Solid-State Spin-Photon Quantum Interface without Spin-Orbit Coupling

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We show that coherent optical manipulation of a single confined spin is possible even in the absence of spin-orbit coupling. To this end, we consider the non-Markovian dynamics of a single valence orbital hole spin that has optically induced spin-exchange coupling to a low-temperature partially polarized electron gas. We show that the fermionic nature of the reservoir induces a coherent component to the hole spin dynamics that does not generate entanglement with the reservoir modes. We analyze in detail the competition of this reservoir-assisted coherent contribution with dissipative components displaying markedly different behavior at different time scales and determine the fidelity of optically controlled spin rotations.

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Quantum dot (QD) spins have emerged as a new paradigm for studying quantum optical phenomena in the solid-state. Motivated by potential applications in quantum information processing, research in this field has focused on understanding spin decoherence induced by the solid-state environment and implementing coherent spin manipulation [1]. While spin does not directly couple to electric fields, electrostatic quantum control has been achieved by employing spin-orbit coupling [2,3] or nuclear field gradients [4–6]. Schemes for *optical* manipulation of solid-state spins considered to date on the other hand rely exclusively on spin-orbit interaction in either the initial or final state of the optical transition [7–12].

In this Letter, we show that the non-Markovian dynamics of spin-exchange coupling between a confined spin and a fermionic reservoir (FR) enables realization of a coherent spin-photon interface. There are two key results conveyed by our work: first, we demonstrate that, contrary to the common conception of treating reservoirs as sources of decoherence, an engineered fermionic reservoir gives rise to a coherent contribution to single QD spin dynamics that could dominate over the decoherence it induces. Second, coherent optical manipulation of a single confined spin is possible in systems having weak or no spin-orbit interaction. In addition to being interesting from a basic optical physics perspective, this latter result implies that single confined spins in emerging material systems such as graphene where the spin-orbit effects are anticipated to be vanishingly small, could be manipulated using optical fields.

We consider here the fidelity of FR-assisted Raman transitions of a generic qubit encoded in the spin state ($|\Uparrow\rangle$, $|\Downarrow\rangle$) of a single QD valence-band (VB) hole [13] (see Fig. 1). We choose units with k_B, $\hbar = 1$. Starting with a neutral QD in equilibrium, initialization may be performed by driving a strongly detuned optical transition to neutral excitonic (X^0) states $|\uparrow\Downarrow\rangle$, $|\downarrow\Uparrow\rangle$ such that the effective final states of the transition coincide with the Fermi edge ε_F. A finite magnetic field along the QD spin quantization axis imparts a Zeeman splitting $\Delta_z \gg T$ on the QD, such that the optically excited QD conduction band (CB) electron in spin state lying $\Delta_z/2$ above ε_F tunnels into the Fermi sea, trapping a hole of opposite spin in the VB. This state is metastable for vanishing tunnel matrix elements between QD VB and FR.

We now discuss interfacing the trapped hole spin with optical fields. Two lasers (with Rabi frequencies $\Omega_\uparrow$, $\Omega_\downarrow$ and detunings $\delta_\uparrow$, $\delta_\downarrow$) couple the two hole spin states virtually to intermediate charged exciton (trion, X^+) states ($|\uparrow\Uparrow\Downarrow\rangle$, $|\downarrow\Uparrow\Downarrow\rangle$) involving two holes in a singlet state and a conduction band (CB) electron whose spin is determined directly by the spin of the hole in the initial state [13]. The X^+ eigenstates lie sufficiently far below the Fermi edge such that it is energetically unfavorable for a QD electron to tunnel out or an extra electron to tunnel in. Employing a Schrieffer-Wolff transformation [14] to eliminate tunnel coupling to lowest order, and in the limit of $\delta_\uparrow$, $\delta_\downarrow \gg \Omega_\uparrow$, $\Omega_\downarrow$, the relevant physics may be described by an effective low-energy Hamiltonian that couples effective (dressed) hole spin states $\hat{\tilde{\sigma}}$ to spin scattering excitations in the FR, where two sequential virtual tunnel processes potentially

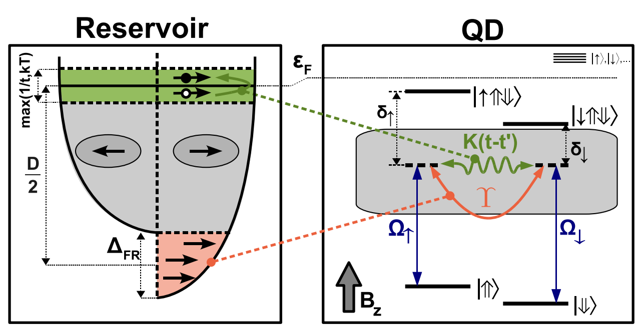

FIG. 1 (color online). Reservoir-assisted spin-photon interface. QD: Lasers (Rabi frequency $\Omega_\downarrow$, $\Omega_\uparrow$) drive strongly detuned ($\delta_\uparrow$, $\delta_\downarrow$) spin-conserving dipole transitions ($|\Downarrow\rangle \leftrightarrow |\downarrow\Uparrow\Downarrow\rangle$, $|\Uparrow\rangle \leftrightarrow |\uparrow\Uparrow\Downarrow\rangle$) between valence and excitonic spin states. Virtually excited states exhibit cotunneling (dashed arrows) with a fermionic reservoir that induces both coherent (Υ) and dissipative non-Markovian [$K(t - t')$] coupling of the two spin transitions.

create an electron-hole excitation in the reservoir. The final effective dynamics are governed by spin-exchange Hamiltonian $H_{\rm sd} = \sum_{kk'} J_{kk'} \hat{\vec{\sigma}} \cdot \hat{\vec{S}}_{kk'}$, a spin-conserving scattering term $H_{\rm dir} = \sum_{kk'} \frac{W_{kk'}}{2} (\hat{c}^\dagger_{k,\uparrow} \hat{c}_{k',\uparrow} + \hat{c}^\dagger_{k,\downarrow} \hat{c}_{k',\downarrow})$ and an exchange interaction-mediated energy shift of the effective hole spin states $H'_0 = \sum_{kk'} \frac{J_{kk'} + W_{kk'}}{2}$. Here $\hat{\vec{S}}_{kk'} = \sum_{ss'} \hat{c}^\dagger_{ks} (\frac{1}{2} \vec{\tau}_{ss'}) \hat{c}_{k's'}$ ($\vec{\tau}$ are Pauli matrices, $\hat{c}^\dagger_{ks}/\hat{c}_{ks}$ fermionic operators for the FR electrons with spin $s = \pm$). The optically induced exchange coupling to lowest order in $\Omega_\sigma(t)/\delta_\sigma$ reads

$$J_{kk'} \approx \frac{\Omega_\uparrow(t)\Omega_\downarrow(t) V_k V_{k'}}{4\delta_\uparrow \delta_\downarrow} \left[\frac{1}{\Delta_d - \varepsilon_k} - \frac{1}{\Delta_0 - \varepsilon_k} \right] + [\mathbf{k} \leftrightarrow \mathbf{k}'] \tag{1}$$

and the corresponding directional coupling is $W_{kk'} = J_{kk'} + \Omega_\uparrow(t)\Omega_\downarrow(t) V_k V_{k'} / [2\delta_\uparrow \delta_\downarrow (\Delta_0 - \varepsilon_k)] + [\mathbf{k} \leftrightarrow \mathbf{k}']$. Here, we assume two-photon resonance condition between the hole spin states, namely $\delta_\uparrow - \delta_\downarrow = \Delta_e$, where Δ_e is the QD CB Zeeman splitting. V_k is the tunnel coupling between QD and FR, $\Delta_d \approx \varepsilon_c + U_{ee} - 2U_{eh} + \delta_L$ and $\Delta_0 \approx \varepsilon_c - 2U_{eh} - \delta_L$ are the energies required to put an extra electron into or remove an electron from the QD; here, $\delta_L = \frac{\delta_\uparrow + \delta_\downarrow}{2}$. U_{ee}, U_{eh} model intradot electron-electron and electron-hole Coulomb interactions in Hartree-Fock approximation, and ε_c is the bare QD CB electron energy.

We consider a two-dimensional electron gas (2DEG) FR. To implement coherent spin rotation, we require that the FR spins are partially polarized along an axis that is tilted by an angle θ with respect to the QD spin axis $\mathbf{z}$. We consider imparting a finite spin splitting $\Delta_{\rm FR}$ on the 2DEG via tunnel-coupling to ferromagnetic gates (yielding realistic $\Delta_{\rm FR} \sim 100\ \mu$eV) [15,16]. Alternatively, spin-3/2 heavy holes in zinc-blende semiconductor QDs show strong anisotropy of the g factor, which may be exploited by an external magnetic field. However, this latter effect intrinsically relies on spin-orbit interaction.

The FR Hamiltonian is given by $H_{\rm FR} = \sum_{ks} \varepsilon_{ks} \hat{c}^\dagger_{ks} \hat{c}_{ks}$ with $\varepsilon_{ks} = \varepsilon_k + \frac{s}{2} \Delta_{\rm FR}$, where spin s denotes a FR electron spin along angle θ. We introduce an effective finite bandwidth $D = 1/(2\rho)$ for FR electrons. Here, ρ is the 2D density of states, and the band is symmetric around the Fermi edge $\varepsilon_F \equiv 0$. We assume weak spin polarization $\Delta_{\rm FR} \ll D$. Within the second order Born approximation in $J_{kk'}$, we obtain a generalized master equation for the dynamics of the QD hole spin, written in the form of Bloch equations:

$$\frac{d}{dt}\sigma^z = -8 \int_0^t dt' K(t - t') \sigma^z(t') + \sin(\theta) \Upsilon \sigma^x(t), \tag{2}$$

$$\frac{d}{dt}\sigma^y = -8 \int_0^t dt' K(t - t') \sigma^y(t') - \cos(\theta) \Upsilon \sigma^x(t), \tag{3}$$

$$\frac{d}{dt}\sigma^x = -8 \int_0^t dt' K(t - t') \sigma^x(t') - \sin(\theta) \Upsilon \sigma^z(t) + \cos(\theta) \Upsilon \sigma^y(t). \tag{4}$$

Here, σ^i are the expectation values of the QD spin (axis $\mathbf{z}$) Pauli operators. Neglecting a spin-independent exchange interaction-mediated energy shift, the reduced QD spin dynamics for $\Delta_{\rm FR}, T \ll D$ described by Eqs. (2)–(4), comprises a coherent spin precession quantified by the rate $\Upsilon = \sum_k (J_{kk} + W_{kk}) \langle \hat{S}^z_{kk} \rangle$ and dissipative dynamics captured by the memory kernel $K(t) = {\rm Re} \sum_{kk'} |J_{kk'}|^2 \langle \hat{S}^z_{k'k}(t) \hat{S}^z_{kk'}(0) \rangle$ [17]. In the following, we will discuss the case of orthogonal FR polarization ($\theta = \pi/2$) [18].

Dissipative and coherent contributions are mediated by distinct parts of the FR. First, the action of the strip of width $\Delta_{\rm FR}$ of excess spins at the lower band edge may be viewed as generating a net magnetic moment of the 2DEG around which the QD spin precesses with rate [19]

$$\Upsilon \approx \tfrac{1}{2} \rho [J_{-(D/2),-(D/2)} + W_{-(D/2),-(D/2)}] \cdot \Delta_{\rm FR}. \tag{5}$$

This action is mediated via energy-conserving cotunnelling processes originating from this strip of electrons through an intermediate doubly occupied QD. This process is coherent in the sense of keeping the QD spin states disentangled from the reservoir states of the polarized strip. The lack of entanglement is ensured by the Pauli exclusion; the absence of free spin states in the polarized strip prohibits modification of the reservoir. This coherent contribution to the system dynamics is valid even in the limit of a single polarized spin.

The dissipative part of the dynamics is mediated by electron-hole scattering processes within a strip of *dynamical* width $\sim \max(1/t, T)$ around the Fermi energy. We focus essentially on the dissipative limit of Kondo effect, in which the QD spin entangles with electrons within the dynamical strip described above and tries to evolve to a singlet with FR electrons but the coupling is too weak to maintain correlations, moving the system towards a maximum uncertainty state $\varrho = (|\Uparrow\rangle\langle\Uparrow| + |\Downarrow\rangle\langle\Downarrow|)/2$. The validity of this description is bounded by time scale $1/T_K$ where $T_K = D\sqrt{J\rho}e^{-1/2J\rho}$ is the Kondo temperature; this bound is effective only for $T < T_K$ [20–22].

We consider the finite-bandwidth memory kernel $K(t) = \int_{-D}^{D} d\omega \cos(\omega t) \mathcal{S}(\omega)$, where $\mathcal{S}(\omega)$ is the FR spectral function, identical in form to a finite-bandwidth spin-boson model [23,24]:

$$\mathcal{S}(\omega) = (J\rho)^2 [1 + n_B(\omega)] T \log\left[\frac{\cosh(\frac{\omega - |\omega| + D}{4T})}{\cosh(\frac{\omega + |\omega| - D}{4T})} \right]. \tag{6}$$

Here, $n_B(\omega) = 1/(e^{\omega/T} - 1)$. As the X^+ state lies well below the Fermi edge, $J_{kk'}$ will be flat around the Fermi edge and can be approximated by $J \equiv J_{k_F k_F}$, as is standard in Kondo physics [20]. Microscopically, ω is the energy of electron-hole excitations in the reservoir. A small-frequency expansion of (6) gives $\mathcal{S}(\omega) = \frac{1}{4} (J\rho)^2 [2kT + \omega + \mathcal{O}(\omega^2/T)]$. The corresponding memory kernel expansion is $K(\tau) = \gamma_M \cdot \delta(\tau) + (J\rho)^2/4\tau^2$, with $\gamma_M = \frac{\pi}{2} (J\rho)^2 kT$ yielding the decay rate within Markov approxi-

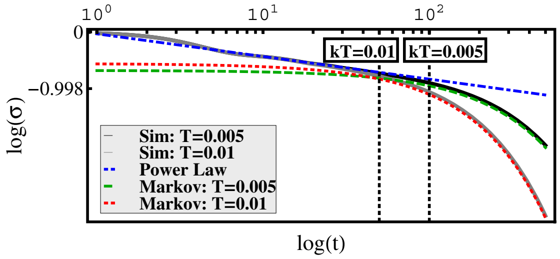

FIG. 2 (color online). Log-log-plot of $\sigma^z(t)$, using $J\rho = 0.005$. Units scaled with respect to D. Thick solid lines show simulation results for $T = 0.005D$ (black) and $T = 0.01D$ (gray), dashed lines are intermediate- and long-time asymptotics. Vertical lines denote expected crossovers at $1/2T$ between non-Markovian and Markovian regimes. $\sigma^z(t)$ displays power-law relaxation for $t < 1/2T$, and long-time exponential decay for $t > 1/2T$.

mation, $\langle\hat{\sigma}^i(t)\rangle \sim e^{-\gamma_M t}$ [25]. This exponential decay is the well-known Korringa relaxation [26] in the high-T limit. We find from inspection of the expansion that exponential decay is only valid for $t \gg 1/2T$. We emphasize here that $\Upsilon > \gamma_M$ is attainable even when $\Delta_{\rm FR} < kT$ since $\Upsilon \propto (J\rho)\Delta_{\rm FR}$ whereas $\gamma_M \propto (J\rho)^2 kT$.

To obtain an analytic expression for $\sigma^i(t)$ in the (non-Markovian) intermediate time scale ($1/D \ll t \ll 1/T$), we set $\Upsilon = 0$ and use a long-time asymptotic expansion of the zero-temperature limit of the memory kernel

$$K(\tau) = \frac{1}{8}(J\rho)^2 D^2 \cos\left(\frac{D\tau}{2}\right)\mathrm{sinc}^2\left(\frac{D\tau}{4}\right) \tag{7}$$

[with $\mathrm{sinc}(x) = \sin(x)/x$]. This form of $K(\tau)$ indicates initial-time oscillations on an ultrashort and experimentally irrelevant time scale of $1/D$. An analytic expression for $t \gg 1/D$ can be obtained by considering the analytic properties of the Laplace transform of the spin dynamics $\sigma(s) = s^{-1}[1 + 8\mathcal{L}\{K(\tau)\}/s]^{-1}$ around $s = 0$:

$$\sigma(s \to 0) \approx \left(\frac{\tilde{J}}{J}\right)^2 \frac{1}{s}\frac{1}{1 - 4(\tilde{J}\rho)^2 \cdot \log(\frac{4s}{D})}, \tag{8}$$

where $\tilde{J} = J/\sqrt{1 + 4(J\rho)^2}$. $\sigma(s)$ has two branch points at $s = \{0, \infty\}$ and we choose the branch cut along the negative real axis. The long-time asymptotic behavior of $\sigma(t)$ is characterized by the functional form of $\sigma(s \to 0)$ along a keyhole contour around $s = 0$ with a radius ϵ in the complex plane and can be obtained by the power-law expansion of the denominator using recursive relationship $\xi \log(s) = 1 - s^{-\xi} + \sum_{n=2}^{\infty} \frac{[-\xi \log(s)]^n}{n!}$ converging in powers of $\xi = 4(\tilde{J}\rho)^2$ by choice of ϵ, which yields to lowest order in ξ and in time domain:

$$\sigma\left(t \ll \frac{1}{T}\right) = \left(\frac{\tilde{J}}{J}\right)^2 \frac{1}{\Gamma(1 - 4(\tilde{J}\rho)^2)}\left(\frac{Dt}{4}\right)^{-4(\tilde{J}\rho)^2} \tag{9}$$

with $\Gamma(x)$ the Gamma function (see Fig. 2). The effect of initial power-law decay for $t \ll 1/T$ on the long-time dynamics $t \gg 1/T$ may be extracted by approximating $\mathcal{S}(\omega)$ by a piecewise linear function $\tilde{\mathcal{S}}(\omega < 2T) = T$, $\tilde{\mathcal{S}}(\omega \geq 2T) = \lim_{T\to 0}\mathcal{S}(\omega)$, from which we obtain long-time dynamics $\sigma(t \gg \frac{1}{T}) = (\tilde{\gamma}_M/\gamma_M)e^{-8\tilde{\gamma}_M t}$ with $\tilde{\gamma}_M = \gamma_M/[1 - 4(J\rho)^2 \log(8T/D)]$ the modified relaxation rate.

We now discuss the quality of optical spin manipulation via Raman transition. Fidelity of spin manipulation is commonly characterized in terms of π/n pulses (i.e., π/n rotations of the QD spin). We focus on performing a $\pi/2$ rotation—more precisely, rotating an initialized spin $|\Uparrow\rangle$ to symmetric superposition $\frac{1}{\sqrt{2}}(|\Uparrow\rangle + |\Downarrow\rangle)$. A measure of fidelity is $F_{\pi/2} = \sqrt{(1 + \langle\sigma^x(\tau_{\pi/2})\rangle)/2}$ [27] with $\tau_{\pi/2}$ the optimal time for a $\pi/2$ rotation. We state fidelity for rectangular pulse shapes and in the limit of separability of dissipative and coherent dynamics. This decoupling is justified in the context of analyzing a π/n pulse by noting that while the reservoir-mediated coherent drive dresses the QD effective hole spin states with splitting $\sim\Upsilon$, π/n rotations probe only time scales $\lesssim 1/\Upsilon$ [and hence do not see the modification of $\mathcal{S}(\omega)$ for energies $\lesssim \Upsilon$]. This approximation breaks down for longer pulse durations. In a non-Markovian regime of operation $\pi/(2\Upsilon) \ll \min(1/T, 1/T_K)$, the fidelity is given by

$$F_{\pi/2} = \sqrt{\frac{1}{2} + \frac{\tilde{J}^2}{2J^2\Gamma(1 - 4(\tilde{J}\rho)^2)}\left(\frac{\pi D}{8\Upsilon}\right)^{-4(\tilde{J}\rho)^2}}. \tag{10}$$

In the Markov regime $1/T \ll \pi/(2\Upsilon)$, it reads

$$F_{\pi/2}^{\rm Markov} = \sqrt{\frac{1}{2} + \frac{\tilde{\gamma}_M}{2\gamma_M}e^{-4\pi\tilde{\gamma}_M/\Upsilon}}. \tag{11}$$

We discuss fidelity in terms of effective exchange coupling $\Lambda = \pi\rho\Omega_\uparrow\Omega_\downarrow V^2/4\delta_\uparrow\delta_\downarrow$ (for constant $V_k = V$). Figure 3 plots the fidelity of a $\pi/2$ pulse across different parameter regimes. We note that while the Markovian result would hint at vanishing dissipation with decreasing temperature, the system in fact approaches a fundamental limit of dissipation with power-law decay. Nevertheless, our results show that arbitrarily high-fidelity coherent spin manipulation is possible either by reducing T for a fixed Λ into the non-Markovian regime, or by reducing Λ for a fixed T and moving the system into the Markovian limit.

We have presented a scheme for coherent optical manipulation of a single confined spin mediated through tunnel coupling to an electron gas. We find that reservoir engineering of such an electron gas, achieved in our case by weak spin polarization, can modify the nature of system reservoir coupling drastically, exploiting the fermionic nature of the reservoir and turning it into a resource that allows for coherent spin manipulation. Pauli exclusion ensures that spin-exchange processes involving electrons originating from low-lying polarized spin states cannot write off information about the hole spin on the reservoir. We show that the microscopic description of such a reservoir at low temperatures gives rise to non-Markovian

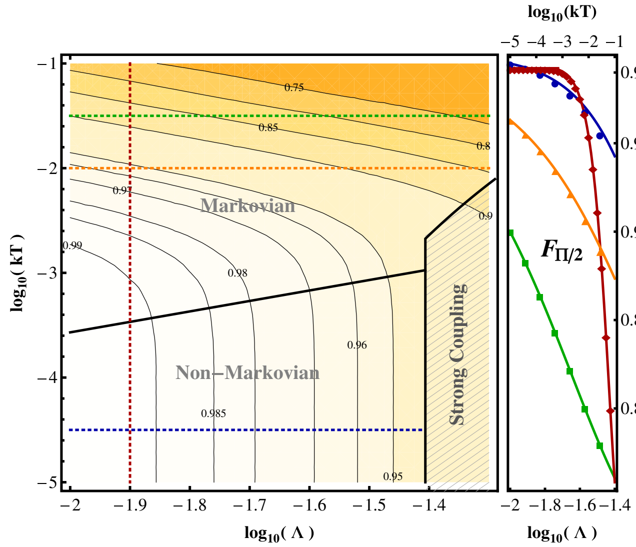

FIG. 3 (color online). Left: Contour plot of fidelity $F_{\pi/2}$ in Markovian and non-Markovian regimes. Each pixel expresses the fidelity of the final state of a $\pi/2$ pulse after time $\tau_{\pi/2}(\Lambda) = \pi/2\Upsilon(\Lambda)$, given parameters T (temperature) and Λ (exchange coupling). Symmetric model with $-\Delta_0 = \Delta_d = 0.5D$ and $\Delta_{FR} = 0.1D$, giving $J = 8\Lambda/\pi D$ and $\Upsilon = 8\Lambda\Delta_{FR}/3\pi D$. All energies are scaled by bandwidth D. Black lines display $\tau_{\pi/2} = 1/2T$ (horizontal line), $\tau_{\pi/2} = 1/T_K$ (vertical line) and $1/T = 1/T_K$. Simulation data [starting from (2)–(6)] is sampled with steps $0.25T$, 0.1Λ. The crossover to non-Markovian relaxation is clearly depicted by the "bending down" of fidelity contour lines, conveying that low-temperature non-Markovian spin manipulation is solely limited by the strength of exchange coupling Λ, manipulation of which displays competition of Λ^2-dependent power-law decay and Λ-dependent coherent spin rotation. Right: Select contour plot cuts with fixed Λ or T (colored dotted lines in both graphs) are compared to analytic results.

dynamics, offering an attractive regime of attaining high-fidelity spin manipulation. On a fundamental level, our findings demonstrate that contrary to the common wisdom, spin-orbit interaction is not necessary for realizing a spin-photon interface.

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