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Coherence and Control of Quantum Registers Based on Electronic Spin in a Nuclear Spin Bath

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We consider a protocol for the control of few-qubit registers comprising one electronic spin embedded in a nuclear spin bath. We show how to isolate a few proximal nuclear spins from the rest of the bath and use them as building blocks for a potentially scalable quantum information processor. We describe how coherent control techniques based on magnetic resonance methods can be adapted to these solid-state spin systems, to provide not only efficient, high fidelity manipulation but also decoupling from the spin bath. As an example, we analyze feasible performances and practical limitations in the realistic setting of nitrogen-vacancy centers in diamond.

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The coherence properties of isolated electronic spins in the solid state are often determined by their interactions with large ensembles of lattice nuclear spins [1,2]. The nuclear spin dynamics is typically slow, resulting in very long bath correlation times. Indeed, nuclear spins are one of the most isolated systems in nature. This allows us to decouple electronic spins from nuclear spins via spin-echo techniques [3,4]. More remarkably, controlled manipulation of the coupled electron-nuclear system allows one to exploit the nuclear spin bath as a long-lived quantum memory [5–7]. Recently this approach has been used to prove a single nuclear qubit memory with coherence times exceeding tens of milliseconds in room temperature diamond [8]. Entangled states composed of one electronic and two nuclear spins have also been probed [9]. The essence of these experiments is to gain complete control over several nuclei by isolating them from the rest of the bath. Universal control of systems comprising a single electronic spin coupled to many nuclear spins has not been proved yet and could yield robust quantum nodes for larger scale quantum information architectures.

In this Letter we describe a technique to achieve universal control of a portion of the nuclear spin bath and use it, together with an electronic spin, to build multiqubit registers. Our approach is based on magnetic resonance control techniques, but there is an essential difference with previously studied small quantum processors, such as NMR molecules. Here the boundary between qubits in the system and bath spins is ultimately dictated by our ability to effectively control the qubits. The interactions between the electronic spin and the nuclear qubit and bath spins are of the same nature, so that control schemes must preserve the desired interactions among qubits while attempting to remove the couplings to the bath. The challenge to overcome is then to resolve individual energy levels for qubit addressability and control, while avoiding fast dephasing due to the uncontrolled portion of the bath.

Before proceeding we note that proposals for integrating small quantum registers into a larger system capable of fault tolerant quantum computation or communication have been explored theoretically [10,11] and experimentally [12,13]. These schemes generally require a communication qubit and a few ancillary qubits per register. The first one couples efficiently to external degrees of freedom (for initialization, measurement and entanglement distribution), leading to easy control but also fast dephasing. The ancillary qubits instead are more isolated and act as memory or ancillas in error correction protocols. While our analysis is quite general as it applies to various systems, e.g., quantum dots in carbon nanotubes [14] or spin impurities in silicon [7], we will focus on the nitrogen-vacancy (NV) centers [4,8,9] in high-purity diamond with low paramagnetic impurity content. These are promising systems to implement hybrid quantum networks due to their long coherence times and good optical transitions that can be used for remote coupling among registers, without overly affecting the coherence of the register’s ancillary qubits [15]. Specifically, nuclear spins can remain coherent for a large number of electronic excitation cycles [8].

We focus on an electronic-spin triplet, as for the NV centers (see Fig. 1). Quantum information is encoded in the $m_s = \{0, 1\}$ Zeeman sublevels, separated by a large zero field splitting (making $m_s$ a good quantum number). Other Zeeman levels are shifted off resonance by an external magnetic field $B_z$ along the N-to-V axis. The Hamiltonian in the electronic rotating frame is

$$\mathcal{H} = \omega_L \sum I_z + S_z \sum A_j \cdot \vec{I} + \mathcal{H}_{\text{dip}}$$

$$= \frac{1}{2} - S_z \omega_L I_z + \sum A_j \cdot \vec{I} + \mathcal{H}_{\text{dip}}, \quad (1)$$

where $S$ and $I^z$ are the electronic and nuclear spin operators, $\omega_L$ the nuclear Larmor frequency, $A_j$ the hyperfine
couplings and $H_{\text{dip}}$ the nuclear dipolar interaction, whose strength can be enhanced by the hyperfine interaction [8]. When the electronic spin is in the $m_s = 1$ state, nearby nuclei in the so-called frozen core [16] are static (since distinct hyperfine couplings make nuclear flip-flops energetically unfavorable) and give rise to a quasistatic field acting on the electronic spin. Other bath nuclei cause decoherence via spectral diffusion [17], but their couplings, setting noise strength and correlation time, are orders of magnitude lower than the interactions used for control. While in the $m_s = 0$ manifold all the nuclear spins precess at the same frequency, the effective frequencies in the $m_s = 1$ manifold, $\omega_I^j$, are given by the hyperfine interaction and the enhanced $g$ tensor [18], yielding a wide range of values. Some of the nuclear spins in the frozen core can thus be used as qubits.

Control is obtained via microwave ($\mu$W) and radiofrequency (rf) fields. The most intuitive scheme, performing single-qubit gates with these fields and two-qubit gates by direct spin-spin couplings, is quite slow, since rf transitions are weak. Another strategy, requiring only control of the electronic spin, has been proposed [19,20]: switching the electronic qubit between its eigenstates induces nuclear spins rotations about two nonparallel axes that generate any single-qubit gate. This strategy is not the most appropriate here, since rotations in the $m_s = 0$ manifold are slow [21]. We thus propose another scheme to achieve universal control, using only two types of gates: (i) One-qubit gates on the electronic spin and (ii) Controlled gates on each of the nuclear spins. The first gate is simply obtained by a strong $\mu$W pulse. The controlled gates are implemented with rf pulses on resonance with the effective frequency of individual nuclear spins in the $m_s = 1$ manifold, which are resolved due to the hyperfine coupling and distinct from the bath frequency [22]. Achievable rf power provides fast enough rotations since the hyperfine interaction enhances the nuclear Rabi frequency when $m_s = 1$ [23]. Any other gate needed for universal control can be obtained combining these two gates. For example, we achieve a single nuclear qubit rotation by repeating the controlled gate after applying a $\pi$-pulse to the electronic spin. Controlled gates between nuclei can be implemented by exploiting the stronger coupling to the electron as shown in Fig. 2(a).

Avoiding direct nuclear interactions is faster as long as the hyperfine coupling is several times larger than the nuclear coupling.

Although selectively addressing ESR transitions is a direct way to perform a controlled rotation with the electronic spin as target, this is inefficient as the number of nuclear spin increases. The circuit in Fig. 2(b) performs the desired operation on a faster time scale.

When working in the $m_s = 1$ manifold, each nuclear qubit is quantized along a different direction and we cannot define a common rotating frame. The evolution must be described in the laboratory frame while the control Hamiltonian is fixed in a given direction for all the nuclei (e.g., the $x$ axis). This yields a reduced rf Rabi frequency $\tilde{\Omega} = \sqrt{\Omega^2 + \gamma_x^2}$ (where $\{\gamma_1, \gamma_2, \gamma_3\}$ define local quantization axes in the $m_s = 1$ manifold and $\Omega_{rf}$ is the hyperfine-enhanced rf frequency). The propagator for a pulse time $t_p$ and phase $\psi$ is

$$U_L(\Omega_{rf}, t_p, \psi) = e^{-i(\tilde{\Omega} t_p - (\lambda - \phi)\sigma_3/2)} e^{-i\tilde{\Omega} t_p \sigma_1/2} e^{-i(\lambda - \phi)\sigma_3/2},$$

where $\{\sigma_3, \sigma_5, \sigma_3\}$ are the Pauli matrices in the local frame and $\tan(\lambda) = \tan\varphi_1 / \cos\varphi_1$. An arbitrary gate $U = R_c(\gamma) R_x(\beta) R_z(\alpha)$ is obtained by combining $U_L$ with an echo scheme (Fig. 3), which not only refocuses the extra free evolution due to the lab frame transformation, but also sets the gate duration to any desired clock time. Fixing a clock time common to all registers is advantageous to synchronize their operation.

In order to refocus the fast electronic-spin dephasing due to the frozen core nuclear spins, we need to embed the control strategy described above in a dynamical decoupling scheme [24] without loosing universal control. The electron-bath Hamiltonian is given by Eq. (1), where the index $j$ now runs over the bath spins. Neglecting for now
the couplings among nuclei, we can solve for the electronic-spin evolution under an echo sequence. By defining $U_0$ and $U_1$ the propagators in the 0 and 1 manifolds and assuming the nuclear spins initially in the identity state (high temperature regime), we calculate the dynamics of the electronic spin, $\rho_e(t) = [1 + (00)|10\rangle\langle 01| + \text{H.c.})p_e(0)$, where $p_e(t) = \text{Tr}[U_1^*U_0U_1^*U_0^*]$ is the function describing the echo envelope experiments \[18,25\]. Since in the $m_t = 0$ subspace all the spins have the same frequency, $f_{ee} = (2\pi/\omega_I) = 0$ and the electron comes back to the initial state. Nuclear spin-spin couplings lead to imperfect echo revival and ultimately to decoherence via spectral diffusion \[17\]. The energy-conserving flip-flops of remote nuclear spins modulate the hyperfine couplings, causing the effective field at the electron to vary in time. The field oscillations can be modeled by a classical stochastic process. The evolution of the electronic spin is thus due to two processes that can be treated separately as they commute: the echo envelope calculated above and the decay due to a stochastic field, approximated by a cumulant expansion \[26\]. For Lorentzian noise with autocorrelation $G(t) = G_0e^{-\left(t/t_c\right)}$ we obtain a spin-echo decay $e^{-\left(\Omega_t^2/3\pi\right)^\dagger}$ for $t \ll t_c$. By using dynamical decoupling techniques \[27\] and selecting a cycle time multiple of the Larmor precession period it is possible to prolong the electronic coherence. Figure 4 shows how to combine the electron modulation with the sequence implementing spin gate effectiveness of these techniques relies on the ability to modulate the evolution on a time scale shorter than the noise correlation time. The noise of the electron-nuclear spin system is particularly adapt to these decoupling schemes. Consider, for example, the NV center. The measured electron dephasing $T_2^*$ time is about 1 $\mu$s in natural diamonds \[18\], as expected from MHz-strong hyperfine couplings. The intrinsic decoherence time $T_2$ can be orders of magnitude longer ($T_2 \approx 600 \mu$s), reflecting the existence of the frozen core, where spin flip flops are quenched. The frozen core radius is about 2.2 nm \[16\] and only spins with hyperfine coupling $\approx 2.5$ kHz contribute to spectral diffusion. Both the inverse correlation time (given by the nuclear dipolar coupling) and the noise rms (given by the coupling to the electron) are of order of few kHz. Dynamical decoupling schemes must thus act on time scales of hundreds of $\mu$s. This in turns sets achievable constraints on the gate speed.

The time of a conditional single nuclear spin rotation must be set so that $\Omega_N$ is much less than the frequency splitting between two neighboring spins (in terms of frequencies). Suppose we want to control an $n$-spin register without exciting bath spins at the Larmor frequency. The minimum frequency splitting between two nuclei in the $m_t = 1$ manifold will be at best $\delta \omega = \frac{\omega_n}{\pi}$ for nuclear frequencies equally spaced and $\omega_M$ the maximum hyperfine-induced effective nuclear frequency. The nuclear frequency spread due to the hyperfine interaction is then $\Delta E_N = \delta \omega$\[1\] - $\frac{\omega_M}{\pi}$. We want the Rabi frequencies to satisfy the constraints: $\Delta E_e > \Omega_e > \frac{n-1}{2} \omega_M > \Omega_{rf}$, where $\Delta E_e = 2g \mu_B B_z$ is the gap to other Zeeman levels ($m_e = -1$ for the NVc) and $\Omega_e$ the $\mu$w power. For a typical choice of 700 G magnetic field, we have $\Delta E_e = 2 \text{GHz}$ and $\omega_L = 0.8$ MHz. Assuming, $\omega_M = 20$ MHz and $n = 4$ spins, we obtain the following parameter window (in MHz) $2000 > \Omega_e > 25, \Omega_{rf} < 5$. The gate clock time can be as short as a few $\mu$s, allowing hundreds of gates in the coherence time.

Since the scheme presented is based on selective pulses, the dominant (coherent) error is due to off-resonance effects. If the Rabi frequency is much smaller than the offset from the transmitter frequency $\delta \omega$, the off-resonance spin will just experience a shift (Bloch-Siegert shift) of its resonance frequency, $\Delta \omega_{BS} = \delta \omega - \frac{1}{T} \int \sqrt{\Omega_{rf}(t)^2 + \delta \omega^2} dt' = -\Omega_{rf}^2 = -\frac{\Omega_{rf}^2}{2\omega_M}$. This results in an additional phase acquired during the gate time that needs to be refocused. Note that this error grows with the register size and constrains $\Omega_{rf}$. When reducing the Rabi frequency to achieve frequency selectivity, we have to pay closer attention to the rotating-wave approximation and consider its first order correction, which produces a shift of the on-resonance spin $\Delta \omega_{RWA} = \Omega_{rf}^2/4\omega_M$. Other sources of errors are evolution of bystander nuclear spins and couplings among spins. More complex active decoupling schemes \[28,29\] can correct for these errors, allowing to use the $m_t = 1$ manifold as a memory.

Advanced techniques like shaped pulses, with amplitude and phase ramping, composite pulses \[30\], optimal control theory or numerical techniques \[31,32\] can provide better fidelity. The analytical model of control serves then as an initial guess for the numerical searches. Pulses found in this way correct for the couplings among nuclei and are robust over a wide range of parameters (such as experimental errors or the noise associated with static fields). Table I shows the results of simulations in a fictitious NV system with 1–4 nuclear spins and effective frequencies in the $m_t = 1$ manifold ranging from 15 to 2 MHz. We searched numerically via a conjugate gradient algorithm for a control sequence performing a desired unitary evolution, varying amplitude and phase of the $\mu$w and rf fields. We then simulated the control sequence in the presence of noise, with contributions from a large, static field and a smaller fluctuating one. The projected fidelities in the...
absence of experimental errors are very high, a sign that
fast modulation of the electronic spin effectively decouples
it from the bath. Pulse robustness to noise is slightly
degraded as the spin number increases: as the noise induces
a spread of the electron frequency, it becomes difficult to
deal with. Pulse robustness to noise is slightly
increased by increasing the control field
intensity. Combining these pulses with dynamical decou-
ing the electronic frequency. The fidelity degradation is still modest
as expected, the optimal gate time increases with register size,
and rf Rabi frequencies are

\[
\text{F} = \frac{1}{\text{Tr}[(U_t^\dagger U_t)/2]}\text{' at the}
\]

optimal gate time. The gate is a π/2 rotation about the local x
axis for qubit 1 in a register of 1-4 nuclear qubits. Noise
parameters are \(T_2^* = 1.5 \mu s\) and \(T_2 = 250 \mu s\); maximum \(\mu\epsilon\nu\)
and rf Rabi frequencies are \(2\pi \times 10\text{ MHz}\) and \(2\pi \times 20\text{ kHz}\). As
expected, the optimal gate time increases with register size,
reflecting both more complex control required in a larger
Hilbert space and weaker hyperfine couplings to distant spins.
We obtained similar results for a CNOT gate between spin 1 and 2.

We thus develop a modular control scheme, scalable to
many registers and applicable to various physical
implementations.

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the David and Lucile Packard Foundation.

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<td>0.9999</td>
<td>0.9992</td>
</tr>
<tr>
<td>(F) (noise)</td>
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<td>0.9995</td>
<td>0.9975</td>
</tr>
<tr>
<td>time (T)</td>
<td>5.0 μs</td>
<td>5.5 μs</td>
<td>6.0 μs</td>
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TABLE I. Simulated fidelities \(\text{F} = \frac{1}{\text{Tr}[(U_t^\dagger U_t)/2\mu]}\) at the
optimal gate time. The gate is a π/2 rotation about the local x
axis for qubit 1 in a register of 1-4 nuclear qubits. Noise
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Division of Atomic, Molecular, and Optical Physics, May
[21] The rotation rates are faster for large external fields, which
however reduce the angle between the two rotation axes, thus increasing the number of switchings.
[22] This leaves open the possibility to implement collective
refocusing schemes on the bath spins.
[26] R. Kubo, in Fluctuation, Relaxation and Resonance in