A High-Fidelity Gateset for Exchange-Coupled Singlet-Triplet Qubits

Pascal Cerfontaine, René Otten, M. A. Wolfe, Patrick Bethke, and Hendrik Bluhm
JARA-FIT Institute for Quantum Information, Forschungszentrum Jülich GmbH and RWTH Aachen University, 52074 Aachen, Germany

Qubit arrays with access to high-fidelity single- and two-qubit gates are a key ingredient for building a quantum computer. As semiconductor-based devices with several qubits become available, issues like residual interqubit coupling and additional constraints from scalable control hardware need to be tackled to retain the high gate fidelities demonstrated in single-qubit devices. Focusing on two exchange-coupled singlet-triplet spin qubits, we address these issues by considering realistic control hardware as well as Coulomb and interqubit exchange coupling that cannot be fully turned off. We use measured noise spectra for charge and magnetic field noise to numerically optimize experimentally realistic pulse sequences and show that two-qubit (single-qubit) gate fidelities of 99.90% (≥ 99.69%) can be reached in GaAs, while 99.99% (≥ 99.95%) can be achieved with vanishing magnetic field noise as in Si.

Large arrays of well-controlled qubits with access to a high-fidelity gateset of single- and two-qubit gates are a key ingredient for building a quantum computer. Qubits based on electron spins in gate-defined quantum dots, which are among the main contenders for scalable quantum computing, have achieved sufficiently high single-qubit gate fidelities exceeding 99.9% [1, 2] using microwave control of individual spins. In contrast, the best two-qubit gate fidelities as well as single-qubit gates in two-qubit devices have not yet exceeded 98% [3, 4], still well below the threshold required for fault-tolerant quantum computation. A similar trend is seen in superconducting qubits, where gate fidelities tend to decrease as more qubits are added [5, 6]. This indicates that it is necessary to treat the implementation of single- and two-qubit gates simultaneously to address major challenges like residual interqubit coupling.

An alternative to single-spin qubits, known as singlet-triplet (S-T0) qubits, uses the \( m_s = 0 \) states of two electron spins. This allows all operations to be achieved with sub-GHz baseband control of the exchange interaction [10, 14], potentially avoiding hardware challenges when scaling up by eliminating the need for microwaves. Their single-qubit operations have already been demonstrated experimentally in GaAs with 99.5% fidelity [15] in close agreement with theoretical predictions of 99.8% [16]. Two-qubit gates have only been demonstrated using capacitive (Coulomb) coupling [17, 18]. The resulting fidelities (70-90%) are much lower than single-qubit gate fidelities due to the relatively weak Coulomb coupling. A promising alternative is to base operations on the much stronger exchange interaction so that two-qubit gates rely on the same ingredients as the single-qubit gates. Hence, comparable fidelities can be expected.

Early theory works on coupled S-T0 qubits [10, 14] have considered single- and two-qubit gates separately, or equivalently assumed exchange coupling that can be freely switched off, while more recent work also includes residual interqubit coupling [19, 20]. These are based on simplified models well suited for conceptual insight. They consider either Coulomb or exchange coupling but not both at the same time. For a significantly higher level of realism, we now include both coupling mechanisms simultaneously, as well as other experimentally relevant effects.

We develop experimentally realistic pulse sequences for single-qubit gates and a controlled-not (CNOT) gate that account for residual interqubit couplings. Their average gate fidelities range from 99.69% to 99.90% in GaAs and from 99.95% to 99.99% in Si devices (with vanishing magnetic field noise). With 99.90% in GaAs and 99.99% in Si, the CNOT fidelity is one to two orders of magnitude better than any experimental result on spin qubits to date. We also investigate how these fidelities scale with the noise strengths and manipulation time, and find a straightforward scaling law allowing our results to be adapted to different device parameters.

Our model includes all effects which we expect to be relevant in a realistic setting in a way that has been shown to predict achievable gate fidelities within a factor two for single-qubit gates [13]. These include interqubit Coulomb coupling, the finite dynamic range of interqubit exchange coupling, electric and magnetic noise with realistic noise spectra, finite pulse rise times and other constraints arising from hardware capabilities. To consider these effects we use a simulation-based approach that constrains the search space only to the extent imposed by hardware limitations. Since this method can be expected to yield the best possible fidelities achievable for the model at hand, we expect that such a comprehensive treatment is of high value for reliably predicting the performance of specific quantum-computing platforms. Thus, our results provide a complete recipe for high-fidelity control of S-T0 qubits that is directly applicable to two-qubit experiments with current technology.

Qubit model – The two-qubit system we consider consists of four linearly-adjacent quantum dots labeled pairwise by \( i, j \in \{1, 2, 3, 4\} \) in a semiconductor heterostructure. During qubit manipulation, each double dot is tuned in the (1, 1) charge state, where \((n, m)\) either represents the number of electrons in dots 1 and 2 or dots...
3 and 4. Metal top gates control the exchange interaction $J_{ij}$ between two adjacent dots (see Fig. 1) by either changing the detuning voltage $\epsilon_{ij}$ (affecting the energy difference between dots $i$ and $j$) or by direct control of the tunnel barrier $b_{ij}$. We focus on controlling the detuning since it is widely used in experiments, provides a good on-off ratio, and the effect of noise is better understood, however our approach can also be adapted to barrier control.

The computational subspace is spanned by the $m_s = 0$ singlet and triplet states of the double dots 1, 2 and 3, 4. We use the basis $|00\rangle = |\uparrow\uparrow\downarrow\downarrow\rangle$, $|01\rangle = |\uparrow\downarrow\uparrow\downarrow\rangle$, $|10\rangle = |\downarrow\uparrow\downarrow\uparrow\rangle$, and $|11\rangle = |\downarrow\downarrow\uparrow\uparrow\rangle$. Two other $m_s = 0$ states, $|\downarrow\uparrow\uparrow\uparrow\rangle$ and $|\uparrow\downarrow\uparrow\uparrow\rangle$, are dynamically accessible via the intermediate exchange $J_{23}$. Any occupation of these leakage states after a gate is a source of error and must be carefully avoided. Leakage into all other spin states has been experimentally shown to be small (on the order of $4 \times 10^{-4}$) due to the Zeeman splitting from an externally applied magnetic field of 500 mT.

In addition to nearest neighbor exchange, each spin experiences a different constant magnetic field $B_i$ that can be realized with micromagnets, gate-voltage-tuning of the electron $g$-factor or via dynamic nuclear polarization (DNP). The complete Heisenberg Hamiltonian of the four spins in terms of the mean magnetic field $B_G = \frac{1}{4} \sum_{i=1}^{4} B_i$ and the average magnetic field gradients across two adjacent dots, $b_{ij} = B_j - B_i$ is given by

$$H = \frac{1}{2} B_G \sum_{i=1}^{4} \sigma_i^{(z)}$$

$$+ \frac{b_{12}}{8} [ -3\sigma_1^{(1)} + \sigma_2^{(2)} + \sigma_3^{(3)} + \sigma_4^{(4)} ]$$

$$+ \frac{b_{23}}{4} [ -\sigma_2^{(1)} - \sigma_3^{(2)} + \sigma_3^{(3)} + \sigma_4^{(4)} ]$$

$$+ \frac{b_{34}}{8} [ -\sigma_3^{(1)} - \sigma_3^{(2)} - \sigma_4^{(3)} + 3\sigma_4^{(4)} ]$$

$$+ \frac{J_{12}}{4} \sigma_1^{(1)} \cdot \sigma_2^{(2)} + \frac{J_{23}}{4} \sigma_2^{(2)} \cdot \sigma_3^{(3)} + \frac{J_{34}}{4} \sigma_3^{(3)} \cdot \sigma_4^{(4)},$$

where $\sigma^{(i)}$ acts on the spin in quantum dot $i$ and all prefactors have units of angular frequency, so that $\hbar = 1$ in the Schrödinger equation. In addition to the Heisenberg Hamiltonian, we also include capacitive coupling between the qubits by adding the empirical model

$$H_c = E_c \frac{\partial J_{12}}{\partial \epsilon_{12}} \frac{\partial J_{34}}{\partial \epsilon_{34}} (I - \sigma_z) \otimes (I - \sigma_z)/4 \quad (2)$$

written in the $\{|T_0\rangle, |S\rangle\}^\otimes 2$ basis. This model reflects the notion that the two-qubit phase acquired by $|S\rangle$ is proportional to the detuning-dependent admixture of $\langle 0,2 |$ charge states in the two hybridized $|S\rangle$ states, which according to first order perturbation theory is given by $\partial J/\partial \epsilon$ for each qubit. The prefactor $E_c$ is the charge coupling energy corresponding to a full transition from $(1, 1)$ to $(0, 2)$ in each double dot.

Control model – For constant $b_{ij} \neq 0$, any desired unitary operation $U_t$ in the computational subspace can be generated by manipulating $\epsilon_{ij}$ and thus $J_{ij}$ as a function of time. However, a straightforward implementation of $U_t$ is complicated by the nonlinear and imperfectly known relation $J_{ij}(\epsilon_{ij})$, noise on the qubit control parameters $\epsilon_{ij}$ and $b_{ij}$, experimental constraints, and residual interqubit coupling. We incorporate these effects as follows when searching for pulse sequences $\epsilon_{ij}(t)$ that realize a given target unitary $U_t$.

In experiments, $\epsilon_{ij}(t)$ is typically controlled by baseband pulses from arbitrary waveform generators (AWGs). Our model includes constraints introduced by the AWG, like bounds, $\epsilon_{\text{min}} \leq \epsilon_{ij}(t) \leq \epsilon_{\text{max}}$, and a fixed sample rate $f_s$. Fixing the sample rate results in a piece-wise constant time trace $\epsilon_{ij,k}$ with time index $k = 1 \ldots N_{\text{seg}}$, where the number of segments $N_{\text{seg}}$ is related to the total gate time $T = N_{\text{seg}}/f_s$. Furthermore, we also consider the AWG’s limited bandwidth which leads to a smooth time trace $\epsilon_{ij}(t)$. We obtain $\epsilon_{ij}(t)$ by convoluting $\epsilon_{ij,k}$ with the impulse response of a typical experimental setup $b_{ij}$. We fix the last 4 samples of each pulse at $\epsilon_{\text{min}}$ to ensure that transients settle nearly completely and do not lead to significant errors in subsequent quantum operations. The phenomenological relation for tilt-control $J_{ij}(\epsilon_{ij}) = J_0 \exp(\epsilon_{ij}/\epsilon_0)$ allows us to obtain $J_{ij}(t)$ using $\epsilon_0$ and $J_0$ from Tab. I. Our model also includes decoherence from $\epsilon_{ij}$ noise, modelled via a full noise spectrum, and quasistatic $b_{ij}$ noise.

Numerics and optimization – We now outline how we determine $\epsilon_{ij,k}$ which generate a target unitary operation $U_t$ with high-fidelity. For given $J_{ij}(t)$ and $b_{ij}$, we first approximate the time-dependent Hamiltonian from Eq. 1 as piece-wise constant, with time steps chosen sufficiently small to cause negligible errors. We compute the matrix exponential of the 6-dimensional $m_s = 0$ subspace by direct diagonalization in each time step and obtain the full unitary operator $U_t$ from $t = 0$ to $T$ as a function of $\epsilon_{ij,k}$ and $b_{ij}$. We define $V_c$ as the truncation of $U_t$ into the four-dimensional computational subspace so that we
can compute coherent leakage $\mathcal{L}_c$ as the distance from unitarity, $\mathcal{L}_c = 1 - \text{tr}(V_c^2 V_c^\dagger)/4$. We also map the operator $V_c$ to the closest unitary $U_c$ in the computational subspace with the same global phase as $U_t$. This ensures that the distance from the target gate, $\Delta = U_t - U_c$ is independent of a global phase and does not contain dynamics in the leakage subspace.

To evaluate the separate effects of the slow noise contributions, we average over a discrete Gaussian distribution of $b_{ij}$ ($\epsilon_{ij}$) with standard deviation $\sigma_{b}$ ($\sigma_{\epsilon}$), and obtain the quantum process $\mathcal{E}_\epsilon$ ($\mathcal{E}_b$) by computing the unitary generated by Eq.\ref{eq:noise} for each noise offset. During gate optimization, we include fast charge noise fluctuations with a white noise spectrum $S_c(f) = S_0$ to obtain the process $\mathcal{E}_\epsilon$ in a computationally efficient manner from a Lindblad equation and a Markov approximation. We quantify the decoherence from each noise source by $\mathcal{I}_0 = 1 - F(U_c, \mathcal{E}_b)$, $\mathcal{I}_0 = 1 - F(U_c, \mathcal{E}_\epsilon)$ and $\mathcal{I}_0 = 1 - F(U_c, \mathcal{E}_b)$, where $F$ is the average gate fidelity \cite{31}.

In order to find the parameters $\epsilon_{ij,k}$ and $b_{ij}$ such that $U_c$ implements a chosen target gate $U_t$ with minimal decoherence and leakage, we use the Levenberg-Marquardt algorithm (LMA) to solve the nonlinear optimization problem

$$\min_{\epsilon_{ij,k}} |\Delta, \mathcal{I}_0, \mathcal{I}_0, \mathcal{I}_0, \mathcal{L}_c|^2.$$  \hspace{1cm} (3)

To speed up the convergence of the LMA, we use analytic derivatives of $U_c$ and $\mathcal{E}_\epsilon$ to efficiently derive all terms in Eq.\ref{eq:noise} with respect to $\epsilon_{ij,k}$ \cite{28}. The chance to find or approach the global minimum is improved by starting the optimization from many random seeds. Accurate control of $b_{ij}$ is often not available in experiments, so we do not explicitly optimize it. Instead, we fix $b_{ij}$ at a few experimentally feasible values and only optimize $\epsilon_{ij,k}$.

After the optimization, more accurate values for the infidelities and leakages are calculated by explicitly computing unitaries for a large number of noise realizations. We now include auto-correlated charge noise with spectra $S_{\alpha}(f) \propto 1/f^\alpha$. In addition to coherent leakage, we also compute the total leakage $\mathcal{L}$ by averaging $\mathcal{L}_c$ over all of the relevant noise contributions, and define the incoherent leakage $\mathcal{L}_1 = \mathcal{L} - \mathcal{L}_c$.

**Parameter values** – We have chosen the experimental singlet-triplet qubit parameters given in Tab.\ref{tab:1} for our optimization, and then generalize our findings to different gate durations and noise strengths as given in Fig.\ref{fig:2}(a) and in the supplementary Fig. S.9 \cite{28}. Since measured charge noise spectra in Si differ greatly between devices \cite{1, 22, 23}, we choose measured GaAs parameters, which are on the low end of what has been measured in Si. The high-frequency spectrum is of the form $S_{\alpha}(f) \propto 1/f^\alpha$ with $\alpha = 0.7$ \cite{22}. While the spectral density of charge noise has not been measured above a few MHz in GaAs or Si, these higher frequency regimes are still important to the gate dynamics. Therefore, we extrapolate the noise spectra with a cautiously optimistic scenario ($\alpha = 0.7$) and a pessimistic scenario ($\alpha = 0$, i.e. white noise). We match each extrapolated spectral curve to $S_0 = 4 \times 10^{-20} \text{V}/\text{Hz}$ at 1 MHz \cite{22}. For quasistatic charge and hyperfine noise we use the parameters given in Tab.\ref{tab:1} where we assume $\sigma_{b} = 0$ as a best case scenario for Si. In GaAs, typical experiments work with $b_{ij}$ from 0.1 ns$^{-1}$ to 7 ns$^{-1}$ using DNP \cite{18}. In Si, gate-voltage-tuning of the electron g-factor \cite{26} or micromagnets \cite{25} do not allow for large gradients, thus we assume values between 0.01 ns$^{-1}$ and 0.7 ns$^{-1}$ are possible. In both materials, leakage can be suppressed by ensuring $b_{23} \gg J_{23}$ as it makes spin exchanges across dots 2 and 3 energetically costly \cite{11}. For interqubit capacitive coupling we consider two cases, taking $E_c = 0$ as a lower and $E_c = 350 \mu\text{eV}$ as an upper bound. The upper bound is estimated from a typical charge stability diagram by determining the distance of the triple points belonging to the $(1,0)-(2,1)$ transition \cite{32}.

**Two-qubit gates** – We now use our qubit model and optimization strategy to search for a CNOT gate. The LMA finds solutions with high probability for $N_{\text{seg}} \geq 30$, typically within $10^4$ iterations, given our objective function and constraints. We calculate the gate fidelities post-optimization using the noise profiles described above, and show a representative pulse sequence in Fig.\ref{fig:4}(a), with $N_{\text{seg}} = 50$ and $E_c = 0$. The pulse duration is 50 ns and the field gradients are $b_{12} = -b_{34} = 1 \text{ns}^{-1}$ and $b_{23} = 7 \text{ns}^{-1}$, where a large $b_{23}$ was chosen to suppress leakage. This pulse exhibits a fidelity of 99.90% for $\alpha = 0.7$ and small leakage, $\mathcal{L} = 1.94 \times 10^{-5}$, while unitary errors are negligible by design of the objective function. We list the different contributions to the infidelity for various parameter sets in Tab.\ref{tab:1}. The first two columns show that for $\alpha = 0$ fast charge noise ($\mathcal{I}_0$) is the dominant contribution to the infidelity due to the higher noise level above 1 MHz, while for $\alpha = 0.7$ slow charge noise ($\mathcal{I}_0$) is most detrimental. In both cases, the gate fidelities are not limited by hyperfine noise ($\mathcal{I}_0$). We assess the gate’s performance for different noise strengths (see Fig. S.9) and find that the infidelity contributions scale quadratically over a wide range of $\sigma_{b}$, $\sigma_{\epsilon}$, and $\sqrt{S_{\alpha,0.7}}$. However, for large noise strengths a fourth

| $\sigma_{c}$ (mV) | $\sigma_{b}$ (mT) | $\sigma_{\epsilon}$ (mV) | $\sigma_{\epsilon_{ij}}$ (mV) | $J_0$ (ns$^{-1}$) | $J_0$ (ns$^{-1}$) |
|------------------|------------------|------------------|------------------|------------------|------------------|
| 8                | 0.3              | 8                | 0                | 1                | 1                |
| 2.072            | 0.272            | -5.460           | -5.460           | 2.460            | 2.460            |
| 1                | 0.1              |                  |                  |                  |                  |
| 500              | 20               |                  |                  |                  |                  |

**TABLE I** Experimental parameters for GaAs and Si. Units of inverse seconds denote angular frequencies. To represent slower gates in Si we scale the measured impulse response \cite{28} by a factor 10 in time, leading to the values for $f_s$ and $J_0$ given here.
order term becomes dominant, indicating that the gate decouples partly from slow noise. The fit coefficients including the fourth order term are given in the supplementary information [28]. We speculate that the quadratic contributions could be reduced further once they dominate the infidelity by improving the dynamical decoupling character. We note that $b_{23}$ noise has a far lesser effect on the gate’s performance than the intraqubit gradient noise since it only enters the Hamiltonian in the leakage subspace, indicating that one need not stabilize the interqubit magnetic field gradient in GaAs. When we repeat the CNOT optimization for $E_c = 350\,\mu\text{eV}$ we find very similar results ($99.92\%$ for $\alpha = 0.7$), suggesting that if capacitive coupling is modelled well it can also be easily compensated by appropriate pulses. Thus we only consider $E_c = 0$ for the further analysis of the GaAs CNOT gate.

In order to investigate the optimal gate speed, Fig.2(b) shows the scaling of the different noise contributions with the AWG’s sample rate $f_s$ for $\alpha = 0.7$. These figures are obtained without reoptimization, we only adjust the time and energy scales while keeping all noise strengths fixed. Thus, they serve as a lower bound for the achievable fidelities. In the presence of only hyperfine noise, faster gates are advantageous, as shown by the yellow curve. The infidelity from quasistatic hyperfine noise scales with $f_s^4$ for lower and $f_s^2$ for higher fidelities, consistent with the dependence on the noise strength. The infidelity from slow charge noise is indifferent to the gate speed, related to the fact that the number of coherent exchange oscillations is constant with respect to the detuning [22]. For pink noise we find $\mathcal{I}_k \propto f_s^{1-\alpha}$, which is consistent with $T_2(\varepsilon) \propto (dJ/d\varepsilon)^{-2}$ for white noise. The blue curve, which combines all noise contributions for $\alpha = 0.7$, indicates that even though the optimization was performed for $\alpha = 0$, the original sample rate $f_s = 1\,\text{GS/s}$ is nearly optimal.

Isotopically purified Si devices may prefer slower gate speeds since the lower magnetic field noise would lower the yellow curve in Fig.2(b). Since field gradients in Si are lower than in GaAs, we use $b_{12} = -b_{34} = 0.1\,\text{ns}^{-1}$ and $b_{23} = 0.7\,\text{ns}^{-1}$. If generating these gradients proves too challenging in Si, it is still possible to suppress leakage by reducing $J_{23}$ and thus increasing the gate duration further. The analysis of the optimized 500 ns long gate is shown in the last two columns of Tab.1 for $E_c = 350\,\mu\text{eV}$, and shows a higher resilience to quasistatic charge noise, suggesting that dynamical decoupling is more effective. The total infidelity is about a factor 10 lower than in GaAs, and reaches $99.99\%$ for $\alpha = 0.7$. If $S_0$ stays fixed at 1 MHz, we expect a further improvement for $\alpha = 1$, which was observed in Si for lower frequencies up to $320\,\text{kHz}$ [1].

**Single-qubit gates –** For closely spaced singlet-triplet qubits, residual interqubit exchange and capacitive coupling can complicate the parallel execution of single qubit gates. Thus, we now optimize generators of the single-qubit Clifford group, $U_x = X_{\pi/2} \otimes I$ and $U_y = Y_{\pi/2} \otimes I$ choosing $N_{\text{seg}} = 20$ (as Ref. [12]). We find that the $X_{\pi/2} \otimes I$ and $Y_{\pi/2} \otimes I$ gates have similar fidelities and thus only report the results for $X_{\pi/2} \otimes I$ here (details

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**Fig. 2.** (color online) (a) CNOT pulse sequence for each controllable exchange $J_{ij}(t)$ with $b_{12} = -b_{34} = 1\,\text{ns}^{-1}$ and $b_{23} = 7\,\text{ns}^{-1}$. Blue traces show the sample values to be programmed to the AWG with 1 ns time resolution. Red traces are convoluted with a measured impulse response, as seen by the qubit. The sequence was optimized for GaAs parameters with $E_c = 0$ and has a fidelity of $99.90\%$ (see Tab.1). (b) Separate infidelity contributions of the optimized CNOT gate as a function of the AWG’s sample rate for $\alpha = 0.7$. We assume the AWG’s rise times scale with the sample rate. The circles (triangles) represent the results obtained for $\alpha = 0$ ($0.7$) in Si and GaAs, while the dotted lines show the expected scaling. (c) Residual coupling compensating $X_{\pi/2} \otimes I$ pulse sequence optimized for $E_c = 350\,\mu\text{eV}$ and $J_{23} = 0.005\,\text{ns}^{-1}$ with $99.69\%$ fidelity in GaAs.
The table lists the infidelity and leakage contributions as defined in the text for two spectral noise densities $S_{\alpha}(f) \propto 1/f^\alpha$ with $S_{\alpha}(1 \text{ MHz}) = 4 \times 10^{-20} \text{ V}^2/\text{Hz}$. The total fidelities with all noise applied simultaneously are $99.75\%, 99.96\%$, and $99.99\%$ for $\alpha = 0$, $0.7$, and $1$, respectively. We obtain $99.69\%$ for $\alpha = 0.7$ and $99.99\%$ for $\alpha = 1$. Without capacitive coupling and residual $J_{33}$, we find a similar fidelity of $99.92\%$ (99.99\%) for $\alpha = 0.7$. A large capacitive coupling of $E_c = 350 \mu\text{eV}$ adds $\sim 0.4\%$ systematic error to the infidelity (see Fig.S.4). The error due to residual exchange scales quadratically with $J_{23}$, and reaches $\sim 10^{-3}\%$ for $J_{23} = 0.005 \text{ ns}^{-1}$ (see Fig.S.5). We can suppress these systematic errors by including residual coupling in the model for the optimization, recovering a high fidelity of $99.69\%$ (99.95\%) listed in Tab.III. Such a residual coupling compensating pulse sequence is shown for $X_{\pi/2} \otimes I$ in Fig.2(c). The resulting gate exhibits a non-idling identity by pulsing $J_{34}$ to suppress noise via dynamical decoupling. Systematic effects of $E_c$ are compensated by keeping the product of $J_{12}$ and $J_{34}$ low since the exponential model for $J_{ij}$ implies $\beta_{ij} \ll |J_{ij}|$ and thus $H_c \propto J_{12}J_{34}$. This can be seen in the interleaving nature of the exchange pulses in Fig.2(c).

Implications – Our work shows that exchange coupling is used to mediate two-qubit gates, interqubit Coulomb coupling must be considered but presents no major obstacle for high-fidelity control of spin-qubit arrays. Using parameters of current GaAs (Si) devices, single-qubit gate fidelities of at least $99.69\%$ ($99.95\%$) and two-qubit gate fidelities of $99.90\%$ ($99.99\%$) are attainable. We find straightforward scaling laws to extrapolate these gate fidelities to different gate durations and noise strengths, which is useful to assess the merits of different material systems and possible gains from improving noise levels. In order to obtain the best possible fidelity, detailed knowledge of the high-frequency charge noise spectrum is beneficial. We expect that our approach can also generate the additional non-Clifford gate required for universal quantum computation with high fidelity.

Our approach can be extended to assess more complex structures controlled by more scalable control hardware, and to determine the implication of qubit inhomogeneities on gate fidelities. In addition, it should be possible to implement our pulses experimentally using an in-situ tune-up procedure which can remove errors due to systematic inaccuracies in the model. For this reason, our results are a strong indication that S-T0 qubits can provide a complete high-fidelity gate set for universal quantum computation.

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