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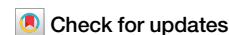
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Pulse shaping for ultra-fast adiabatic quantum gates

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A fundamental challenge in quantum computing is to increase the number of operations within the qubit coherence time. While this can be achieved by decreasing the gate duration, the use of shorter signals increases their bandwidth and can cause leakage into energetically separated states. A common method to suppress leakage for short pulses is the Derivative Removal by Adiabatic Gate (DRAG) method, which, however, relies on IQ modulation of radio-frequency (RF) signals, thus cannot be applied to the baseband signals, e.g., for semiconductor spin qubits. This paper proposes a novel technique, Delayed Leakage Reduction (DLR), that suppresses leakage at targeted frequencies even for baseband control by using time-delayed repetitions of the control signal to enable rapid, high-fidelity operations. We apply DLR on the adiabatic CZ gate between two spin qubits and achieve fidelities exceeding 99.9% within 9.4 ns for a resonance frequency difference of only 100 MHz. Towards the experimental realization of the proposed control method, we also assess the impact on the fidelity of the sampling rate, noise, and signal quantization of the electronic hardware generating the control pulse, thus setting the minimum hardware requirements for any experimental demonstration.

Semiconductor spin qubits^{1,2} are among the leading candidates for realizing scalable quantum computers thanks to their compatibility with classical CMOS fabrication technology^{3–7} and their high-temperature operation^{3,8–11}. While their potential for quantum computing has been demonstrated in various experiments^{12–30}, high-fidelity single and two-qubit gates are essential to achieve a fully operational quantum computer^{31,32}.

High-fidelity two-qubit gates between spin qubits are realized using the electrically tunable exchange coupling J , which arises from the overlap of the wavefunctions² and can be precisely controlled by electric signals applied to the gate electrodes. The exchange interaction can be used to implement the maximally entangling $\sqrt{\text{SWAP}}$, controlled X (CNOT), and the controlled Z (CZ) quantum gates, depending on the single-qubit Zeeman energies. The experimental implementation of $\sqrt{\text{SWAP}}$ operations is often challenging due to the stringent requirements on the resonance frequency difference ΔE_Z , $J \gg \Delta E_Z$ ^{33,34}, and rapid pulse rise times⁸. The highest reported fidelities were realized using the CZ gate²², which is implemented by varying the exchange interaction adiabatically with respect to ΔE_Z to suppress qubit excitations^{34,35}. In practice, the exchange interaction J is controlled by tuning the tunnel coupling between the two qubits via barrier gate electrodes, and ΔE_Z is realized either by local differences in the g-factor^{36–40} or designed via anisotropic magnetic field via micromagnets^{23,41,42}.

The CZ gate fidelity depends on the duration, resonance frequency splitting, decoherence, and shape of the control pulse^{24,35}. While fast gate times reduce decoherence, nonadiabatic pulses can cause the fidelity to rapidly oscillate as a function of the pulse duration due to interference with excited states. Therefore, high-fidelity CZ-gate implementations require precise timings commensurate with the frequency $\sqrt{\Delta E_Z^2 + J^2}$ ^{8,33,35,43,44}. In contrast, adiabatic pulses gradually vary the Hamiltonian parameters such that the system can follow the changes, resulting only in a phase accumulation during the evolution. Recent experiments employing adiabatic pulses have demonstrated two-qubit-gate fidelity beyond 99%^{22,24,25,27,29,30}.

The standard approach for mitigating errors in gate operations follows the optimal control theory^{45,46}, in which optimal pulses are derived using strategies commonly referred to as shortcuts to adiabaticity^{35,47–86}. For instance, the use of window functions for the pulse shape can improve the fidelity limited by non-adiabatic errors, as swap-like (flip-flop) errors³⁵ are related to the spectrum of the control pulse^{35,87}. A stronger suppression of non-adiabatic dynamics can be realized by using IQ modulation, such as DRAG⁴⁸ or oscillating signals around zero⁸⁸. However, since the exchange interaction for CZ gates in spin qubits is always a positive-valued baseband signal, these methods are not applicable.

In this work, we demonstrate the relation between the gate fidelity of an adiabatic CZ gate and the spectrum of the control pulse. Using

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the gained insight, we introduce a simple time-domain pulse-shaping strategy, Delayed Leakage Reduction (DLR), that suppresses spectral leakage to other energy levels using only baseband pulses. Our method combines two identical control signals with a time delay between them, specifically engineered to suppress leakage at targeted frequencies that could otherwise degrade the fidelity. Furthermore, we use our method to achieve a gate time of only 9.4 ns for $\Delta E_Z = 100$ MHz with fidelity $>99.9\%$, thus approaching the maximum possible gate speed. Additionally, we analyze the impact of a limited sampling rate for the adiabatic control signal on the fidelity to assess the feasibility of our proposed method in a practical experimental setup. By determining the minimum sampling rate required for adiabatic pulses, we remarkably find that the proposed DLR method only marginally increases the required sampling rate compared to traditional window methods.

In the Results section, we start by introducing the methods necessary to compute the fidelity and excitation rate of the evolution based on the spectrum of the control signal and ΔE_Z . The Delayed Leakage Reduction introduces our new pulse shaping method, DLR, that suppresses the spin flip-flop errors to undesired states. In the Nyquist criterion on adiabatic sampling, we derive and estimate the minimum sampling rate for adiabatic signals. Later, the Impact of noise introduces an approximate small-signal model that allows us to map nonlinearities and noise in the voltage domain into the corresponding exchange pulse. Finally, in the Discussion, we summarize our results and present prospective future work.

Results

Analysis of CPHASE gates between spin qubits

The controlled Z (CZ) gate is obtained as a special case of the controlled phase (CPHASE) gate, which naturally arises from the exchange interaction between spins. More precisely, the CZ operation can be achieved by composing a CPHASE gate with phase π and single-qubit Z operations

$$\text{CZ} = Z_1(-\pi/2)Z_2(-\pi/2)U_{\text{CPHASE}}. \quad (1)$$

This decomposition highlights that the two gates differ only by local, single-qubit phases, which often can be implemented virtually and do not affect entanglement.

The Hamiltonian of two interacting spin qubits via the exchange interaction can be written in the standard computational basis $\{|00\rangle, |01\rangle, |10\rangle, |11\rangle\}$ as³⁴

$$H_{\text{exc}} = \begin{pmatrix} E_z & 0 & 0 & 0 \\ 0 & (-\Delta E_Z - J)/2 & J/2 & 0 \\ 0 & J/2 & (\Delta E_Z - J)/2 & 0 \\ 0 & 0 & 0 & -E_z \end{pmatrix}, \quad (2)$$

where $E_z = (f_1 + f_2)/2$ and $\Delta E_Z = f_1 - f_2$ are the average and difference of the qubit resonance frequencies f_1 and f_2 , respectively. J represents the exchange interaction that can be controlled by detuning or tunneling the coupling of the spins. The time-evolution of this Hamiltonian directly implements $\sqrt{\text{SWAP}}$ and CPHASE operations depending on the choice of ΔE_Z and the waveform used for $J(t)$. To facilitate a more transparent analysis of the system dynamics, we decompose the total Hamiltonian H_{exc} into three mutually commuting components

$$H_{\text{RF}} = \begin{pmatrix} E_z & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & -E_z \end{pmatrix}, \quad (3)$$

$$H_{\text{CPHASE}} = \begin{pmatrix} 0 & 0 & 0 & 0 \\ 0 & -J/2 & 0 & 0 \\ 0 & 0 & -J/2 & 0 \\ 0 & 0 & 0 & 0 \end{pmatrix}, \quad (4)$$

$$H_{\text{SWAP}} = \begin{pmatrix} 0 & 0 & 0 & 0 \\ 0 & (-\Delta E_Z)/2 & J/2 & 0 \\ 0 & J/2 & (\Delta E_Z)/2 & 0 \\ 0 & 0 & 0 & 0 \end{pmatrix}. \quad (5)$$

By separating the Hamiltonian in this manner, we gain direct access to the individual unitary evolutions, enabling a clearer understanding of how each term contributes to the overall gate operation and how undesired dynamics can be suppressed. Since all three components commute, the total unitary corresponding to H_{exc} is given by the product of the time-evolution of each individual Hamiltonian. The Hamiltonians H_{CPHASE} and H_{SWAP} can be used to implement CPHASE and SWAP operations respectively, and H_{RF} implements the same single-qubit Z gate on both qubits simultaneously. As these joint single-qubit operations can be calibrated out by tracking the time and the resonance frequency or compensated by individual follow-up gates, we will ignore them for the remainder of this article.

To achieve high fidelity CPHASE (SWAP) operations, the effect of the other Hamiltonian, i.e., H_{SWAP} (H_{CPHASE}) must be suppressed, i.e., its resulting unitary should be as close as possible to the identity operator (I). For the CPHASE operation H_{CPHASE} , the corresponding time-evolution operator is given by

$$U_{\text{CPHASE}} = \begin{pmatrix} 1 & 0 & 0 & 0 \\ 0 & e^{i\pi \int_0^t J(t)dt} & 0 & 0 \\ 0 & 0 & e^{i \int_0^t J(t)dt} & 0 \\ 0 & 0 & 0 & 1 \end{pmatrix}. \quad (6)$$

By choosing $\int 2\pi \cdot J(t) dt = \pi$, the diagonal terms are $(1, i, i, 1)$, so that a CZ gate can be implemented using Eq. (1).

A high-fidelity CPHASE gate can be realized if the time evolution of the SWAP Hamiltonian U_{SWAP} is equal to the identity operator up to single-qubit phase errors. Before addressing the general case of a time-varying J , we briefly analyze the case for the constant J

$$U_{\text{SWAP}} = \frac{1+ZZ}{2} + \cos(\pi f_{\text{osc}} t) \frac{1-ZZ}{2} - i \sin(\pi f_{\text{osc}} t) \times \left[\frac{\Delta E_Z}{2f_{\text{osc}}} (IZ - ZI) + \frac{J}{2f_{\text{osc}}} (XX + YY) \right], \quad (7)$$

where $PQ = P \otimes Q$ with $P, Q = I, X, Y, Z$ being two-qubit Pauli operators and $f_{\text{osc}} = \sqrt{\Delta E_Z^2 + J^2}$. Residual single-qubit phase errors (ZI and IZ) can be compensated in this case with cascaded virtual single-qubit $Z(\theta)$ gates²⁴. However, the presence of the non-diagonal terms (XX and YY) leads to leakage to other energy levels and prevents the adoption of simple phase compensations. To suppress the non-diagonal terms, three commonly employed strategies are available. Firstly, by choosing $J \ll \Delta E_Z$, the amplitude of the non-diagonal terms in Eq. (7) will be $J/f_{\text{osc}} \approx J/\Delta E_Z \ll 1$. This reduces the oscillation amplitude, enabling the desired fidelity, but requires a longer gate time to accumulate sufficient phase in Eq. (6) due to the smaller value of J or asks for a larger frequency separation ΔE_Z , which can be challenging to achieve in practice. Secondly, the gate time can be chosen as an integer multiple of f_{osc} , which leads the $\sin(\pi f_{\text{osc}} t)$ to be zero at the end of the evolution. However, this method, called synchronization, requires extremely precise pulse timing, especially for faster gates³⁵. Lastly, pulse shapes for $J(t)$ that follow classical window functions can be used^{35,37}.

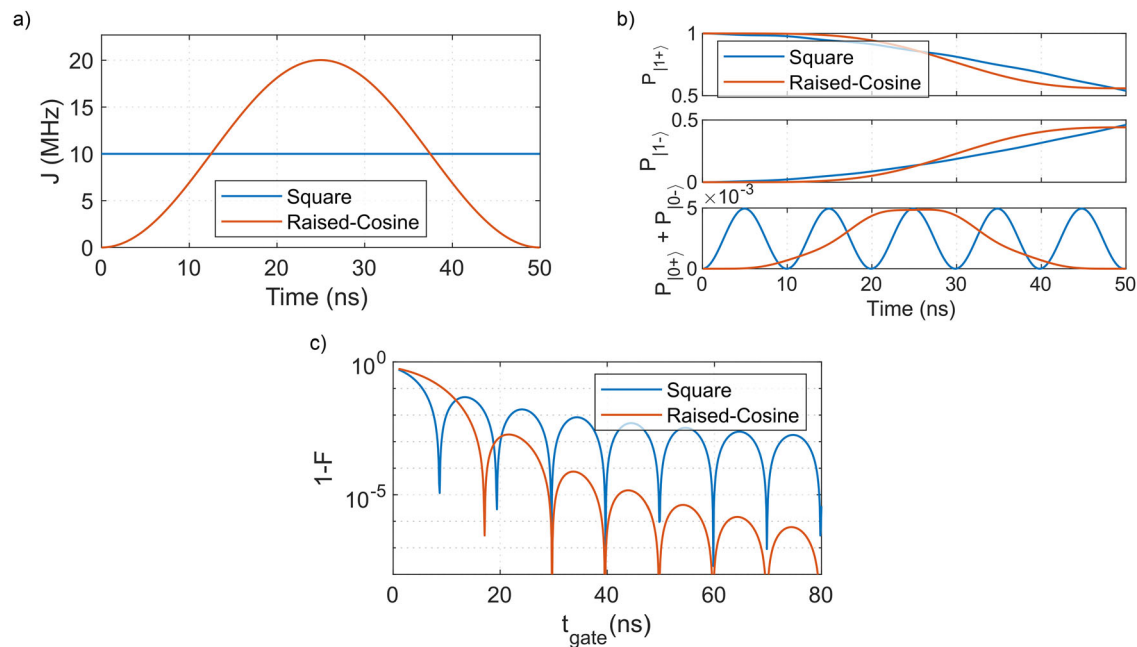


Fig. 1 | Impact of the pulse shape on the CZ gate. Simulated CZ operation on H_{exc} with $\Delta E_Z = 100$ MHz for square (blue) and raised-cosine (red) J control signals. **a** Time evolution of $J(t)$, with duration $t_{\text{gate}} = 50$ ns and amplitude tuned for $\int 2\pi \cdot J(t) dt = \pi$. **b** State probabilities of $P_{|1+⟩}$, $P_{|1-⟩}$ and $P_{|0+⟩} + P_{|0-⟩}$ during the

operation. **c** Fidelity of the CZ operation when sweeping the gate duration (t_{gate}) for both square (blue) and raised-cosine (red) shaped J signals, assuming perfect single-qubit phase compensation. For each point, the amplitude of the signal is optimized for the highest fidelity.

This method enforces slow changes of the eigenvectors of the system, thus resulting in a higher likelihood of ending in the initial state. We will examine this technique in more detail in the next subsection, where the assumption of a time-constant J is dropped.

Relating non-adiabatic errors to spectral properties

In Fig. 1, we compare the pulses and the resulting evolutions for the two commonly adopted strategies, synchronization and the usage of window functions for pulse shaping. In an ideal CPHASE operation, only the transition from $P_{|1+⟩}$ to $P_{|1-⟩}$ should occur, with both populations reaching 1/2 at the end of the gate. However, the presence of the undesired SWAP components induces additional oscillations in the residual population $P_{|0+⟩} + P_{|0-⟩}$, and may prevent $P_{|1+⟩}$ and $P_{|1-⟩}$ from reaching the aimed 1/2. Although $P_{|1+⟩}$ and $P_{|1-⟩}$ states can be corrected to 1/2 via single qubit Z calibration (see the Phase calibration, final gate time, and average gate fidelity in the Methods section), compensating the residual state populations $P_{|0+⟩} + P_{|0-⟩}$ requires the implementation of real single qubit operations.

Furthermore, choosing a smaller $J/\Delta E_Z$ ratio will reduce the amplitude of the oscillations in the residual state at the cost of longer gates for fixed ΔE_Z . As an alternative, since the residual state evolution is periodic with frequency f_{osc} for the square pulse, a perfect gate can be achieved by synchronizing the gate time with the gate oscillations^{33,35,43,44}. However, using a perfect square pulse requires a wide bandwidth and precise timing, as highlighted in Fig. 1c. In contrast, a smooth adiabatic pulse, e.g., a raised-cosine pulse, greatly decreases the residual state error (see Fig. 1b) during the evolution, thus allowing for faster gate times without strict requirements on the timing of the control signal. As a result, adiabatic pulses longer than 15 ns can achieve a higher gate fidelity than the square pulses (see Fig. 1c). However, shorter high-fidelity gates cannot be implemented straightforwardly using a raised-cosine (or other window function) pulse³⁵.

To better understand the characteristic dynamic of the adiabatic evolution³⁵, we make use of an approximate analytical expression of the non-adiabatic error (see the Explicit expressions of errors in the Methods

section)

$$\epsilon = |\text{tr}(\mathcal{E}|01\rangle\langle 10|)|^2 \approx \left| \pi \int_0^{t_{\text{gate}}} J(t) e^{i2\pi(\alpha_1(t) - \alpha_2(t))} dt \right|^2. \quad (8)$$

where α_1 and α_2 are the accumulated phase of the eigenstates $|01\rangle$ and $|10\rangle$, i.e., $\alpha_{1,2}(t) = \pm \int_0^t f_{\text{osc}}(t')/2 dt'$, and t_{gate} is the gate duration. This expression can be interpreted as the control signal interacting with the frequency of one state ($e^{i2\pi\alpha_1}$) and being observed in a rotating frame of the other qubit state ($e^{-i2\pi\alpha_2}$). To make this formulation more practical for spectral analysis, we rewrite the phase difference in the exponent by introducing and subtracting ΔE_Z . This isolates a term that oscillates at ΔE_Z as

$$\epsilon \approx \left| \pi \int_0^{t_{\text{gate}}} J(t) e^{i2\pi \int_0^t (f_{\text{osc}}(t') - \Delta E_Z) dt'} e^{i2\pi t \Delta E_Z} dt \right|^2. \quad (9)$$

For $J \ll \Delta E_Z$, the difference $f_{\text{osc}}(t) - \Delta E_Z$ becomes negligible and the integral simplifies. Under this approximation, the error can be directly related to the magnitude of the Fourier spectrum of $J(t)$ at frequency ΔE_Z

$$\epsilon \approx |\pi S(-\Delta E_Z)|^2. \quad (10)$$

where

$$S(f) = \int_0^{t_{\text{gate}}} J(t) e^{-i2\pi f t} dt. \quad (11)$$

is the Fourier transform of $J(t)$.

The average gate infidelity of the CZ gate is then given by (see the Phase calibration, final gate time, and average gate fidelity in the

Methods section)

$$1 - F \approx \frac{2}{5} (\epsilon + 2\delta\Phi_f^2 + \delta\Phi^2 + \delta\theta^2), \quad (12)$$

where $\delta\Phi_f = \pi - 2\pi \int_0^{t_{\text{gate}}} J(t) dt$ is the exchange miscalibration, $\delta\Phi = (\theta_1 + \theta_2 - \pi) - 2\pi \tau_z t_{\text{gate}}$ the miscalibration of the rotating frame, $\delta\theta = (\theta_1 - \theta_2) - 2\pi \int_0^{t_{\text{gate}}} \sqrt{\Delta E_z^2 + J^2(t)} dt$ the miscalibrated phase corrections, and $\theta_{1,2}$ is the single-qubit phase gate $Z(-\theta_{1,2})$ for qubit 1 and 2.

In Fig. 2, the fidelity computed using Eqs. (8) and (10) is compared with the simulation results. For different gate durations t_{gate} , the amplitude of the signal is calibrated such that $\int_0^{t_{\text{gate}}} 2\pi J(t) dt = \pi$ is fulfilled. This leads to an increase in J for shorter pulses, and the approximation $f_{\text{osc}}(t) - \Delta E_z \approx 0$ does not hold anymore. As a consequence, Eq. (10) loses its accuracy for shorter gates, while Eq. (8) fits the simulations better. This comparison confirms that the analytical model accurately captures the fidelity behavior, especially for longer gate durations. It is important to note that the spectral analysis of gate fidelity is equally applicable to square pulses. Such a spectrum exhibits notches at approximately every integer multiple of $1/t_{\text{gate}}$ reflecting the synchronization condition, and high fidelity gates are obtained when one of these notches aligns with frequency $-\Delta E_z$.

This spectral perspective reframes the challenge of minimizing gate infidelity as a problem of pulse shaping or spectral design. In the next section, we introduce a novel time-domain technique—Delayed Leakage Reduction (DLR)—designed to suppress spectral leakage at critical frequencies.

Delayed leakage reduction

To minimize leakage at ΔE_z according to Eq. (10), a delayed replica of the control signal can be combined with the original control signal to cancel out spectral components at the critical frequency $f = -\Delta E_z$, thus maximizing the gate fidelity. If $g(t)$ is the original pulse used to modulate J , e.g. the raised-cosine pulse described in the previous section, the DLR pulse is defined as

$$g'(t) = g(t) + g(t - t_d) \quad (13)$$

where t_d is the delay between two pulses. In the Fourier transform, this leads to a frequency-dependent phase shift of the delayed replica as

$$S(f) = [1 + e^{-i2\pi f t_d}] \int_0^{t_{\text{gate}}} g(t) e^{-i2\pi f t} dt \quad (14)$$

$$= [1 + e^{-i2\pi f t_d}] G(f). \quad (15)$$

where $G(f)$ is the Fourier transform of $g(t)$. By choosing $t_d = \frac{1}{2\Delta E_z}$, a notch in the spectrum at $S(-\Delta E_z)$ is created, thus suppressing the non-adiabatic error. Figure 3 demonstrates an example of a 50 ns DLR-shaped pulse, where $g(t)$ is chosen as a 45 ns raised-cosine pulse and $t_d = 5$ ns (Fig. 3a), thus showing a clear notch at $f_{\text{notch}} = \frac{1}{2t_d} = 100$ MHz (Fig. 3b). In Fig. 3c, we compare the resulting fidelity of the CPHASE operation to a straightforward DLR implementation (DLR-static) for commonly used window functions. While DLR-static outperforms the raised-cosine beyond 30 ns, the fidelity of the DLR pulse drops significantly for shorter gate times due to $f_{\text{osc}} \approx \Delta E_z$ becoming invalid. To overcome such a limitation, more advanced DLR techniques are proposed in the following.

As J changes during the gate operation, so does f_{osc} , especially for shorter gate durations. To cope with this, the DLR-average strategy places the notch in the spectrum at the average of the maximum and minimum value of $f_{\text{osc}}(t)$ during the operation. However, calculating the required t_d from $f_{\text{osc}}(t)$ is not straightforward since $f_{\text{osc}}(t)$ is a function of $J(t)$, which is also a function of t_d . Therefore, $f_{\text{osc}}(t)$ is first calculated with a raised-cosine pulse for a given gate time, and t_d is computed accordingly. While this is not

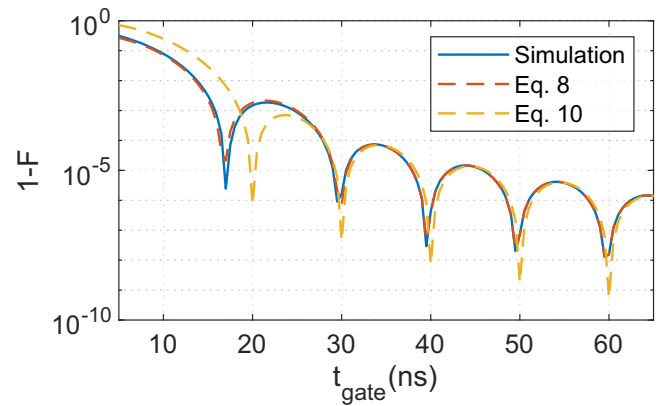


Fig. 2 | Benchmarking the analytical fidelity estimation with simulations. The simulated fidelity of the CPHASE gate using a raised-cosine pulse is compared to error rate equations as a function of gate time. The raised-cosine pulse amplitude is optimized for each gate duration to ensure precise phase accumulation, and ΔE_z is 100 MHz.

an exact solution, DLR-average can reach lower infidelities with faster gates than the DLR-static strategy (Fig. 3c).

For the last method, DLR-dynamic, we use the time-dependent delay $t_d(t) = \frac{1}{2f_{\text{osc}}(t)}$. Since in this method $J(t)$ and $f_{\text{osc}}(t)$ are recursive as for DLR-average, $f_{\text{osc}}(t)$ is first calculated with a simple Raised-Cosine pulse. To preserve the gate time, we redefine $g(t)$ as

$$g(t, t_{\text{gate}}) = \begin{cases} 0, & t < 0 \\ \frac{1}{2} \left(1 - \cos\left(\frac{2\pi t}{t_{\text{gate}}}\right) \right), & 0 \leq t \leq t_{\text{gate}} \\ 0, & t > t_{\text{gate}} \end{cases} \quad (16)$$

Then the pulse for $J(t)$ is given by $g(t, t_{\text{gate}} - t_d(t)) + g(t - t_d(t), t_{\text{gate}} - t_d(t))$, where t_{gate} is the required gate time. As shown in Fig. 3c, using the DLR-dynamic (purple), after a gate time of 9.4 ns, the infidelity never exceeds 8.9×10^{-4} .

Although those strategies can improve the fidelity below 15 ns with respect to DLR-static, the appearance of an infidelity hump between approximately 20–25 ns (see Fig. 3c) does not change. To explain this, we find that for a gate time of 20–25 ns, the lowest-frequency notch of the DLR pulse spectrum is only at 80–100 MHz. Since the slope of the main lobe is steeper than the other lobes, the notch introduced by the DLR-envelope cannot properly attenuate the spectrum.

Although the control signal discussed so far has been constructed only by adding a delayed replica, it is also possible to use multiple delayed replicas with different amplitudes and delays to customize the spectrum for specific applications, e.g., by reducing the leakage to various energy levels or for stronger suppression.

Numerical optimization for other windows with DLR

Adopting a more generic pulse for the J -control signal, such as a window function, may become analytically challenging, because, for instance, window functions with non-zero initial or final value may introduce discontinuities during the DLR pulse due to the delayed replica. Such discontinuities can cause a rapid change in f_{osc} during the operation, thus invalidating Eq. (9)³⁵. However, the optimal delay t_d for a DLR-static pulse can be numerically identified for a given gate time and window shape by directly simulating the fidelity performance of each sample waveform. Thanks to its simplicity, DLR can be implemented using arbitrary window functions. In Fig. 4c, the best fidelity for the given t_{gate} is shown for Hamming, Kaiser, and the raised-cosine. A hump similar to the one observed in Fig. 3c can be clearly seen around 20–25 ns for the same reason. Due to the different shapes of the main lobe for the three different functions, each pulse

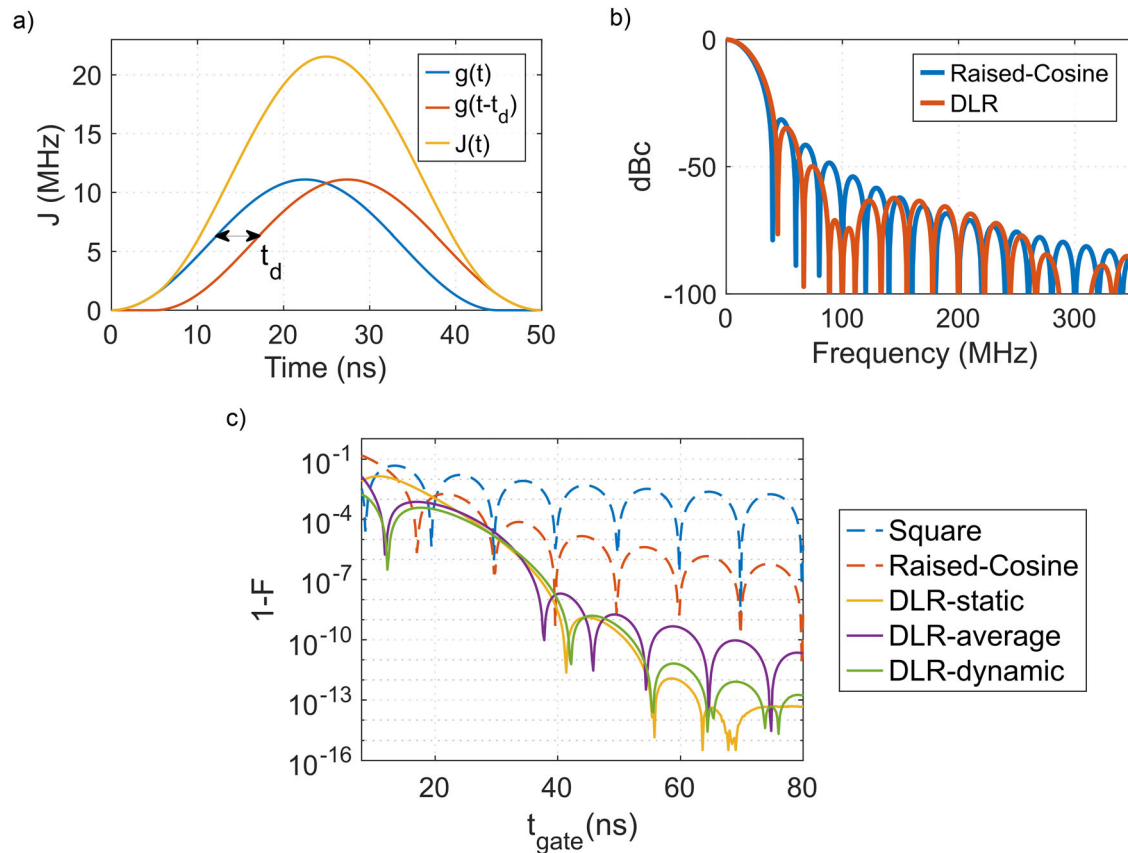


Fig. 3 | DLR pulse and its impact on fidelity. **a** Example of a 50 ns signal generated with DLR. It consists of two raised-cosine pulses, $g(t)$ and $g(t-t_d)$, which have a delay of $t_d=5$ ns between each other to create a notch at $\Delta E_Z=100$ MHz. **b** The frequency spectrum of two $J(t)$ signals, a raised-cosine and DLR, with a 50 ns gate duration. Since DLR consists of two shorter pulses, its main and side lobes are wider. Notches at integer multiples of $\frac{1}{2t_d}=100$ MHz are clearly visible. **c** Fidelity of the CPHASE operation in a two-qubit system with $\Delta E_Z=100$ MHz for different gate durations

(t_{gate}). For every gate time and signal shape, the amplitude of $J(t)$ is tuned for maximum fidelity, according to Eq. (6). The DLR-static and DLR-average introduce a spectral notch at fixed frequencies, namely 100 MHz and $\frac{1}{2}(100 + \max[f_{\text{osc}}(t)])$, respectively. In contrast, the DLR-dynamic continuously adapts the notch frequency throughout the gate operation, following the instantaneous oscillation frequency $f_{\text{osc}}(t)$.

has a different hump over different values of t_{gate} . The noisy behavior in those results is attributed to numerical noise, e.g., a too large step size in the t_d sweep.

The results demonstrate that high-fidelity operations exceeding 99.99% are achievable across a broad range of gate times, provided that the notch frequency is appropriately tuned. The error range of highly tuned t_d for each window can be found in Fig. 12d–f.

Comparing DLR and DRAG

The Derivative Removal by Adiabatic Gate (DRAG) protocol is one of the most common quantum control methods used to reduce leakage to higher energy levels. This is achieved by adding correction terms to the control signals to mitigate the leakage to unwanted levels. The first-order DRAG pulse is constructed by adding the derivative of the original pulse, scaled by a specific coefficient, to the original signal

$$J(t) = g(t) + i\beta \frac{d}{dt}g(t), \quad (17)$$

where β is a scaling coefficient. The resulting spectrum is

$$S(f) = G(f) + i2\pi f(i\beta G(f)) \quad (18)$$

$$= G(f)[1 - 2\pi f\beta]. \quad (19)$$

where $G(f)$ is the spectrum of $g(t)$.

Setting $\beta = -\frac{1}{2\pi\Delta E_Z}$ ensures $S(-\Delta E_Z) = 0$, minimizing leakage at the oscillation frequency. In Fig. 5b, the frequency spectrum of DRAG and DLR is compared. While the spectrum of the DRAG technique is suppressed on a single side, the DLR technique achieves suppression at both sides.

In Fig. 5c, the DLR-average and DRAG techniques applied to the raised-cosine pulse are compared as a function of the gate duration. For a fair comparison, β is also chosen to create the notch at the average of the maximum and minimum values of f_{osc} . The plots for both methods are similar, except that DLR follows the DRAG with a small shift in gate duration. A closer inspection indicates that the value of the shift is approximately the time delay t_d of our DLR protocol. As we have discussed in the Delayed Leakage Reduction, DLR expands the spectrum due to the use of two pulses that are t_d shorter than the original pulse, thus expanding the spectrum of the resulting pulse. This expansion explains why the DRAG pulse outperforms DLR slightly.

However, from a practical standpoint, implementing a DRAG pulse requires access to the complex domain, e.g., by adopting IQ bandpass modulation as usually done for transmon qubits^{48,89}, which is infeasible for baseband pulses such as the exchange interaction for the CPHASE gate implementations²⁴ and similar situations with limited control. Additionally, DLR features an always symmetric pulse shape. This symmetry could be particularly advantageous in scenarios such as simultaneous spin-qubit driving⁹⁰, where control pulses must selectively address a target qubit without inducing leakage or excitations in neighboring qubits.

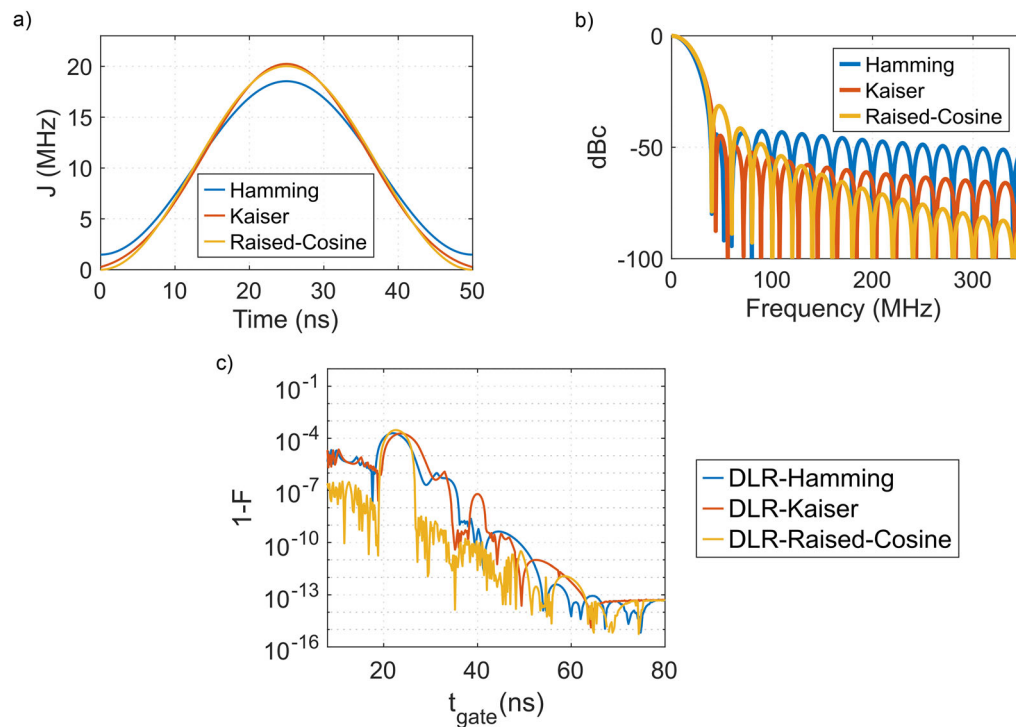


Fig. 4 | DLR pulse shape comparison. **a** Hamming, Kaiser ($\beta = 6.14$) (see the Window Functions for Spectral Shaping in the Methods section), and raised-cosine window functions with 50 ns gate time. **b** Frequency spectrum of the window functions in (a). **c** Infidelity of the CPHASE gate controlled by the DLR-static pulse using each of the window functions in (a). For each window function, the DLR-static

pulse is used to derive the gate fidelity. Among all delays generating a notch between 80–130 MHz (see the Additional Data of DLR Numerical Optimization in the Methods section), the one with the minimum infidelity has been selected and constitutes a data point in the plot.

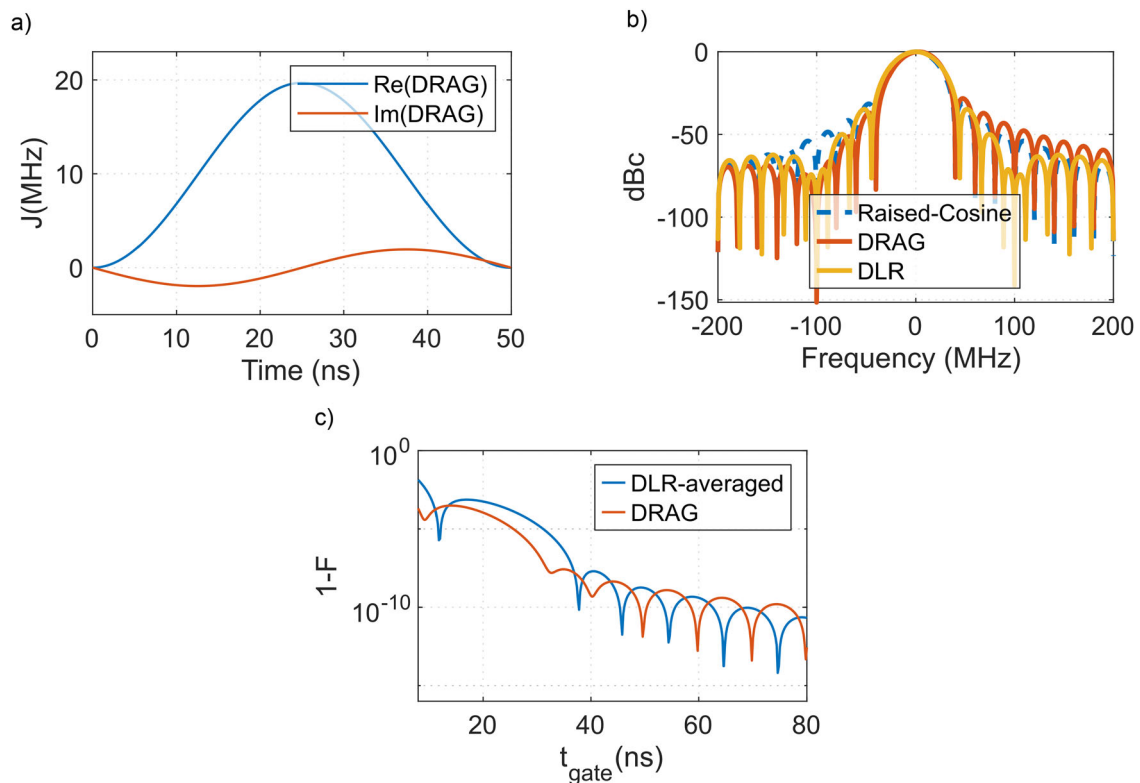


Fig. 5 | Comparison between DLR and DRAG pulses. **a** Example of J pulse generated with DRAG where $\beta = 1/(2\pi\Delta E_Z)$, with $\Delta E_Z = 100$ MHz. **b** Comparison of the DRAG and DLR frequency spectrum with 50 ns t_{gate} and notch at 100 MHz.

c Infidelity of DLR and DRAG signals for different gate durations (t_{gate}). The signal amplitude is optimized for each t_{gate} to achieve accurate π phase accumulation.

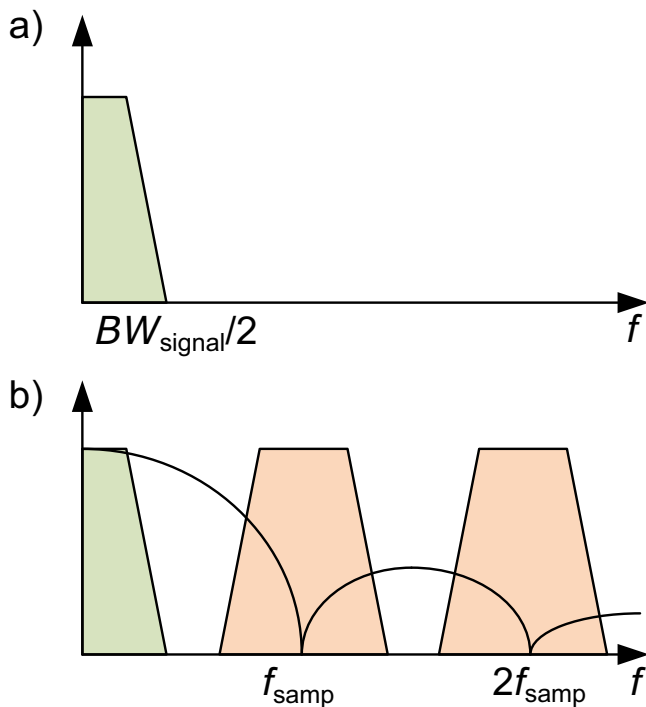


Fig. 6 | Spectral effects of signal sampling. Demonstration of the Nyquist sampling criterion in the frequency domain. **a** The spectrum of an arbitrary baseband signal, with its bandwidth indicated. **b** The spectrum after sampling, showing the replication of the signal at integer multiples of the sampling frequency. The sinc-shape envelope is introduced by zero-order hold sampling.

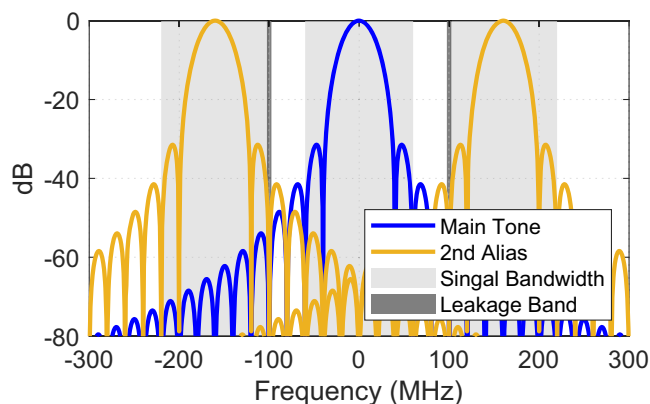


Fig. 7 | Sampling an adiabatic pulse. An example of how Eq. (21) determines the minimum sampling frequency for a raised-cosine pulse with a duration of 50 ns. With an infidelity threshold $F_{th} = 0.01\%$, BW_{signal} is defined as the frequency at which the spectral power drops by $10 \cdot \log_{10}(F_{th}) = -40$ dB with respect to the main lobe, leading to $BW_{\text{signal}} = 120$ MHz. Since higher-order aliases are further from the baseband, it is sufficient to examine only the second alias. With Eq. (21), the minimum sampling frequency of 160 MHz ensures that the second alias lies just beyond the oscillation frequency, thereby avoiding spectral overlap. Sampling at frequencies above this threshold further separates the alias from ΔE_Z , keeping the sampling-induced leakage below the target infidelity.

Consequently, although DLR performs slightly worse than DRAG, its straightforward time-domain implementation and symmetric spectrum make it more suitable for specific applications.

Nyquist criterion on adiabatic sampling

In practice, adiabatic pulses for J will be generated through the use of signal generators, such as arbitrary wave generators (AWGs), that reproduce a

time-sampled version of $J(t)$. However, sampling at a rate lower than the Nyquist rate can increase leakage or excitations due to the appearance of unwanted oscillations at frequencies f_{osc} . While choosing a sampling frequency much higher than the minimal requirement could be a safe option, it significantly increases the power and cost of the electronic interface, thus limiting the scalability and the feasibility of such a solution. Therefore, determining the minimum sampling rate required to preserve adiabaticity of the evolution is of utmost importance.

We start by investigating the effect of finite sampling on the spectrum $S(f)$ of the DLR pulse. After sampling a signal in the time domain, aliases of the signal are created at integer multiples of the sampling frequency, as shown in Fig. 6. Furthermore, assuming that the analog electronics reconstructing the pulse adopt a zero-order hold, the sinc(f) envelope modulates the entire spectrum. The frequency spectrum of $S(f)$, sampled with the sampling frequency f_{samp} can then be written as^{91–93}.

$$S_{\text{samp}}(f) = \frac{\sin(\pi f / f_{\text{samp}})}{\pi f / f_{\text{samp}}} \sum_{k=-\infty}^{\infty} S(f - f_{\text{samp}} \cdot k). \quad (20)$$

As explained in the Methods section, the fidelity of the resulting operation can be estimated by estimating $S_{\text{samp}}(-\Delta E_Z)$. To simplify the expressions, we focus only on the dominant terms $k = 0, -1$ and ignore the remainder. The contribution of a limited sample rate to the infidelity can be calculated by computing the amplitude of the second alias at $-\Delta E_Z$. To keep the infidelity below a certain threshold, the sampling frequency should be chosen as

$$f_{\text{samp}} \geq \Delta E_Z + BW_{\text{signal}}/2. \quad (21)$$

Here BW_{signal} is the bandwidth of the original signal $S(f)$ that must be chosen based on the target infidelity threshold, as shown in the example in Fig. 7. However, this analysis does not yet include the effect of the sinc function over the spectrum, which can further reduce the effect of sampling on the infidelity. Additionally, the usage of anti-alias filters or any bandwidth limitation can further reduce the spectral leakage and improve the gate fidelity. However, the bandwidth of the voltage signal must be high enough to faithfully reproduce the desired exchange pulse. Due to the approximately exponential relation between voltage and exchange ($J(t) \propto \exp(2 \cdot \alpha \cdot V_B(t))$), with α being a lever arm, the relevant bandwidth requirement can be determined via numerical inversion (see the Impact of noise section).

To further investigate the limits in the choice of the sampling frequency, Fig. 8 shows the fidelity of the CPHASE operation generated with a 50 ns raised-cosine pulse (which has a $BW_{\text{signal}} = 120$ MHz for $F_{th} = 0.01\%$) for different sampling frequencies and ΔE_Z . High-fidelity operations are guaranteed between two boundary conditions: adiabatic leakage (Eq. (10)) and minimum sampling frequency (Eq. (21)). Outside this region, high-fidelity operations can still be observed under specific conditions. Firstly, the $J/\Delta E_Z$ ratio becomes smaller with large ΔE_Z , thus lowering the SWAP oscillation amplitude. Secondly, when the sampling period matches the oscillation period, the control signal effectively forms a sequence of diabatic pulses, which can also yield high fidelity. This condition is spectrally equivalent to the sinc envelope in Fig. 7, exhibiting a notch at ΔE_Z , thereby suppressing leakage at that frequency. However, these regimes come with practical limitations. Large ΔE_Z can cause strong frequency crowding, thereby constraining addressability and increasing crosstalk in multi-qubit systems. Moreover, the fidelity of diabatic operations is highly sensitive to precise timing and parameter calibration, making them less robust in experimental settings.

Impact of noise

So far, we have neglected the impact of decoherence due to omnipresent charge noise superimposed on the electrode voltage. The exchange

interaction (J) as a function of the barrier voltage (V_B) is well approximated as^{23,24}

$$J = J_{\text{off}} \cdot e^{2 \cdot \alpha_{\text{leverarm}} \cdot V_B}, \quad (22)$$

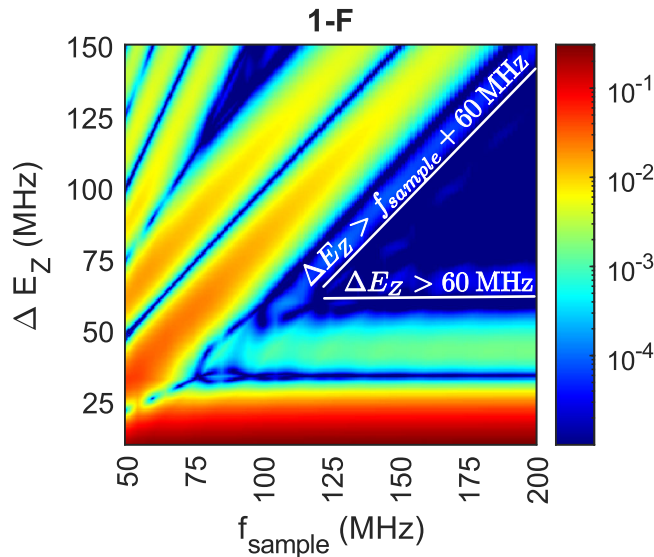


Fig. 8 | Impact of sampling frequency on the fidelity of the CZ gate. Fidelity of CZ operation generated with a 50 ns raised-cosine signal in a two-qubit system with different ΔE_Z and f_{sample} . In the figure, it can be seen that infidelity is preserved when conditions in Eq. (10) and Eq. (21) are met. To increase the visibility of the region, the lowest infidelity is limited by 10^{-5} . The plot with raw data can be found in Fig. 13.

where α_{leverarm} is the lever arm and J_{off} is the residual exchange. In our simulations, the pulses are calibrated such that J_{off} does not introduce over-rotation errors from the residual phase accumulation for long adiabatic pulses as in ref. 35. Additionally, in experimental settings, residual exchange during idle periods can be mitigated using decoupling protocols⁹⁰, or by physically increasing the separation between the two qubits to lower J_{off} ⁹⁴.

The strong nonlinearity between $J - V_B$ implies that analytical assessment of nonidealities — such as quantization and voltage noise — becomes intractable when working directly in the voltage domain. This is crucial to address the spin flip-flop (ϵ) and phase accumulation errors ($\delta\Phi_J$) in Eq. (12). For spin flip-flop errors ϵ , the amplitude of the spectrum of the exchange pulse is crucial. To quantify such effects, we consider small fluctuations in barrier voltage caused by the aforementioned non-idealities. For small voltage fluctuations, we can linearize (22) around the operating point, resulting in

$$\frac{dJ}{dV_B} = 2 \cdot \alpha_{\text{leverarm}} \cdot J_{\text{off}} \cdot e^{2 \cdot \alpha_{\text{leverarm}} \cdot V_{B,OP}} \quad (23)$$

$$dJ = 2 \cdot \alpha_{\text{leverarm}} \cdot J_{OP} \cdot dV_B \quad (24)$$

where J_{OP} and $V_{B,OP}$ are the operating point for the exchange interaction and the barrier voltage, respectively.

To show the validity of this approach, Fig. 9 compares the spectrum of the voltage and exchange pulses in the presence of $1/f$ voltage noise. In the following section, we will use the DLR-average pulse to compare DLR with other pulses. As shown in the Derivations of Noise Impacts on Fidelity in the Methods section, the noise level in the voltage domain can be converted to the exchange-pulse domain using the small-signal approximation. Then ϵ term in Eq. (10) can then be approximated as (see the Derivations of Noise

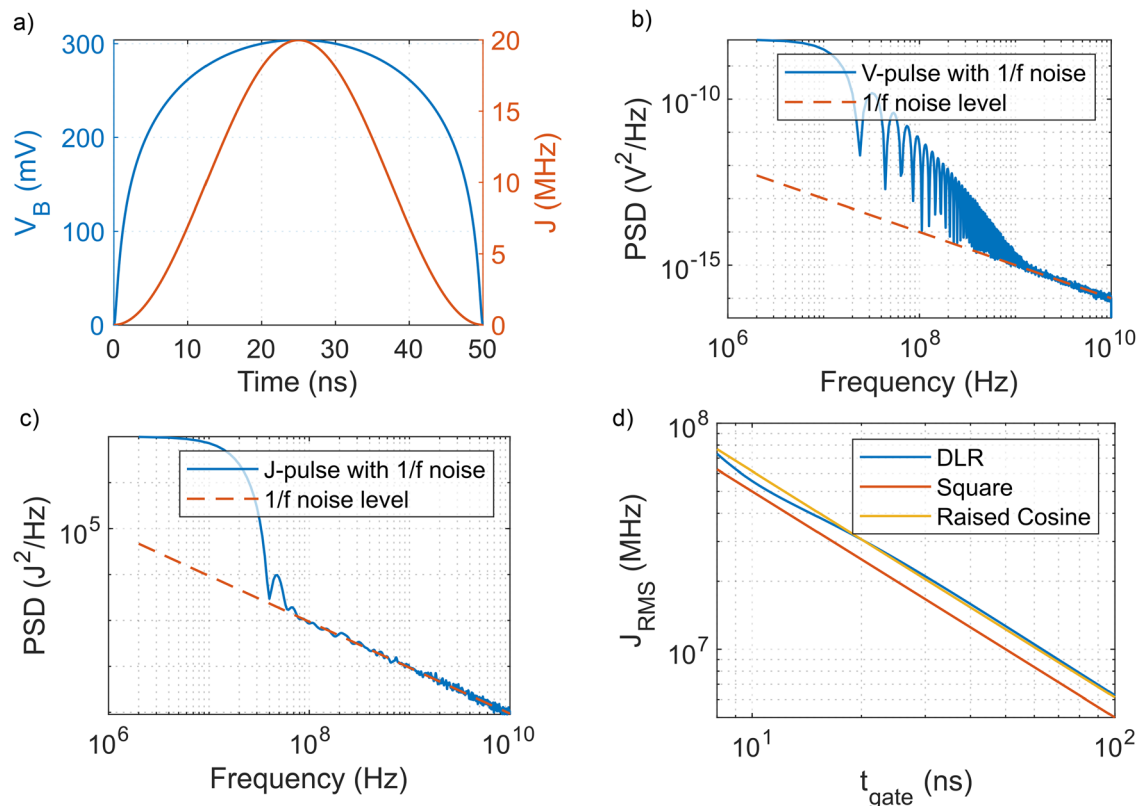


Fig. 9 | Noise analysis of exchange pulses. **a** Voltage pulse V_B generates a raised-cosine exchange pulse of 50 ns with $\alpha_{\text{leverarm}} = 12.5$ and $J_{\text{off}} = 10$ kHz. Simulated PSD of a raised-cosine **(b)** voltage and **(c)** exchange pulse under the presence of $1/f$ noise.

The plots show the average power over 100 noise realizations. **d** RMS of the different exchange pulses scales inversely proportionally with the gate time.

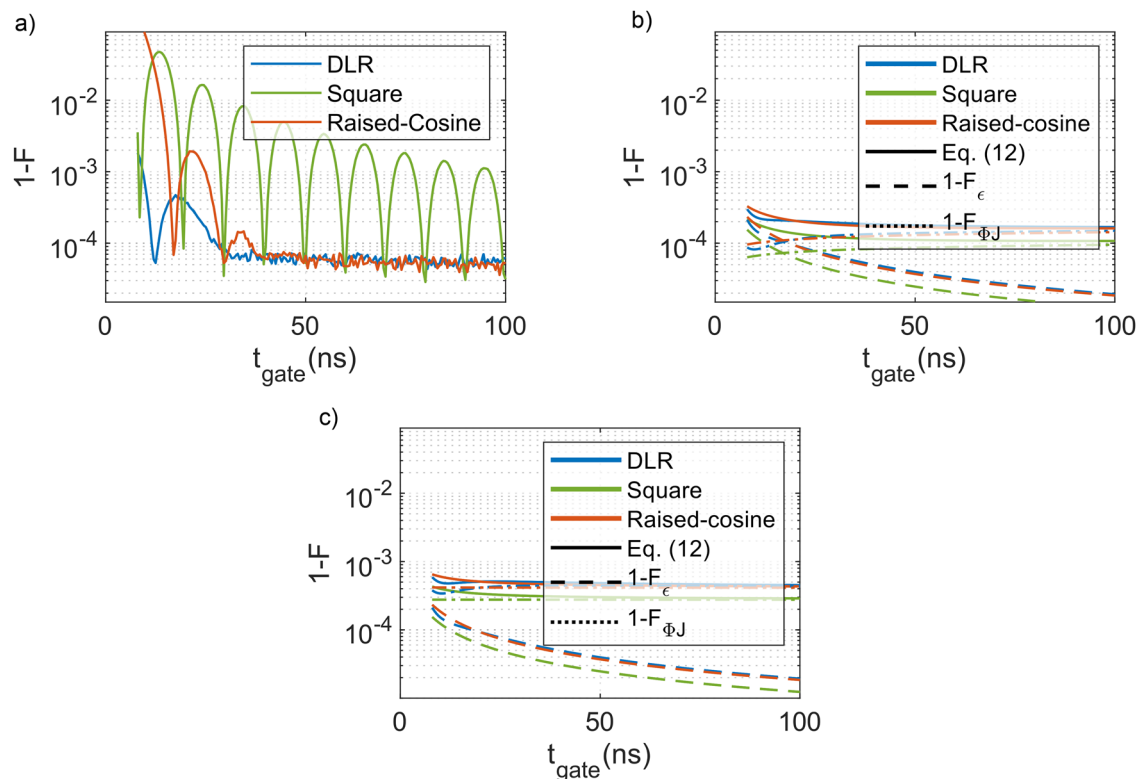


Fig. 10 | Impact of noise on gate fidelity. **a** Simulated DLR, square, and raised-cosine pulses in the presence of $1/f$ noise with $P_1 = 0.1 \mu\text{V}^2/\text{Hz}$ applied to a 50 ns voltage pulse with $\alpha_{\text{leverarm}} = 12.5$ and $J_{\text{off}} = 10 \text{ kHz}$, and 100 repetitions. **b** The fidelity limitation imposed by $1/f$ noise is computed analytically using Eq. (12), where ϵ and

Φ_J are estimated using Eq. (24). To illustrate the individual contributions of ϵ ($1-F_{\epsilon}$) and Φ_J ($1-F_{\Phi_J}$), each term is also substituted independently into Eq. (12). **c** To show the impact of noise for a given measurement time, analysis in (b) is repeated where $\sigma_{1/f} = \sqrt{P_1 \cdot \ln(t_{\text{meas}} \cdot f_c)}$ ($t_{\text{meas}} = 0.1 \text{ s}$).

Impacts on Fidelity in the Methods section):

$$\epsilon \approx \pi^2 \frac{P_1 \cdot t_{\text{gate}}}{\Delta E_Z} \cdot (2 \cdot \alpha_{\text{leverarm}} \cdot J_{\text{RMS}})^2. \quad (25)$$

where P_1 (V^2/Hz) is the power spectral density (PSD) of flicker noise at 1 Hz and t_{gate} is the pulse duration. Figure 9.d shows the J_{RMS} for different exchange-pulse shapes, which determines the impact of noise according to (25). Since all pulses satisfy $\int_0^{t_{\text{gate}}} 2\pi \cdot J(t)dt = \pi$, their RMS values remain close to one another.

Noise will also impact the fidelity via the phase accumulation errors (see the Derivations of Noise Impacts on Fidelity in the Methods section)

$$\delta\Phi_J = J_{\text{RMS}} \cdot 2 \cdot \alpha_{\text{leverarm}} \cdot t_{\text{gate}} \cdot \sigma_n. \quad (26)$$

For $1/f$ noise, σ_n can be found as $\sigma_{1/f} = \sqrt{P_1 \cdot \ln(t_{\text{gate}} \cdot f_c)}$ where f_c is the corner frequency of the $1/f$ noise.

Figure 10 compares the numerical simulation results with our analytical approximation. In Fig. 10b, the fidelity is computed using the analytical approximation, considering only $1/f$ noise, therefore neglecting any flip-flop error. The results indicate that ϵ and Φ_J have opposite trends as a function of the gate duration t_{gate} . On one hand, shorter pulses increase the error contribution due to ϵ , since J_{RMS} decreases with t_{gate} . On the other hand, longer gate times increase the accumulated $1/f$ noise $\sigma_{1/f}$, thereby increasing Φ_J . However, if we consider the $1/f$ noise not only during the gate operation but also during measurement, as in Fig. 10c, its impact becomes approximately gate-time independent.

Discussion

This work introduces the time-domain pulse-shaping strategy Delayed Leakage Reduction (DLR), which enables high-fidelity, ultra-fast adiabatic operations in systems constrained to baseband control, such as two-qubit

gates for spin qubits. By analytically linking the gate fidelity to the spectral properties of the control signal, we reformulate the fidelity optimization problem as a spectral design problem. The proposed DLR technique introduces targeted spectral notches through time-delayed pulse replicas, effectively suppressing transition errors at critical frequencies without requiring IQ modulation. Simulation results demonstrate that DLR achieves fidelity exceeding 99.9% within gate durations as short as 9.4 ns, while only marginally increasing the hardware sampling requirements. Additionally, we have shown that the impact of charge noise is approximately constant with gate time, thus not a limiting factor for our short gate durations.

While DLR is primarily designed to suppress spectral leakage at a single transition frequency, it can be extended to multi-level systems by applying it sequentially. In this approach, each DLR application is tuned to place a spectral notch at a different target frequency. However, since each additional DLR pulse shortens the effective gate duration, the spectrum of the resulting signal broadens, and previously placed notches may shift. This shift can be compensated by ensuring that the product of the notch frequency and the gate duration remains constant, i.e., $f_{\text{notch},1} \cdot t_{\text{gate},1} = f_{\text{notch},2} \cdot t_{\text{gate},2}$. This relation allows earlier DLR stages to be adjusted accordingly, preserving the intended spectral suppression across multiple transitions.

The constraint on the minimum sampling frequency for adiabatic control gives a theoretical bound that ensures spectral aliases remain outside of the region that causes spin flip-flop errors. This informed design method prevents over-design of the control hardware, which can cost unnecessary power consumption and system complexity, particularly for large-scale systems. Given the limited power budget available in a cryogenic environment, avoiding an over-designed cryogenic controller is also an important consideration. Thanks to this method, both room temperature and cryogenic electronics can be designed with this information, and help to keep adjusting the power and complexity while suppressing leakage. To further suppress the spectral leakage, a low-pass filter can be used to attenuate alias components that overlap with the leakage band. Moreover, employing first-

order-hold reconstruction instead of zero-order-hold introduces linear interpolation between sampling points, resulting in a spectral envelope proportional to $\text{sinc}^2(\pi f/f_{\text{sample}})$ rather than $\text{sinc}(\pi f/f_{\text{sample}})$ ^{91,92}. These modifications significantly reduce the amplitude of aliasing components, thereby improving gate fidelity under finite sampling constraints.

Thanks to its conceptual generality and compatibility with baseband control, the proposed DLR technique is not only applicable to spin qubits but can also be extended to other qubit platforms and adiabatic processes. By enabling ultra-fast, high-fidelity gate operations that surpass the speed of current state-of-the-art two-qubit spin-qubit gates, DLR offers a versatile and scalable tool for advancing the development of next-generation quantum processors.

Methods

Window functions for spectral shaping

In this appendix, we summarize the mathematical definitions of the three window functions employed for shaping the exchange control envelope $J(t)$, Raised-Cosine, Kaiser, and Hamming windows. These functions are widely used in signal processing to suppress spectral leakage and to tailor the energy distribution of time-domain control signals.

Each window function is defined over the interval $t \in [0, t_{\text{gate}}]$, where t_{gate} denotes the total gate duration. Outside this interval, i.e., for $t < 0$ or $t > t_{\text{gate}}$, the window function is identically zero:

$$w(t) = 0, \text{ for } t < 0 \text{ or } t > t_{\text{gate}}. \quad (27)$$

Raised-Cosine window. The Raised-Cosine window is defined as

$$w_{\text{RC}}(t) = \frac{1}{2} \left[1 - \cos \left(2\pi \frac{t}{t_{\text{gate}}} \right) \right], \quad (28)$$

for $t \in [0, t_{\text{gate}}]$. This window exhibits smooth onset and termination, effectively suppressing high-frequency components in the spectral domain.

Kaiser window. The Kaiser window introduces a tunable shape parameter β that controls the trade-off between main-lobe width and side-lobe suppression. It is defined as

$$w_{\text{K}}(t) = \frac{I_0 \left(\pi \beta \sqrt{1 - \left(\frac{2t}{t_{\text{gate}}} - 1 \right)^2} \right)}{I_0(\pi \beta)}, \quad (29)$$

where $I_0(\cdot)$ denotes the zeroth-order modified Bessel function of the first kind. Larger values of β yield stronger side-lobe suppression at the cost of broader main lobes.

Hamming window. The Hamming window is a fixed-shape window function defined as

$$w_{\text{H}}(t) = 0.54 - 0.46 \cos \left(2\pi \frac{t}{t_{\text{gate}}} \right), \quad (30)$$

for $t \in [0, t_{\text{gate}}]$. Compared to the Raised-Cosine window, the Hamming window offers an improved main-lobe width at the expense of slightly higher side-lobe levels.

All window functions are normalized such that

$$\int_0^{t_{\text{gate}}} w(t) dt = t_{\text{gate}}, \quad (31)$$

ensuring consistent energy scaling across different pulse shapes.

Numerical simulations

Simulation scripts are developed using the SPINE⁹⁵ toolkit. To solve the time-dependent evolution of the qubit system, the Schrödinger equation

$$i\hbar \frac{d}{dt} |\psi(t)\rangle = H(|\psi(t)\rangle) \quad (32)$$

is solved iteratively as

$$U(t + \Delta t) = e^{-\frac{i}{\hbar} H(t + \Delta t) \Delta t} U(t), \quad (33)$$

where the quantum operator U is calculated with small time steps Δt by assuming that the change of H between steps is negligible. The average gate fidelity is given by

$$F(U_{\text{ideal}}, U_{\text{sim}}) = \frac{\text{Tr}[\mathcal{U}_{\text{ideal}}^\dagger \mathcal{U}_{\text{sim}}] + d}{(d+1)d} \quad (34)$$

$$= \frac{\text{Tr}[\mathcal{U}_{\text{ideal}}^\dagger \mathcal{U}_{\text{sim}}]^2 + d}{(d+1)d}, \quad (35)$$

where the superoperators are calculated as

$$\mathcal{U} = U^* \otimes U. \quad (36)$$

Explicit expressions of errors

Following³⁵, the total imperfect time evolution of an adiabatic quantum gate can be written as

$$U_{\text{gate}}(t) = \mathcal{D}^\dagger(t) U_{\text{ad}}(t) \mathcal{E}(t) \mathcal{D}(0) \quad (37)$$

for all times t in the interval $[0, t_{\text{gate}}]$. Here, U_{ad} is the desired operation to perform, $\mathcal{E}$ is the error, and $\mathcal{D}$ is the unitary transformation from the computational basis into the instantaneous eigenbasis. Remark that for our case $\mathcal{D}(0) = \mathcal{D}(t_{\text{gate}}) = 1$, thus we will ignore the terms in the following. For adiabatic dynamics

$$U_{\text{ad}}(t) = \exp \left[-\frac{i}{\hbar} \int_0^t H_{\text{ad}}(t') dt' \right] \quad (38)$$

$$= \sum_m e^{i\alpha_m(t)} |m(t)\rangle \langle m(t)|. \quad (39)$$

where we used the spectral decomposition, introduced the phases $\alpha_m = -\frac{i}{\hbar} \int_0^t \epsilon_m(t') dt'$, eigenvectors $|m\rangle$, and eigenenergies $\epsilon_m(t)$. The error term can be expressed as a Magnus expansion³⁵

$$\mathcal{E}(t) = \exp \left[-\frac{i}{\hbar} \sum_{k=1}^{\infty} H_{\text{Mk}} \right], \quad (40)$$

where the lowest order term is given by integration over the error Hamiltonian

$$H_{\text{M1}} = \frac{1}{\hbar} \int_0^t H_{\text{error}}(t') dt'. \quad (41)$$

The error Hamiltonian is given by³⁵

$$H_{\text{error}} = \sum_{m_1, m_2} e^{i(\alpha_{m_2} - \alpha_{m_1})} \langle m_1(t) | H_{\text{total}} - H_{\text{ad}} | m_2(t) \rangle \quad (42)$$

and finally the error rates for the transition $\mathcal{O}_{m_1, m_2} = |m_2\rangle\langle m_1|$ by

$$\epsilon_{m_1, m_2} = \left| \text{Tr} \left(\mathcal{E} \mathcal{O}_{m_1, m_2} - \mathcal{O}_{m_1, m_2} \right) \right|^2 \approx \left| \frac{1}{\hbar} \int_0^{t_{\text{gate}}} \langle m_1(t) | H_{\text{total}}(t) - H_{\text{ad}}(t) | m_2(t) \rangle e^{i(\alpha_{m_2}(t) - \alpha_{m_1}(t))} dt \right|^2 \quad (43)$$

$$-H_{\text{ad}}(t) | m_2(t) \rangle e^{i(\alpha_{m_2}(t) - \alpha_{m_1}(t))} \quad (44)$$

For the coupling between two spin qubits, the Hamiltonians (in frequency units) read

$$H_{\text{ad}} = \frac{\hbar}{2} \begin{pmatrix} 2E_z & 0 & 0 & 0 \\ 0 & f_{\text{osc}} - J & 0 & 0 \\ 0 & 0 & -f_{\text{osc}} - J & 0 \\ 0 & 0 & 0 & -2E_z \end{pmatrix} \quad (45)$$

$$H_{\text{total}} - H_{\text{ad}} = \frac{\hbar}{2} \begin{pmatrix} 0 & 0 & 0 & 0 \\ 0 & 0 & \frac{i}{2\pi} \frac{\dot{J}(t) \Delta E_z}{f_{\text{osc}}^2} & 0 \\ 0 & -\frac{i}{2\pi} \frac{\dot{J}(t) \Delta E_z}{f_{\text{osc}}^2} & 0 & 0 \\ 0 & 0 & 0 & 0 \end{pmatrix}, \quad (46)$$

where $f_{\text{osc}} = \sqrt{J^2 + \Delta E_z^2}$ is the eigenfrequency and $\dot{J}$ the time derivative of the exchange frequency. The single non-adiabatic error is then given by³⁵

$$\epsilon = |\text{Tr}(\mathcal{E} | 01 \rangle \langle 10 |)|^2 = \left| \frac{\Delta E_z}{2} \int_0^{t_{\text{gate}}} \frac{\dot{J}}{f_{\text{osc}}^2} e^{-2\pi i \int_0^t f_{\text{osc}}(t') dt'} dt \right|^2. \quad (47)$$

The expression from the main text, Eq. (8) can be recovered by using integration by parts assuming $J(0) = J(t_{\text{gate}}) = 0$

$$\epsilon = \left| \pi \int_0^{t_{\text{gate}}} f_{\text{osc}}(t) \arctan \left(\frac{J(t)}{\Delta E_z} \right) e^{-2\pi i \int_0^t f_{\text{osc}}(t') dt'} dt \right|^2 \quad (48)$$

$$\approx \left| \pi \int_0^{t_{\text{gate}}} J(t) e^{-2\pi i \int_0^t f_{\text{osc}}(t') dt'} dt \right|^2. \quad (49)$$

In the last line, we used $f_{\text{osc}}(t) \arctan \left(\frac{J(t)}{\Delta E_z} \right) = J(t) + \mathcal{O} \left(\frac{J(t)^3}{\Delta E_z^3} \right)$.

Second-order Magnus expansion correction

The correction from second-order Magnus expansion is given by Eq. (55) where the explicit expressions read

$$\delta = 4\pi^2 \text{Im} \left[\int_0^{t_{\text{gate}}} dt \int_0^t dt' g(t) g^*(t') e^{-i2\pi(\alpha(t) - \alpha(t'))} \right] \quad (50)$$

$$g(\tau) = -\frac{\Delta E_z \dot{J}(\tau)}{2\pi f_{\text{osc}}^2(\tau)} \quad (51)$$

$$\alpha(\tau) = \int_0^\tau dt' f_{\text{osc}}(t'). \quad (52)$$

$$H_{M2} = \frac{i}{2\hbar^2} \int_0^{t_{\text{gate}}} dt \int_0^t dt' [H_{\text{error}}(t), H_{\text{error}}(t')] = i2\pi^2 \int_0^{t_{\text{gate}}} dt \int_0^t dt' (g(t) g^*(t') e^{-i2\pi(\alpha(t) - \alpha(t'))}) \quad (53)$$

$$-g(t') g^*(t) e^{i2\pi(\alpha(t) - \alpha(t'))} \begin{pmatrix} 0 & 0 & 0 & 0 \\ 0 & 1 & 0 & 0 \\ 0 & 0 & -1 & 0 \\ 0 & 0 & 0 & 0 \end{pmatrix} \quad (54)$$

$$= \begin{pmatrix} 0 & 0 & 0 & 0 \\ 0 & \delta & 0 & 0 \\ 0 & 0 & -\delta & 0 \\ 0 & 0 & 0 & 0 \end{pmatrix}, \quad (55)$$

Phase calibration, final gate time, and average gate fidelity

The average gate fidelity of the total gate evolution $U_{\text{gate}}(t_{\text{gate}})$ with respect to the targeted quantum gate $U_{\text{ideal}} = Z_1(-\theta_1)Z_2(-\theta_2)\text{CZ}$ is then given by

$$F(U_{\text{ideal}}, U_{\text{gate}}) = F[U_{\text{ideal}}, \mathcal{D}^\dagger(t_{\text{gate}}) U_{\text{ad}}(t_{\text{gate}}) \mathcal{E} \mathcal{D}(0)], \quad (56)$$

$$\approx F[U_{\text{ideal}}, U_{\text{ad}}(t_{\text{gate}}) \exp \left(-\frac{i}{\hbar} (H_{M1} + H_{M2}) \right)], \quad (57)$$

$$= F[\text{CZ}, Z_1(\theta_1)Z_2(\theta_2)U_{\text{ad}}(t_{\text{gate}}) \exp \left(-\frac{i}{\hbar} (H_{M1} + H_{M2}) \right)], \quad (58)$$

where we used the approximation and simplifications made in the Explicit expressions of errors in the Methods section, and the properties of the trace. We note that the two single-qubit phase gates $Z_1(-\theta_1)$ and $Z_2(-\theta_2)$ can be done digitally by using the virtual gate method. Using further $\text{CZ} = \text{CZ}'$, the fidelity can be expressed as

$$F(U_{\text{ideal}}, U_{\text{gate}}) = \quad (59)$$

$$F[1, \exp(-i\Phi_f Z_1 Z_2) \exp(-i\delta\Phi(Z_1 + Z_2)/2) \exp(-i\delta\theta(Z_1 - Z_2)/2) \exp(-\frac{i}{\hbar} (H_{M1} + H_{M2}))], \quad (60)$$

where $\delta\Phi_f = \pi - 2\pi \int_0^{t_{\text{gate}}} J(t) dt$ is the exchange miscalibration, $\delta\Phi = (\theta_1 + \theta_2 - \pi) - 2\pi E_z t_{\text{gate}}$ the miscalibration of the rotating frame, and $\delta\theta =$

$(\theta_1 - \theta_2) - 2\pi \int_0^{t_{\text{gate}}} \sqrt{\Delta E_z^2 + J^2(t)} dt$ the miscalibrated phase corrections.

Assuming small errors and using the explicit error Hamiltonians (41) and (55), we find the following compact expressions for the infidelity by expanding Eq. (60) up to second order

$$1 - F(U_{\text{ideal}}, U_{\text{gate}}) \approx \frac{2}{5} [\epsilon + 2\Phi_f^2 + (\delta\theta + \delta)^2 + \delta\Phi^2]. \quad (61)$$

Setting $\delta = 0$, we recover formula (12) from the main text.

Numerical phase calibration

While $\int_0^{t_{\text{gate}}} J(t) dt$ generates phase accumulation of π , this operation also shifts the oscillation frequencies of the qubits during the operation. This can be

calibrated by a single-qubit $Z_{1,2}$, explained in the Phase calibration, final gate time, and average gate fidelity in the Methods section. The rotation angle of this calibration can be calculated as

$$\theta_{1,2} = \Delta E_Z \mp \sqrt{\Delta E_Z^2 + J^2}. \quad (62)$$

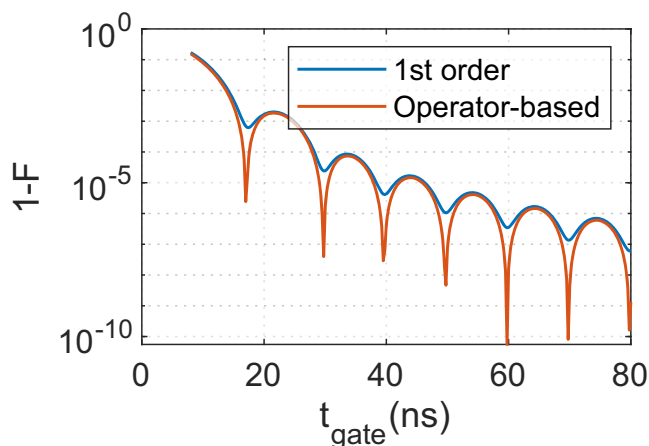


Fig. 11 | Phase calibration. Infidelity ($1-F$) of the CPHASE gate plotted against gate time t_{gate} , comparing the 1st-order approximation (blue) and the operator-based approach (orange).

Then, $Z_1(\theta_1)$ and $Z_2(\theta_2)$ can be applied to qubits. However, during the simulations, we realized that Eq. (62) doesn't give the best results. To improve the accuracy, a higher-order Magnus expansion can be used.

While numerically solving these equations may not be the most straightforward method, θ is derived from the evolution operator. The unitary propagator in a rotating reference frame, U , is determined as described in Numerical Simulations in the Methods section. Subsequently, the angles $\theta_{1,2}$ are calculated as $\theta_1 = \pi + \angle U_{11} - \angle U_{22}$ and $\theta_2 = \pi + \angle U_{44} - \angle U_{33}$, where $U_{\text{row, column}}$ denotes the matrix elements of U . Consequently, the difference in angle between the states $|0\rangle$ and $|1\rangle$ for a single qubit is π .

In Fig. 11, a comparison of Eq. (62) and the operator-based method. It can be seen that the accuracy of the CPHASE is improved, especially on an integer multiple of the SWAP oscillation periods (≈ 10 ns).

Comparison with DRAG

In the DRAG protocol, the spectral notches are constructed using IQ control. We emulate this here by considering complex exchange signals J . In particular, we modify the exchange Hamiltonian in Eq. (2) with

$$H_{\text{exc}} = \begin{pmatrix} E_z & 0 & 0 & 0 \\ 0 & (-\Delta E_Z - |J|)/2 & \bar{J}/2 & 0 \\ 0 & J/2 & (\Delta E_Z - |J|)/2 & 0 \\ 0 & 0 & 0 & -E_z \end{pmatrix} \quad (63)$$

Additional data of DLR numerical optimization

Due to the initial discontinuity present in the Hamming and Kaiser window functions, Eq. (10) no longer holds. Consequently, we resort to a direct

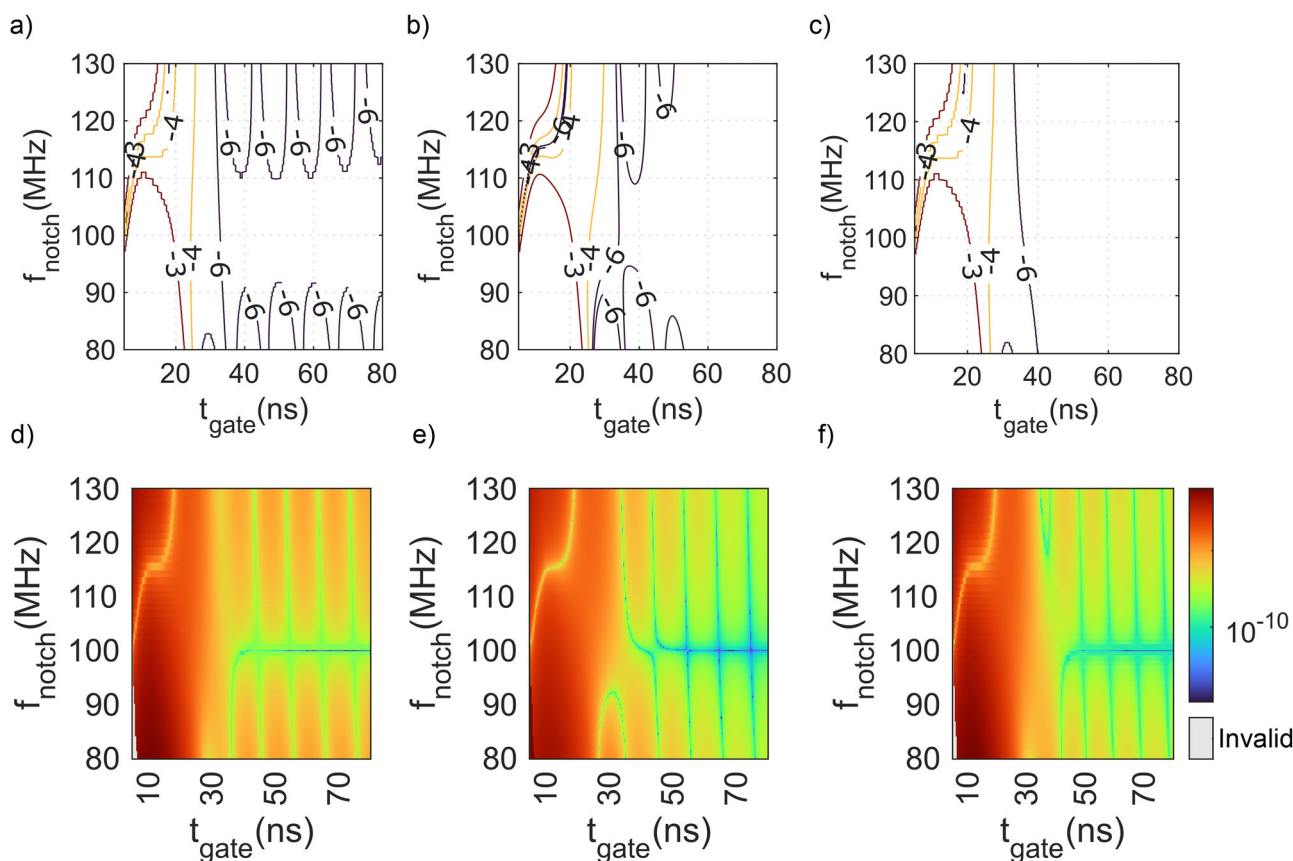


Fig. 12 | DLR with different windows and gate durations. The notch created on the frequency spectrum of the signal can be chosen arbitrarily. To calibrate the DLR numerically, a range of notch frequencies is tested for different gate times. Here, infidelities are shown for **a** raised-cosine, **b** Hamming, and **c** Kaiser windows. The

corresponding fidelity boundaries at 0.1%, 0.01%, and 0.0001% ($1-10^x$ where $x = -3, -4, -6$) are shown in **d-f**, providing a visual reference for the performance thresholds across the parameter space.

numerical optimization approach, where the delay parameter t_d in the DLR protocol is swept between 3.84 ns and 6.25 ns, corresponding to a notch frequency f_{notch} ranging from 130 MHz to 80 MHz, respectively. The resulting gate fidelities for each combination of gate time and f_{notch} are shown in Fig. 12.

For all window functions considered, longer gate durations consistently yield optimal performance when $f_{\text{notch}} \approx \Delta E_Z \approx 100$ MHz. The vertical stripe patterns observed in the fidelity maps arise from non-adiabaticity conditions, where residual oscillations synchronize with the gate duration. For shorter gate times, however, the assumption $f_{\text{osc}} \approx \Delta E_Z$ breaks down, and the optimal choice of f_{notch} becomes window-function dependent.

Beyond 35 ns, as f_{osc} increases, higher values of f_{notch} tend to yield improved performance. However, at sufficiently large f_{osc} at the lower gate time, the main spectral lobe of the pulse begins to overlap with ΔE_Z , leading to increased leakage. In this regime, selecting a lower f_{notch} becomes advantageous. The optimal fidelity achieved for each gate duration is summarized in Fig. 4c.

Additional data of sampling frequency

Simulated fidelity of the CPHASE operation generated by a 50 ns raised-cosine pulse in a two-qubit system, as a function of the qubit frequency separation ΔE_Z and the sampling frequency f_{sample} . The data shown is identical to Fig. 8; however, no clipping filter has been applied to the fidelity values. In Fig. 8, fidelity values below 10^{-5} were clipped for improved visibility, while here the raw fidelity values are retained. This allows for a more

accurate representation of the fidelity landscape, particularly in regions where the error rate falls below the clipping threshold.

Derivations of noise impacts on fidelity

For a given $V_0(J_0)$ and voltage noise (V_n), exchange fluctuations can be approximated using Eq. (24) and yielding:

$$J_n = J_0 \cdot 2 \cdot \alpha_{\text{leverarm}} \cdot V_n \quad (64)$$

Thanks to the modulation theorem $(x(t)y(t) \Rightarrow X(f) \otimes Y(f))$, its spectrum is

$$S_{Jn} = (2 \cdot \alpha_{\text{leverarm}}) \cdot S_{J0} \otimes S_{Vn}. \quad (65)$$

Since the applied pulse has negligible spectral components at ΔE_Z to achieve high fidelity, we can approximate the spectrum S_{J0} as a Dirac pulse centered at DC with associated power equal to the power of the pulse.

Thus, $\epsilon \approx |\pi S(-\Delta E_Z)|^2$ (Eq. (10)) can be found with the Fourier spectrum of V_B ,

$$\epsilon \approx |\pi S_{Vn}(-\Delta E_Z) \cdot 2 \cdot \alpha_{\text{leverarm}} \cdot J_{\text{RMS}}|^2. \quad (66)$$

For $1/f$ voltage noise with power spectral density S_1 at 1 Hz, it can be approximated as

$$\epsilon \approx \pi^2 \frac{S_1^2}{\Delta E_Z} \cdot (2 \cdot \alpha_{\text{leverarm}} \cdot J_{\text{RMS}})^2. \quad (67)$$

The phase accumulation errors can be found as

$$\delta\Phi_J = \pi - 2\pi \int_0^{t_{\text{gate}}} J(t) dt \quad (68)$$

$$= \int_0^{t_{\text{gate}}} dJ(t) dt \quad (69)$$

$$\approx J_{\text{RMS}} \cdot 2 \cdot \alpha_{\text{leverarm}} \cdot \int_0^{t_{\text{gate}}} dV_B(t) \cdot dt. \quad (70)$$

For white noise, the $\int_0^{t_{\text{gate}}} dV_B(t) \cdot dt$ can be approximated using its constant power spectral density, giving

$$\int_0^T V_{\text{white}}(t) dt \sim \mathcal{N}(0, P_0 T), \quad (71)$$

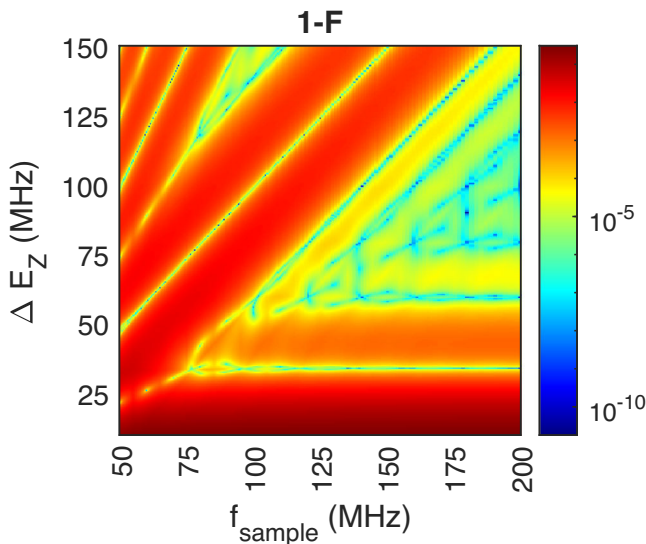
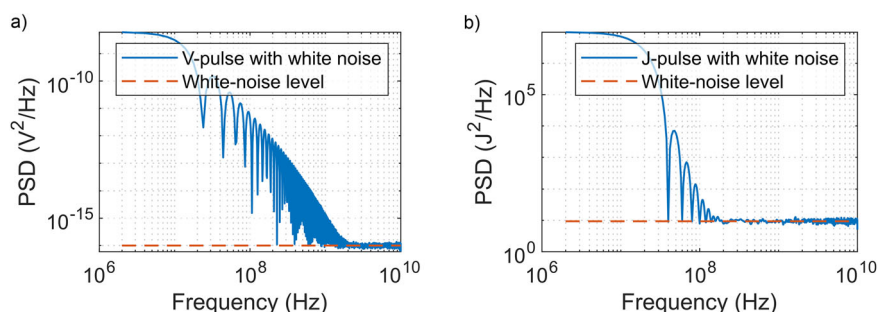


Fig. 13 | Raw data of the impact of sampling frequency on the fidelity of the CZ gate. Computed fidelity of the CPHASE gate generated with a 50 ns Raised-Cosine signal in the two-qubit system with varying ΔE_Z and f_{sample} .

Fig. 14 | Impact of white noise on exchange pulse.

Simulated PSD of a raised-cosine **a** voltage and **b** exchange pulse under the presence of white noise on 50 ns voltage pulse with $\alpha_{\text{leverarm}} = 12.5$ and $J_{\text{off}} = 10$ kHz. The plots show the average power over 100 noise realizations. The exchange pulse is generated numerically, then its spectrum is calculated. White noise level on the exchange spectrum is analytically calculated by following Eq. (24), where J_{OP} is chosen as the RMS of the raised-cosine exchange pulse (12.25 MHz).



where P_0 is the one-sided white-noise spectral density, while the $1/f$ component grows approximately as

$$\text{Var} \left(\int_0^T V_{1/f}(t) dt \right) \approx P_1 \cdot \ln(t_{\text{gate}} \cdot f_c) \quad (72)$$

where f_c is the cut-off frequency, and P_1 is the PSD at 1 Hz.

White and quantization noise on the voltage

In Figs. 13, 14, we show how a white noise level on the voltage can be mapped to the exchange spectrum. Similarly, the signal quantization noise from the signal generators can also be approximated as white noise⁹³ and similarly mapped to the exchange spectrum. Then, the requirements on the exchange spectrum can be translated into requirements on the voltage spectrum, such as its purity noise level. In turn, using the standard expression linking the Signal-to-Noise and Distortion Ratio (SNDR)⁹³ and the effective number of bits (ENOB) for the digital-to-analog converter, i.e., the AWG (Arbitrary Waveform Generator) adopted in the experiment:

$$\text{SNDR} = \text{ENOB} \cdot 6.02 \text{ dB} + 1.76 \text{ dB}, \quad (73)$$

the AWG required resolution can be derived.

Data availability

The datasets generated during the current study are publicly available at <https://doi.org/10.5281/zenodo.16707554>⁹⁶.

Code availability

The code used throughout this manuscript is available at <https://doi.org/10.5281/zenodo.16707554>⁹⁶. The repository includes scripts for reproducing all figures and numerical results presented in this work.

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Author contributions

I.P. developed the theoretical model, performed the simulations, and performed the analysis. M.R.-R. supervised and contributed to the theoretical framework. F.S. and R.W.J.O. supervised the project. I.P. has written the manuscript with the input from F.S., R.W.J.O., and M.R.-R.. All authors reviewed the manuscript.

Competing interests

The authors declare that a patent application related to this work has been filed. The inventors are Ilker Polat, Ramon W. J. Overwater, Fabio Sebastiano, Maximilian Rimbach-Russ, and Masoud Babaie. The Dutch patent application has been submitted under application number NL2040125 and is currently pending. The patent covers aspects of the Dynamic Leakage Reduction (DLR) technique presented in this manuscript.

Additional information

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