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Entanglement-buildup through charged-exciton decay in a semiconductor quantum dot

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We analyze the decay of a single-electron charged exciton in a quantum dot embedded in a field effect structure. We show how the quantum properties of the charged exciton are transferred through tunneling and relaxation to the spin entanglement between electrons in the dot and contact, carefully examine the proper theoretical description of the underlying scattering processes, and identify the pertinent disentanglement mechanisms.

1 Introduction Ultrafast semiconductor spectroscopy provides a unique laboratory for the investigation of coherence and scattering effects in solids [1]. Recent years have seen spectacular examples, e.g. the observation of phonon quantum beats [1] or the buildup of screening [2], and have revealed dephasing and relaxation times of the order of femto- and picoseconds for carrier-carrier and carrier-phonon interactions, respectively. In semiconductor quantum dots, sometimes referred to as *artificial atoms*, the strong confinement in all three spatial directions results in a substantial suppression of decoherence: on the one hand, Coulomb interactions among carriers do not result in scattering but only give rise to energy renormalizations of the electron–hole few-particle states [3]; on the other hand, at low temperatures phonon-mediated dephasing of the electron–hole states is of only minor importance. These remarkable features render quantum dots ideal candidates for the solid-state implementation of quantum engineering, [4] e.g., for the purpose of quantum computation or quantum communication.

Such future technology requires devices operating at the single-quantum level. Prominent examples are the quantum-dot based single-photon [5] and single-electron [6] turnstile devices. In both cases a single quantum dot is optically excited and consequently decays through emission of single quanta, i.e. photons or electron–hole pairs, which is monitored by photon detection or as a photo current. Furthermore, when the excited state can decay to two, usually spin-degenerate states of lower energy, it becomes possible to establish entanglement through the cascade process. This has been utilized for the proposals of turnstile entangled-photon [7] and entangled-electron [8] devices. In this paper, we first discuss the latter device recently proposed by ourselves. We develop a density-operator approach of Lindblad form for its theoretical description, which allows for a convenient solution of the complete cascade decay process through unraveling. In addition to our original paper [8], we carefully examine the underlying scattering processes and discuss the foundations and limitations of our theoretical approach.

The proposed structure (Fig. 1a) is similar to the one used by Zrenner et al. [6] it consists of a single electron-doped quantum dot inside a field-effect structure; such doping can be achieved by applying an external bias voltage which transfers an electron from a nearby *n*-type reservoir to the dot

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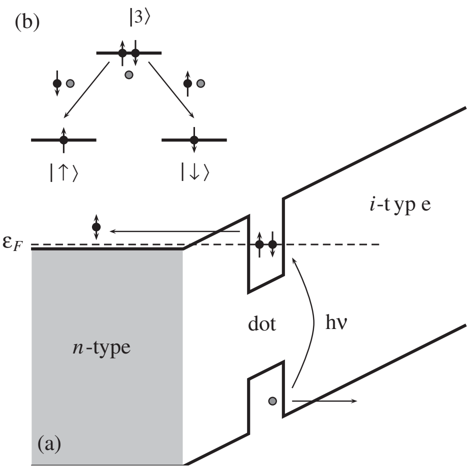

Fig. 1 a) Schematic band diagram of the proposed structure. b) Level scheme of the spin-degenerate electron states $|\sigma=\uparrow,\downarrow\rangle$ and the charged-exciton state $|3\rangle$ in the dot.

[3], where further charging is prohibited because of the Coulomb blockade. Optical excitation of this structure then results in the excitation of a *charged exciton*, i.e., a complex consisting of two electrons and a single hole [3]. Since within the field-effect structure the charged exciton is not a stable configuration, in a consequent step one electron and hole will tunnel out from the dot to the nearby contacts; here, the system can follow two pathways, where either the electron in the dot has spin-up and the one in the reservoir spin-down orientation or vice versa. According to the laws of quantum mechanics, the total state of the system thus becomes a superposition of these two configurations where the spins of the dot and reservoir electron are maximally entangled.

2 Master equation of Lindblad form In our theoretical approach the dot is characterized by (Fig. 1b) the spin-degenerate electron ground states σ and the negatively charged exciton state $|3\rangle$, i.e., a Coulomb renormalized complex consisting of two electrons and a hole. As for the reservoirs we assume fermionic field operators $C_\sigma^\dagger(\omega)$ and $D^\dagger(\omega)$ [8], which, respectively, describe the creation of an electron (spin orientation σ) or hole with energy ω. Since we are dealing with an *open system* [9], i.e., a system interacting with its environment, we have to adopt a density-operator formalism. With $H_0 = H_0^d + H_0^R$ the Hamiltonian for the decoupled dot and reservoir states and V the tunnel coupling [8], the time evolution of the density operator ρ is governed by [9]

$$\dot{\rho}_t^I \cong -\int_0^t dt' \, [V^I(t), [V^I(t'), \rho_{t'}]] \,, \tag{1}$$

where the superscript I denotes the interaction representation according to H_0. Quite generally, it requires the Markov and adiabatic approximations [1] to bring Eq. (1) to Lindblad form. We shall discuss both approximations in slightly more detail.

2.1 Markov approximation Within the Markov approximation, in Eq. (1) $\rho_{t'}$ is replaced by ρ_t. Thus, the time evolution of the density operator at a given time is completely determined by ρ_t itself; as discussed in detail in Ref. [10], this approximation is appropriate for systems where higher-order contributions to Eq. (1) are negligible, or in other words, where scatterings occur seldomly on a time scale given by H_0.

2.2 Adiabatic approximation The adiabatic approximation, which replaces the memory kernel in Eq. (1) by a delta-like memory, is more subtle. Similarly to Fermi's golden rule, this procedure gives generalized scattering rates between asymptotic states. However, conceptual difficulties arise in case of final-state ambiguities, such as for the level scheme of Fig. 1b, where the charged exciton $|3\rangle$ can decay to either $|\downarrow\rangle$ or $|\uparrow\rangle$. What is the final state of such a decay? To establish entanglement, a quantum-mechanical superposition of the two decay paths must be formed; on the other hand, in a simple-minded framework of Fermi's golden rule the system would have to relax through just one of the possible decay channels and would thus not allow the entanglement buildup.[1]

We will now show how to overcome this problem. First, through the tunnel coupling V a quantum coherence between dot and reservoir is established. This initial buildup of quantum coherence, which, according to Zurek [11], may be called a *premeasurement*, does not lead to decoherence. However, the tunnel-generated electron–hole pair does not propagate freely but is subject to environment interactions, e.g., inelastic phonon scatterings of the hole which enters with a high excess energy into the contact. Here, the situation is different with respect to the tunnel-induced quantum-coherence buildup because of our ignorance about the details of the environment state. This ignorance is usually accounted for by tracing over the environmental degrees of freedom [9, 11], and leads to a wavefunction collapse. According to Zurek, it is precisely this triangle dot–reservoir–environment and our ignorance about the detailed environment state before and after scattering that introduces decoherence. We can account for this ignorance by tracing in Eq. (1) over degrees of freedom which we are not interested in, e.g., those of the hole. The remainder of our procedure is then more technical and quite standard. It is based on the observation that the V^I's evolve with different frequency components. Thus, when performing the partial trace we encounter a short-lived memory kernel, which, under appropriate conditions and in the spirit of the random-phase approximation, can be approximated as delta-like.[2] Apparently, this procedure does not invoke any conceptual difficulties regarding the proper choice of final states and validates the use of the adiabatic approximation for the study of our present concern.

Before proceeding, let us briefly comment on the consequences of our analysis for the entangled-photon creation through biexciton decay [7]. The corresponding level scheme corresponds to that of Fig. 1b, where $|3\rangle$ is replaced by the biexciton state and $|\sigma\rangle$ by the spin-degenerate single-exciton ones. Again we are faced with a final-state ambiguity, which can be used to establish a photon-dot entanglement. However, in contrast to the tunnel-generated electron–hole pair the photon propagates without further environment interactions, thus, in principle, not allowing the aforementioned tracing procedure. On the other hand, if we detect the emitted photon through broadband photocounting, as usually done in such experiments, we keep ignorant about the details of the photon state. Theoretically this is accomplished by tracing over the unknown photon degrees of freedom, which again validates the use of the adiabatic approximation [12].

3 Unraveling of the cascade decay The charged-exciton decay can thus be described by a master-equation approach of Lindblad form [8, 9] $\dot{\rho} \cong -i(H_{\mathrm{eff}}\rho - \rho H_{\mathrm{eff}}^\dagger) + \sum_i L_i \rho L_i^\dagger$, with $H_{\mathrm{eff}} = H_o - \frac{i}{2}\sum_i L_i^\dagger L_i$ and L_i the Lindblad operators accounting for the different scattering channels, i.e., tunneling and scatterings in the reservoirs (see Ref. [8] for details). Conveniently it is solved through unraveling, and we obtain [8, 9, 12]

$$\begin{aligned}\rho(t) \cong p^{(0)}(t)\,|3\rangle\,\langle 3| + \mathrm{const} \times \sum_{\sigma\sigma'} \int_{\epsilon_F}^{\omega_c} \mathrm{d}\omega\,\mathrm{d}\omega'\, C_\sigma^\dagger(\omega)\,|\bar{\sigma}\rangle\,\langle\bar{\sigma}'|\,C_\sigma(\omega') \\ \times\,(p^{(1)}(t) + p^{(2)}(t)\,\delta(\omega-\omega') + p^{(3)}(t)\,\delta(\omega-\omega')\,\delta_{\sigma\sigma'})\,. \qquad (2)\end{aligned}$$

Here, $p^{(i)}(t)$ account for the probabilities that the system: (0) has remained unscattered; (1) one electron and hole have tunneled out from the dot, and that the electron in the reservoir has suffered a (2) spin-unselective, and finally (3) spin-selective scattering. Most importantly, we observe in Eq. (2) that spin-unselective scatterings destroy the phase coherence (i.e., dephasing) but do not affect the degree of spin entanglement. *Thus, the decay of an optically excited charged-exciton indeed generates a*

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[1] Note that the situation is particularly cumbersome when the two spin states are not perfectly degenerate, e.g., due to magnetic fields or to band mixing. Here, the final states not only differ with respect to spin orientation but also to energy.

[2] More specifically, it has to be guaranteed that the environment memory is short-lived on the timescale of the coherent dynamics.

robust spin entanglement between the electron in the dot and reservoir. Only spin-selective scatterings, which are expected to occur on a much longer time scale [13], will eventually destroy the spin entanglement.

4 Summary In conclusion, we have proposed a scheme for an optically triggered spin entanglement of electrons in semiconductors. It consists of a single-electron doped quantum dot embedded in a field-effect structure. Optical excitation of an additional electron–hole pair (charged exciton) is transferred through tunneling to a photocurrent, where the spins of the electrons in the dot and reservoir are maximally entangled. This entanglement is robust against dephasing and relaxation processes which are not spin-selective, and thus benefits from the long spin lifetimes in semiconductors. The proposed device might be useful in future quantum information applications to establish entanglement between spatially separated sites.

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