Orbital and Spin Dynamics of Single Neutrally-Charged Nitrogen-Vacancy Centers in Diamond

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The neutral charge state plays an important role in quantum information and sensing applications based on nitrogen-vacancy centers. However, the orbital and spin dynamics remain unexplored. Here, we use resonant excitation of single centers to directly reveal the fine structure, enabling selective addressing of spin-orbit states. Through pump-probe experiments, we find the orbital relaxation time (430 ns at 4.7 K) and measure its temperature dependence up to 11.8 K. Finally, we reveal the spin relaxation time (1.5 s) and realize projective high-fidelity single-shot readout of the spin state (≥ 98%).

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Defect centers in solids are a promising class of systems for quantum science and technology. They combine bright optical transitions, access to long-lived electronic and nuclear-spin registers, and compatibility with solid-state device engineering. Of particular prominence is the negatively-charged nitrogen-vacancy center (NV−) in diamond, which has enabled recent advances in quantum information science and quantum sensing.

Alongside NV−, the nitrogen-vacancy defect can exist in both the neutral- (NV0) and positive- (NV+) charge states. These additional charge states can be used as a resource in a number of applications, such as spin-to-charge conversion for improved spin-state readout, classical data storage in NV ensembles, and deliberate charge-state switching for improved nuclear-spin coherence under ambient conditions.

Conversely, for experiments based upon NV−, undesired conversion to NV0 can be a hindrance: Active charge-state initialization protocols have been used to counter this. For quantum networks, stochastic conversion from NV− to NV0 is an important decoherence mechanism for nuclear-spin quantum memories.

Despite the importance of NV0, understanding of many of its properties remains elusive. In particular, the orbital- and spin-dynamic timescales are unknown. Also, recent magnetic circular dichroism (MCD) measurements on ensembles give insight into the NV0 fine structure, no direct observation has been reported.

Building an understanding of the system and its associated dynamic processes is important for improving control in NV quantum devices. Moreover, the knowledge gained may offer new insights into the physics of other impurities in solids.

Finally, NV0 may prove to be a powerful quantum system in its own right.

Here, we develop protocols combining resonant excitation of both NV0 and NV−. We apply these novel protocols to reveal the orbital and spin dynamics of single NV0 centers in diamond as well as to realize initialization and single-shot readout of the NV0 spin state. We perform our measurements on single NV centers at cryogenic temperatures; see Fig. 1(a). The NV center is addressed with microwave (mw) pulses (NV− ground-state spin transitions) as well as with polarization-controlled λred = 637 nm (NV− zero-phonon line (ZPL)) and λyellow = 575 nm (NV0 ZPL) laser light. We apply an axial magnetic field of Bz = 1890(5) G to induce significant Zeeman splitting.

The ZPL of the NV0 center has been conclusively attributed to this defect. A combination of ab initio calculations and symmetry arguments led to the proposal of ground states of 2E symmetry, which can be optically excited to a 2A2 manifold. An additional metastable 4A2 quartet state was also predicted and has been observed by electron paramagnetic resonance (EPR) measurements under excitation of the NV0 ZPL.

A splitting of the transitions of the two orbital states E+ and E− has been measured. However, the associated fine structure has not been observed in PL or EPR measurements.

We start by performing spectroscopy using the experimental procedure sketched in Fig. 1(b). For each frequency step, we (1) probabilistically prepare the emitter in NV0 by applying strong laser excitation resonant with the NV− ZPL, in combination with weak mw driving to induce the conversion NV− → NV0. We then (2) apply polarized
yellow light, during which time all single photons above 650 nm are integrated. The measured spectra [Fig. 1(c)] show four transitions—the first direct spectroscopic observation of the NV$^0$ fine structure. These observations validate the model of Barson et al. [16], and we hence follow their theoretical description below. Under the secular approximation, the ground-state Hamiltonian of NV$^0$ can be described by

$$H = g \mu_B \tilde{S}_z B_z + l \mu_B \tilde{L}_z B_z + 2\lambda \tilde{L}_z \tilde{S}_z + e_{\perp} (\tilde{L}_- + \tilde{L}_+).$$  (1)

$g$ is the spin $g$ factor, $\mu_B$ is the Bohr magneton, $l$ is the orbital $g$ factor, $\lambda$ is the spin-orbit interaction parameter, and $e_{\perp}$ is the perpendicular strain parameter. $\tilde{L}_z = \sigma_z$ and $\tilde{S}_z = \frac{1}{2} \sigma_z$ are the orbital and spin operators, respectively, defined in terms of the Pauli matrix $\sigma_z$, with $\tilde{L}_\pm = \frac{1}{\sqrt{2}}(\tilde{X} \pm i \tilde{Y})$ are the orbital operators defined within the basis of the strain eigenstates $\{X, Y\}$. The $z$ axis is defined parallel to the NV axis.

The resulting level structure is presented in Fig. 1(d). The $^2E$ ground state is composed of a pair of doublet states with opposite spin-orbit parity (lower spin-orbit branch: $\{|+, \downarrow\}, \{-, \uparrow\}$; upper spin-orbit branch: $\{|-, \downarrow\}, \{|+, \uparrow\}$). The degeneracy of each doublet is lifted by orbital- and spin-Zeeman contributions under the applied magnetic field. Conversely, the $^2A_2$ excited state exhibits no spin-orbit structure but is rather split by the spin-Zeeman effect alone. These contributions lead to four spin-conserving transitions. The contributing ground state for each observed transition is indicated in Fig. 1(c).

We find that the luminescence of the transitions depends significantly on the polarization of the excitation light [see Fig. 1(c)]. Differing transition amplitudes for orthogonal polarizations can be attributed to optical selection rules that are strongly dependent on $e_{\perp}$ [16,19]. Based upon these observations, we develop a method to extract $e_{\perp}$ and simultaneously the fine structure parameters of the NV$^0$ Hamiltonian [19]. By fitting spectra from three individual NV centers against our theoretical model, we find $l = 0.039(11)$ and $\lambda = 4.9(4)$ GHz. These values are roughly a factor of 2 larger than those found previously using NV-ensemble MCD measurements [36].

Crucially, the data in Fig. 1(c) show that resonant optical excitation in this magnetic field regime allows for state-resolved addressing, enabling the heralded preparation of specific states and investigation of the system dynamics. To date, only the excited-state lifetime $\tau_{\text{exc}}$ of 21 ns has been reported [38]. Here, we investigate the orbital- and spin-relaxation timescales of the ground state, $\tau_{\text{orbit}}$ and $\tau_{\text{spin}}$, respectively; see Fig. 1(d).

In order to unambiguously measure the dynamics of NV$^0$, we design and implement a charge-resonance (CR) protocol that realizes high-fidelity heralded preparation into NV$^0$, with the $\lambda = 575$ nm laser resonant with a chosen optical transition; see Fig. 2(a). The CR protocol (1) can be broken down as follows. First, a heralding signal confirms preparation in NV$^-$, with the $\lambda = 637$ nm lasers on resonance with the NV$^-$ transitions. Next, a strong red optical pulse induces charge state conversion, after which a chosen NV$^0$ transition is excited with yellow light. If the photon counts obtained during the “NV$^0$ check” exceed a preset threshold, the protocol is completed. Further details are given in Supplemental Material [19].

After the CR protocol, we perform the experimental sequence on NV$^0$ (2). Finally, we detect whether undesired conversion to NV$^-$ occurred during the experimental sequence and then perform readout of the NV$^0$ state (3). The number of repetitions of the experimental sequence (2) is chosen to minimize the overhead from the CR protocol while maintaining an NV$^3$ population above 85% and ranges from $N = 15$ to 1000 dependent upon the used yellow power. Note that the CR protocol prepares a specific spin state of the NV$^0$ center. For circular polarization, we typically start the experiment by heralding the $\downarrow$ spin state. For linear polarization, however, due to their close spectral vicinity, the CR check heralds either the $\downarrow$ or $\uparrow$ spin state.

In Fig. 2(b), we show time-resolved pump measurements. Here, the yellow laser is gated by an acousto-optic modulator (AOM), with a measured rise and fall time of 30(5) and 7(1) ns, respectively. Upon opening the AOM,
we observe a rapid increase in fluorescence due to optical cycling, which is then damped as population is pumped out of the driven state. By fitting the steady-state fluorescence counts for $L (H)$ polarization, we extract a saturation power of 2.5(2) [1.8(1)] nW and saturation counts of 105(2) [103(2)] kcts/s; see Fig. 2(c). As the optical power is increased, coherent optical Rabi oscillations are observed. In Fig. 2(d), we plot the fitted frequency of these oscillations, revealing the expected $\sqrt{P_{\text{yellow}}}$ dependence. When the AOM is closed, the fluorescence decays with $\tau_{\text{exc}} = 22(1) \text{ ns}$ [inset in Fig. 2(b)], which is consistent with the literature [38].

To uncover the recovery timescale after pumping, we turn to pump-probe spectroscopy. Example time traces are shown in Fig. 3(a). The resulting data are well described by an exponential recovery with a single timescale associated with how fast the system relaxes [19] once illumination is turned off. At the base temperature of our cryostat $[T = 4.65(3) \text{ K}]$, we extract $\tau_{\text{recovery}} = 0.43(6) \mu$s. We attribute these fast dynamics to orbital relaxation processes, i.e., $|+\rangle \leftrightarrow |-\rangle$ and $\tau_{\text{orbit}} = \tau_{\text{recovery}}$.

We repeat the pump-probe measurements across a range of temperatures. The fitted recovery times are shown as rates $R_{\text{recovery}} = 1/\tau_{\text{recovery}}$ in Fig. 3(b). After an initial linear increase, a rapid increase is observed at higher temperatures. At these higher temperatures, the required time resolution exceeds the AOM switching time constants, which we take into account in the fitting procedure [19]. The initial linear increase ($\propto T$) can be attributed to single-phonon processes, while high-order processes appear to govern the recovery rate at higher temperatures [39,40]. Here, we fit individually to a two-phonon Raman process ($\propto T^2$) and a two-phonon Orbach process ($\propto \exp[-\Delta/k_B T]$), giving $A = 0.53(3) \text{ MHz/K}$ [$A = 0.54(2) \text{ MHz/K}$] and $B = 1(1) \times 10^7 \text{ MHz}$ [$B = 1(4) \times 10^7 \text{ MHz}$] for the lower (upper) branch.

FIG. 2. Time-resolved resonant pump measurements. (a) Experimental sequence consisting of preparation (1), measurement (2), and readout (3) parts. (b) Fluorescence of NV$^0$ when driving the lower spin-orbit branch with $H$ polarization for $P_{\text{yellow}} = 2.4, 4.1, 10.2 \text{ nW}$ (bottom to top) averaged over at least $1 \times 10^6$ repetitions. Measurements have a timing resolution of 250 ps and are offset for clarity. Solid red lines are simulations of the full system dynamics with our theoretical model [19]. Inset: Decay of fluorescence counts after the AOM is closed. (c) Steady-state (ss) fluorescence counts as a function of $P_{\text{yellow}}$, for $H$ (squares) and $L$ (circles) polarization. The dashed line is a fit to the recovery behavior [19]. (b) Recovery rate $R_{\text{recovery}}$ as a function of the cryostat temperature. Circles (squares) describe data measured on the lower (upper) spin-orbit branch. Error bars for $R_{\text{recovery}}$ correspond to 1 s.d. fit errors. The solid lines are fits of form $f(T) = AT + B \exp[-\Delta/k_B T]$, giving $A = 0.53(3) \text{ MHz/K}$ [$A = 0.54(2) \text{ MHz/K}$] and $B = 1(1) \times 10^7 \text{ MHz}$ [$B = 1(4) \times 10^7 \text{ MHz}$] for the lower (upper) branch.
preparation of $\downarrow$ by applying 25 nW for 250 $\mu$s and proceed when more than 25 photons are detected. After a delay of 0.1 ms, we perform a charge-state check with red excitation, followed by a second yellow readout (again, 25 nW for 250 $\mu$s); see Fig. 2(a) (3). We then repeat this experiment, but with a delay of 10 s between the yellow readouts, allowing for relaxation processes to occur. The resulting histograms are shown in Fig. 4(a).

In the first case (dark colors), we observe a single dominant population which can be modeled by a Poissonian distribution with mean photon count 25.2(2) and that we attribute to $\downarrow$. In the second case (light colors), we additionally observe a second distribution, again modeled as a Poissonian distribution with mean photon count 0.171(4). A charge-state measurement of NV$^-$ performed before each readout shows that only a small fraction of the population ($P_{NV^-} \sim 1\%$) is found in the unwanted charge state—which we discard from the histograms—and that the majority of low-count events can be attributed to a dark state of NV$^0$. As the populations evolve without laser excitation, the dark state must be part of the ground-state manifold; we therefore assign this state to the second spin state $\uparrow$. A readout threshold of five photons [solid line, Fig. 4(a)] discriminates the two spin states.

We now sweep the delay time between initialization and readout. The measured populations of $\downarrow$ ($P_\downarrow$) and $\uparrow$ ($P_\uparrow$) are plotted in Fig. 4(b), showing relaxation to a mean population of 0.494(6). The data are consistent with a spin-1/2 $T_1$ process of characteristic timescale $\tau_{\text{spin}} = 1.51(1)$ s. Note that the observed value is a lower bound of the intrinsic spin relaxation, as it may be limited by leakage of resonant laser light. By setting the initial and long-time population in $\downarrow$ to be 1 and 0.5, respectively, we obtain a lower bound for the single-shot readout fidelity, $F_{\text{RO}} = \frac{1}{2}(F_\downarrow + F_\uparrow) \geq 98.2(9)\%$, where $F_\downarrow$ is the probability to assign $|s\rangle$ after preparing $|s\rangle$ [19].

To investigate the cycling nature of the driven optical transition, we now repeat the measurement under 5 nW of resonant yellow excitation; see Fig. 4(c). We find that $P_\downarrow$ decreases on a timescale faster than can be explained by spin relaxation alone, showing that the optical excitation induces spin pumping. Possible spin-mixing channels are given either in the $^2A_2$ excited state or via an intersystem crossing, which might be offered by the $^4A_2$ state. We also find a significant increase of $P_{NV^+}$ due to optically induced charge conversion [19,42]. However, this slows once $\downarrow$ is depleted, as $\uparrow$ is a dark state for optical excitation. Beyond this, $P_\uparrow$ reduces with $\tau_{\text{spin}}$, and charge conversion continues. We find a high state preparation fidelity for $\uparrow$ of $99.9^{+10.0}_{-10.0}$% after 600 ms but with an absolute population in the NV$^0$ $\uparrow$ state of only 22(2)%.

To reveal the respective rates, we develop a three-level rate equation model that we fit to our data, using the measured spin-relaxation time as a fixed input [solid lines, Fig. 4(c)] [19]. For the applied power of 5 nW, we extract characteristic timescales of 27(1) ms [90(4) ms] for the charge conversion (spin-pumping) process. From this, we can estimate the cyclicity of the $\downarrow$ state within this regime to be $0.98(8) \times 10^5$ cycles, mainly limited by recharging to NV$^-$ [19].

In a second experiment, the 5 nW yellow excitation is stroboscopically interleaved with strong NV$^- \rightarrow$ NV$^0$ ionization pulses [19]; see Fig. 4(d). Again, we observe a gradual decrease of $P_\downarrow$ and an increase of both $P_\uparrow$ and $P_{NV^-}$, but then $P_{NV^-}$ growth stops and even inverts. This observation can be explained via the picture that the removal of an electron from NV$^-$ prepares a random spin state in NV$^0$, eventually populating the dark state $\uparrow$. Competing rates between this spin-selection process and spin relaxation lead to the observed steady-state populations. We again fit a three-level rate equation model, using the previously obtained parameters as fixed inputs [19], and extract a timescale for ionization of 18(4) ms. The rate equation model does not accurately describe the behavior at long timescales, which is likely due to a reduction of the NV$^0$ spin-relaxation time under red excitation and strong NV$^-$ microwave driving [19].

As a final step, we develop a master equation simulation to capture the full dynamics of the NV$^0$ center [19]. In Fig. 2(b), we plot the simulated excited-state population (solid line), using the uncovered NV$^0$ timescales and spectral properties. We match the Rabi frequency to the
measured optical power and further include a spectral average over a Gaussian distribution of detuning values with $\text{FWHM} = 2\pi \times 20 \text{ MHz}$. We find excellent agreement with our experimental fluorescence data, emphasizing a consistent understanding of the NV$^0$ dynamics.

In conclusion, we have developed a novel toolbox for the study and control of single neutrally-charged NV centers in diamond. We have uncovered the dynamic timescales and study and control of single neutrally-charged NV centers in orbitals or by feedback based upon the NV$^0$ spin. We have demonstrated single-shot readout and initialization by microwave spin locking in both orbitals or by feedback based upon the NV$^0$ spin. Finally, at reduced temperatures that suppress the orbital dynamics, NV$^0$ may prove to be a powerful system for quantum technologies in its own right.

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