ .

## VI. EXPERIMENTAL REALIZATION

To be specific we propose a GaSb/InAs/GaSb structure (as realized in Ref. 15) with InAs QW width of  $6.28 \text{ nm}$  and a length of  $\approx 100 \text{ nm}$  driven by pulses with a strength of 10



kV/cm and duration of 1 ps. The nonzero  $k_z = \pi/w$  leads to a spin-flip transition between subbands with a magnitude linear in  $k_z$  while the energies for subbands are proportional to  $k_z^2$ . In this configuration our model is valid since  $w/L \ll 1$ . For smaller  $L$ , we reduce  $w$ . The spin precession is however controlled by the pulses and hence our SFET is still operational. For  $L$  is less than the mean-free-path, dephasing caused by impurities is negligible. Scattering from (acoustic) phonons is spin independent causing a current relaxation within tens of picosecond<sup>17</sup> at few kelvins which is larger than our transport time. Since the pulses are weak the second transverse subband is not reached, hence intersubband transitions play no role. Also, multiphoton processes are subsidiary for such weak pulses. Finally we remark, symmetric pulses do not lead to a net currents ( $\gamma = 0$ ) whereas a similar effect is achievable via quantum interferences in (the higher energy) interband, one-photon, and two-photon absorption.<sup>18</sup> Our photoinduced current for the first dynamic is sizable. For the above device based on GaSb/InAs/GaSb the static current is usually  $\approx 1.07 \mu\text{A}$ . The ratio between the induced and the static current is  $\Delta = 2\bar{p}/\hbar k_f$  which varies in a large range. For  $F = 10 \text{ kV/cm}$  and  $t_p = 1 \text{ ps}$  we find  $\Delta = 7.6$  which means an induced current of  $\approx 8 \mu\text{A}$ .

## VII. CONCLUSIONS

In summary we studied theoretically the photoinduced spin dynamics in a quantum wire with a Rashba spin-orbit interaction. For an efficient and a subpicosecond control of the spin dynamics the pulses have to be shaped appropriately such that in effect a linear momentum boost is transferred to the charge carrier. For linear polarized photons we find that if the photon polarization axis is along the wire's direction, the phase coherence in different spin channels is maintained, even though the charge carriers are speeded up. In the case that the photon pulse polarization is perpendicular to the wire we predict an additional spin precession comparable to that induced by the static Rashba SOI. The photon-induced precession is tunable in a considerable range by scanning the parameters of the pulse electric field.

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