Dynamics of Hole Singlet-Triplet Qubits with Large $g$-Factor Differences

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The spin-orbit interaction (SOI) allows electrical manipulation of individual spins and has therefore become a key ingredient for the realization of fully electrically controlled spin qubits [1,2]. For electrons in Si it is rather weak and synthetically boosted by means of micromagnets [3,4]. For holes, on the other hand, it is an intrinsic property which allows to perform electron dipole spin resonance (EDSR) measurements [1,2,5–9]. In Ge it is particularly strong leading to Rabi frequencies beyond 100 MHz [7,10,11]. SOI for holes can be linear or cubic in the wave vector $k$, with nanowire qubits favoring the former type while planar qubits the latter [12–14]. The SOI is not only important for single spin but also for singlet-triplet qubits as it causes an intrinsic mixing between the heavy hole (HH) and light hole (LH) bands and thereby locally affects the $g$ factors of the individual spins allowing to drive $S - T_0$ oscillations [15]. In combination with an extrinsic Rashba type SOI caused by the structural inversion asymmetry induced by the heterostructure, it also mixes the $S$ and $T_-$ states contributing therefore to a measurable avoided crossing $\Delta_{ST_-}$.

Here, we investigate this avoided crossing for a double quantum dot (DQD) Ge hole spin system and gain insight into the interplay between SOI and the $g$-factor anisotropy and their consequences on qubit dynamics.

A scanning electron microscope (SEM) image of the device under consideration is depicted in Fig. 1(a) and further details can be found in Ref. [15]. A two-dimensional hole gas is embedded in a Ge/SiGe heterostructure and additional TiPd top gates confine a DQD and a charge sensor (CS). For qubit state selective read out we rely on Pauli spin blockade (PSB) combined with Ohmic reflectometry [16]. Fast detuning pulses are applied to gates LB and RB through an arbitrary waveform generator (AWG) with a pulse-rise time of $\tau_{\text{rise}} \approx 2$ ns. Throughout this work we apply a small magnetic field in a plane perpendicular to the axis connecting the two dots (DQD axis), $B = [B \cos(\theta), 0, B \sin(\theta)]$, where $\theta$ describes the tilt angle from the in-plane direction. We tune the DQD to a charge transition between an effective $(2,0) \leftrightarrow (1,1)$ state, with $(n_L, n_R)$, where $n_L$ ($n_R$) denotes the effective hole number in the left (right) QD [Fig. 1(c)]. The tunnel coupling between the dots is described by $t_C$ while the energy detuning between the $S(2,0)$ and $S(1,1)$ state is parametrized by $\epsilon$. Each QD is characterized by an out-of-plane and an in-plane $g$ factor, $g_\perp$ and $g_\parallel$, respectively. However, the dynamics of singlet-triplet qubits is only

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induced when the system (the complete Hamiltonian $H_{\text{tot}}$) is driven by the Zeeman energy $\mu_B B$ and the separation time $\tau_S$. We apply a rapid pulse of duration $\tau_p = 65$ ns and varying $\epsilon$ in order to extract the individual $g$ factors without the need for EDSR. We fix the magnetic field at $|B| = 2\, \text{mT}$ and observe the oscillations vs $\tau_S$ while rotating $B$. We use a fast pulse ($\tau_p = \tau_{\text{rise}} = 2$ ns) in Fig. 2(a) and a ramped pulse with ramp time $\tau_R = 100$ ns in Fig. 2(c). In both cases we pulse to $\epsilon = 4\, \text{meV}$ for a duration $\tau_R$. From the fast Fourier transform (FFT) in Figs. 2(b) and 2(d) we extract the oscillation frequency $f_{S-T_0}$ (orange dots) and $f_{S-T_-}$ (pink dots). We notice that $f_{S-T_-}$ is very close to $f_{S-T_0}$ and $f_{S-T_0}$ is almost $2\, \text{meV}$, which is in line with the observations in Fig. 1(f). Moreover, in Fig. 2(d) the FFT power vanishes for $\theta \approx 60^\circ$, indicating that the ramp time $\tau_R$ induces a completely diabatic passage over the avoided crossing. This is in line with Fig. 1(e) where we observed a charge sensor in the measurement point as a function of $\epsilon$ and $B$ for $\theta = 90^\circ$, $60^\circ$, and $10^\circ$, respectively. A high phase signal corresponds to a larger triplet return probability. In the out-of-plane direction we observe the expected funnel shape of the $S-T_-$ anticrossing [24]. At $60^\circ$ we similarly observe a typical funnel shape, however, we notice the line to be fainter, which indicates a smaller $\Delta_{ST_-}$. The picture drastically changes towards the in-plane direction where the $S-T_-$ avoided crossing develops interference fringes with a pattern resembling a butterfly; 2 components can be attributed to $S-T_-$ oscillations at low detuning and $S-T_0$ oscillations becoming more prominent at high detuning. The angular anisotropy of the funnel pattern, further exemplified in the Supplemental Material [17], Fig. S5, is the main focus of this work and requires knowledge of the full Hamiltonian and therefore an understanding of the interplay between the $g$-factor anisotropy and the spin-flip element $t_{\text{SF}}$.

In order to extract the $g$-factor anisotropy we rely on singlet-triplet oscillations. After initialization in $S(2,0)$, appropriate pulses to (1,1) induce either $S-T_0$ or $S-T_-$ oscillations. The probability to maintain the initial eigenstate of the system after a sweep with ramp time $\tau_R$ is given by the Landau-Zener formula $P_{\text{LZ}} = \exp[-(2\pi\Delta_{ST}^2/\hbar v)]$ [25,26], where $v = |dE/dt| = |dJ(\epsilon)/d\epsilon|_{\epsilon=e^*} = \epsilon(\Delta_\epsilon/\tau_R)$ is the velocity calculated at $\epsilon = e^*$ and $J(\epsilon) = (\epsilon^2/4) + 2\tau_c - (\epsilon/2)$ is the exchange energy [Fig. 1(b)] [15]. If $v$ satisfies the diabatic condition ($P_{\text{LZ}} \approx 1$) $S-T_0$ oscillations with a frequency $f = (1/\hbar) \sqrt{J^2 + (g^* \mu_B B)^2}$ will be favored. With $P_{\text{LZ}} < 1$ the $S-T_0$ oscillations are suppressed and the qubit is initialized in a superposition of $S$ and $T_-$. After a time $\tau_S$ the system is pulsed back to the measurement point where another nondiabatic passage will cause an interference between the two states [27]. The accumulated phase difference is then given by $\phi = 2\pi f_{S-T_-} \tau_S \approx (\tau_S/\hbar)J - (1/2)g^* \mu_B B$ [28] [Fig. 1(b)]. As the oscillation frequency of the $S-T_0$ ($S-T_-$) qubit depends on $g^*$ ($g^*$) [Fig. 1(b)] we can extract the individual $g$ factors without the need for EDSR. We fix the magnetic field at $|B| = 2\, \text{mT}$ and observe the oscillations vs $\tau_S$ while rotating $B$. We use a fast pulse ($\tau_p = \tau_{\text{rise}} = 2$ ns) in Fig. 2(a) and a ramped pulse with ramp time $\tau_R = 100$ ns in Fig. 2(c). In both cases we pulse to $\epsilon = 4\, \text{meV}$ for a duration $\tau_R$. From the fast Fourier transform (FFT) in Figs. 2(b) and 2(d) we extract the oscillation frequency $f_{S-T_0}$ (orange dots) and $f_{S-T_-}$ (pink dots). We notice that $f_{S-T_-}$ is very close to $f_{S-T_0}$ and $f_{S-T_0}$ is almost $2\, \text{meV}$, which is in line with the observations in Fig. 1(f). Moreover, in Fig. 2(d) the FFT power vanishes for $\theta \approx 60^\circ$, indicating that the ramp time $\tau_R$ induces a completely diabatic passage over the avoided crossing. This is in line with Fig. 1(e) where we observed a sensitive to differences in, or the average of, the Zeeman energies of the dots, and hence we define $g^\pm = g^x \pm g^y$ as the $g$-factor difference and sum. The energy spectrum of the system (the complete Hamiltonian $H_{\text{tot}}$ is derived in the Supplemental Material [17] Sec. VI) is depicted in Fig. 1(b) as a function of $\epsilon$. At $\epsilon = e^*$ the $S$ and $T_-$ states anticross.

We start by mapping out $\Delta_{ST_-}$ as a function of magnetic field angle by varying the magnetic field strength $B$ and $\epsilon$ [24]. We initialize the system deep in (2,0) in a singlet state [point I in Fig. 1(c)], then pulse quickly to (1,1) where the spins are separated (Sep). Mixing between $S$ and $T_-$ is induced when $\epsilon \approx e^*$. In the end we measure the spin state inside the PSB triangle ($M$). The resulting triplet return probability depends both on the size of the avoided crossing and the separation time $\tau_S$. We apply a rapid pulse of duration $\tau_p = 65$ ns and varying $\epsilon$ in order to extract the individual $g$ factors without the need for EDSR. We fix the magnetic field at $|B| = 2\, \text{mT}$ and observe the oscillations vs $\tau_S$ while rotating $B$. We use a fast pulse ($\tau_p = \tau_{\text{rise}} = 2$ ns) in Fig. 2(a) and a ramped pulse with ramp time $\tau_R = 100$ ns in Fig. 2(c). In both cases we pulse to $\epsilon = 4\, \text{meV}$ for a duration $\tau_R$. From the fast Fourier transform (FFT) in Figs. 2(b) and 2(d) we extract the oscillation frequency $f_{S-T_0}$ (orange dots) and $f_{S-T_-}$ (pink dots). We notice that $f_{S-T_-}$ is very close to $f_{S-T_0}$ and $f_{S-T_0}$ is almost $2\, \text{meV}$, which is in line with the observations in Fig. 1(f). Moreover, in Fig. 2(d) the FFT power vanishes for $\theta \approx 60^\circ$, indicating that the ramp time $\tau_R$ induces a completely diabatic passage over the avoided crossing. This is in line with Fig. 1(e) where we observed a
FIG. 2. (a) and (c) Oscillation amplitude of the singlet state in the measurement point as a function of separation time and magnetic field angle at $B = 2$ mT for $\tau_{\theta} = 2$ and $\tau_{R} = 100$ ns, respectively. (b) and (d) FFT of (a) and (c) revealing the oscillation frequency anisotropy. The orange and pink dotted lines are fit to our model. We find a small offset of $\pm 0.04$ ns, $\tau_{\theta}$, for HH states due to a special Ising-type SOI [33], while this type of SOI does not induce out-of-plane terms $\delta b_{z}$ for HH states due to a special Ising-type SOI [32], the inhomogeneous dephasing times extracted for $B_{z}$ of $\approx 700$ ns at 1 mT in Ref. [15] give an upper limit for the hyperfine component $\delta b_{z} < 2$ neV, suggesting that the effects of the nuclear spin bath may safely be neglected.

The lines arising in the FFT plots can be fit by the energy splitting between the three lowest lying states of the system depicted in Fig. 2(e) with $g_{+}^{\perp} = 12.00$, $g_{+}^{R} = 2.04$, $g_{+}^{\parallel} = 0.10$, $g_{+}^{\parallel} = 0.43$, and $t_{C} = 11.38 \mu$eV. The latter is extracted from exchange oscillation measurements (see Supplemental Material [17] Fig. S2).

Interestingly $|g_{+}^{\perp}| > |g_{+}^{R}|$ while $|g_{+}^{\perp}| < |g_{+}^{\parallel}|$. This means that the $g$ factors in the out-of-plane direction have the same sign while they exhibit opposite signs in the in-plane direction. To understand this observation we investigate the effect of the dot geometry on the $g$ factors. As is shown in Supplemental Material [17] Sec. VIII by using the semimicroscopic Luttinger-Kohn Hamiltonian as a starting point, the effects of the intrinsic HH-LH mixing and an elliptical confinement potential can combine to yield $g$-factor renormalizations. While the correction to the out-of-plane $g$ factor is $|\delta g_{\perp}| < 10^{-2}$ for the values considered and hence negligible, the in-plane $g$ factor can be altered considerably,

$$
g_{\perp} = g_{\perp}^{0} - \xi_{1} - \frac{\hbar(\omega_{x} - \omega_{y})}{\hbar(\omega_{x} + \omega_{y})} - \xi_{2}\Delta.
$$

Here, $\xi_{1} \approx 20.3$ and $\xi_{2} \approx 6.0$ are material specific constants, $\Delta$ is the HH-LH splitting, $\hbar\omega_{x,y}$ are the in-plane confinement energies, and $g_{\perp}^{0} = 0.2$ for Ge [30]. It can be seen from Fig. 2(f) that the in-plane $g$-factor corrections can be negative in one dot but not in the other for opposite elliptical confinement.

Electrostatic simulations of the DQD potentials arising from the applied gate voltages [29], not accounting for random disorder potentials, confirm the differently shaped dots. In fact, both dots appear elongated with the major axis of the dots being almost perpendicular to each other [Fig. 2(g) shows the calculated hole density $n_{h}$ and Supplemental Material [17] Sec. IX gives details about the simulation].

We now turn to extract $t_{SO}$ by analyzing $\Delta_{ST_{-}}$ in more detail. We perform Landau-Zener sweeps at $|B| = 20$ mT and extract $\Delta_{ST_{-}}$ from $P_{LZ}$ [Fig. 2(b)] and repeat this for different $\theta$. We vary $\tau_{R}$ during the first passage over the avoided crossing, creating a superposition of $S$ and $T_{-}$, and keep the return sweep diabatic in order to maintain this superposition [Fig. 3(a) and inset of Fig. 3(b)]. The extracted $\Delta_{ST_{-}}$ is reported for different $\theta$ in Fig. 3(c). In general, $\Delta_{ST_{-}}$ may depend on effects influencing the hole spins such as the $g$-factor differences in the two dots, the SOI and possible effective magnetic field gradients caused by the hyperfine interaction [31]. While the hyperfine interaction can result in a strong out-of-plane hyperfine component $\delta b_{z}$ for HH states due to a special Ising-type form [32], the inhomogeneous dephasing times extracted for $B_{z}$ of $\approx 700$ ns at 1 mT in Ref. [15] give an upper limit for the hyperfine component $\delta b_{z} < 2$ neV, suggesting that the effects of the nuclear spin bath may safely be neglected.

In planar HH DQD systems the SOI can be parametrized by a real in-plane spin-orbit vector $t_{SO} = (t_{x}, t_{y}, 0)$. Such in-plane spin-flip tunneling terms stem from the cubic Rashba SOI [33], while this type of SOI does not induce out-of-plane terms $t_{z}$. In a basis in which the total Hamiltonian is diagonal in the absence of the SOI and $g$-factor differences, the $S - T_{-}$ splitting has the form [34]
The single LZ passage pulse sequence (inset) leads to a singlet return probability \( P_S \) that decays exponentially with the ramp time \( \tau_R \). A fit to the Landau-Zener transition formula (black dashed line) allows to extract \( \Delta_{ST} \). (c) \( \Delta_{ST} \) as a function of magnetic field angle. The extracted \( \Delta_{ST} \) is fit to Eq. (2) with \( t_x \) and \( t_y \) as fitting parameters (solid blue line). The black dashed line represents the maximum \( \Delta_{ST} \) as a function of \( \theta \) that can be reliably measured by a single LZ passage. The light colored data points are, therefore, excluded from the fit. (inset) Comparison between the analytical result \( \Delta_{ST} \) and numerical simulation (squares) for \( \Delta_{ST} \) around \( \theta = 0 \). Here, the analytical expression fails due to the small in-plane Zeeman energies.

(d) Comparison between the two contributions to \( \Delta_{ST} \).

\[
\Delta_{ST} = \left| \Delta_{SO} \sin \left( \frac{\Omega}{2} \right) + \Delta_{EZ} \cos \left( \frac{\Omega}{2} \right) \right|, \tag{2}
\]

where the spin-orbit splitting \( \Delta_{SO} \) and the Zeeman splitting \( \Delta_{EZ} \) due to anisotropic site-dependent g tensors read

\[
\Delta_{SO} = t_x + it_x \frac{g^+_{\perp} \sin \theta}{\sqrt{\left( g^+_{\parallel} \cos \theta \right)^2 + \left( g^+_{\perp} \sin \theta \right)^2}}, \quad \Delta_{EZ} = \frac{\mu_B B}{4\sqrt{2}} \frac{\left( g^+_{\parallel} g^+_{\perp} - g^+_{\parallel} g^+_{\perp} \right) \sin(2\theta)}{\sqrt{\left( g^+_{\parallel} \cos \theta \right)^2 + \left( g^+_{\perp} \sin \theta \right)^2}}, \tag{3}
\]

and \( \Omega = \arctan(2\sqrt{t_z/e}) \) is the mixing angle at the anticrossing. The analytical result (2) agrees well with the numerical results obtained by exact diagonalization of the system Hamiltonian for all \( \theta \) except in a narrow region around \( \theta = 0 \) [inset of Fig. 3(c)]. We attribute these deviations to the small in-plane Zeeman energies which violate the assumption of an isolated two-level system made when deriving (2) (see Supplemental Material [17] Sec. VI). Because of the opposite sign \( g \)-factor corrections in the dots the Zeeman splitting \( \Delta_{EZ} \) can be the dominant contribution to \( \Delta_{ST} \), exceeding the spin-orbit splitting by one order of magnitude at small angles. Even when the magnetic field has a large out-of-plane component, the effect of different \( g \) factors can contribute crucially to \( \Delta_{ST} \) [Fig. 3(d)].

The extracted \( \Delta_{ST} \) in Fig. 3(c) can be fit by the model with \( t_x \) and \( t_y \) as free parameters and \( t_C, g^+_{\parallel}, g^+_{ \perp}, g^-_{\parallel}, g^-_{ \perp} \) extracted from previous measurements. Between \(-25^\circ\) and \(25^\circ\) the splitting seems to drop to zero as the Landau-Zener assumptions of diabatic return sweeps are not met and an extraction of \( \Delta_{ST} \) is not accurate. The black dashed line corresponds to the maximum \( \Delta_{ST} \) that allows a diabatic passage with a rise time of \( 2 \) ns of our pulses \( (P_{LZ,max} = 0.99 = \exp[-(2\pi \Delta_{ST,max}^2/h\tau)]) \). The model fits the dark blue data points with \( t_x = 129.0 \pm 18.0 \) and \( t_y = -369.8 \pm 13.8 \) neV, yielding the total spin-flip tunneling element \( t_{SO} = \sqrt{t_x^2 + t_y^2} = 392.0 \) neV.

Having characterized all the elements in the Hamiltonian from independent measurements we can now reproduce the funnel measurements in Fig. 1 (Fig. 4). In particular the sharper line at \( \theta = 60^\circ \) [Fig. 4(b)] as well as the \( S - T_- \) oscillations for \( \theta = 10^\circ \) [Fig. 4(c)] reflect what we observe in the data. Even with \( t_{SO} = 0 \) the in-plane \( g \)-factor difference induces \( S - T_- \) oscillations [Fig. 4(d)] further confirming its dominant role in determining the size of \( \Delta_{ST} \).

In conclusion, we have demonstrated that the \( g \)-tensor anisotropy and, in particular, the in-plane \( g \)-factor...
difference can lead to a considerable contribution to $\Delta_{SO}$ in the in-plane direction. However, Landau-Zener sweeps and singlet-triplet oscillations measured in different magnetic field directions allowed us to distinguish the Zeeman induced coupling from the spin-orbit induced coupling and, thereby, infer the magnitude and orientation of $t_{SO}$. We reconstructed the experimental data in our simulations confirming the validity of our theoretical model. This understanding of the interplay between $t_{SO}$ and the in plane $g$-factor difference opens the possibility to operate hole singlet-triplet qubits at sweet spots, for example with orthogonal axis [34]. Our work, therefore, provides important insight into the spin-orbit interaction of hole spin double quantum dot devices and lays the foundation for the design of future hole spin qubit experiments.

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[1] V. N. Golovach, M. Borhani, and D. Loss, Electric-dipole-induced spin resonance in quantum dots, Phys. Rev. B 74, 165319 (2006).

[2] D. V. Bulaev and D. Loss, Electric Dipole Spin Resonance for Heavy Holes in Quantum Dots, Phys. Rev. Lett. 98, 097202 (2007).

[3] X. Wu et al., Two-axis control of a singlet-triplet qubit with an integrated micromagnet, Proc. Natl. Acad. Sci. U.S.A. 111, 11938 (2014).

[4] J. Yoneda et al., A quantum-dot spin qubit with coherence limited by charge noise and fidelity higher than 99.9%, Nat. Nanotechnol. 13, 102 (2018).

[5] R. Maurand et al., A CMOS silicon spin qubit, Nat. Commun. 7, 13575 (2016).

[6] A. Crippa et al., Electrical Spin Driving by g-Matrix Modulation in Spin-Orbit Qubits, Phys. Rev. Lett. 120, 137702 (2018).

[7] H. Watzinger, J. Kukučka, L. Vukušić, F. Gao, T. Wang, F. Schäffler, J.-J. Zhang, and G. Katsaros, A germanium hole spin qubit, Nat. Commun. 9, 3902 (2018).

[8] N. W. Hendrickx, W. I. L. Lawrie, L. Petit, A. Sammak, G. Scappucci, and M. Veldhorst, A single-hole spin qubit, Nat. Commun. 11, 3478 (2020).

[9] N. W. Hendrickx, D. P. Franke, A. Sammak, G. Scappucci, and M. Veldhorst, Fast two-qubit logic with holes in germanium, Nature (London) 577, 487 (2020).

[10] F. N. M. Froning, L. C. Czamenski, O. A. H. van der Molen, A. Li, E. P. A. M. Bakkers, D. M. Zumbühl, and F. R. Braakman, Ultrafast hole spin qubit with gate-tunable spin-orbit switch functionality, Nat. Nanotechnol. 16, 308 (2021).

[11] K. Wang et al., Ultrafast operations of a hole spin qubit in germanium, Nat. Commun. 13, 206 (2022).

[12] R. Moriya et al., Cubic Rashba Spin-Orbit Interaction of a Two-Dimensional Hole Gas in a Strained-Ge/SiGe Quantum Well, Phys. Rev. Lett. 113, 086601 (2014).

[13] S. Bosco, M. Benito, C. Adelsberger, and D. Loss, Squeezed hole spin qubits in germanium quantum dots with ultrafast gates at low power, Phys. Rev. B 104, 115425 (2021).

[14] Z. Wang, E. Marcellina, A. R. Hamilton, J. H. Cullen, S. Rogge, J. Salti, and D. Culcer, Optimal operation points for ultrafast, highly coherent hole spin-orbit qubits, npj Quantum Inf. 7, 54 (2021).

[15] D. Jirovec et al., A singlet-triplet hole spin qubit in planar Ge, Nat. Mater. 20, 1106 (2021).

[16] R. J. Schoelkopf, P. Wahlgren, A. A. Kozhevnikov, P. Delsing, and D. E. Prober, The radio-frequency single-electron transistor (RF-SET): A fast and ultrasensitive electrometer, Science 280, 1238 (1998).

[17] See Supplemental Material at http://link.aps.org/supplemental/10.1103/PhysRevLett.128.126803 for pulse calibrations, additional data and details of the theoretical model derivation, which includes Ref. [18–23].

[18] X. Marie, T. Amand, P. Le Jeune, M. Paillard, P. Renucci, L. E. Golub, V. D. Dymnikov, and E. L. Ivchenko, Hole spin quantum beats in quantum-well structures, Phys. Rev. B 60, 5811 (1999).

[19] A. V. Trifonov et al., Strong enhancement of heavy-hole landé factor $q$ in ingaas symmetric quantum dots revealed by coherent optical spectroscopy, arXiv:2103.13653.

[20] G. Granger, G. C. Aers, S. A. Studenikin, A. Kam, P. Zawadzki, Z. R. Wasilewski, and A. S. Sachrajda, Visibility study of $S - T_1$ landau-zeren-stückelberg oscillations without applied initialization, Phys. Rev. B 91, 115309 (2015).

[21] L. A. Orona, J. M. Nichol, S. P. Harvey, C. G. L. Bøttcher, S. Scappucci, and M. Veldhorst, A single-hole spin qubit, Nat. Nanotechnol. 5, 487 (2020).

[22] J. Danon and Y. V. Nazarov, Pauli spin blockade in the presence of strong spin-orbit coupling, Phys. Rev. B 80, 041301(R) (2009).

[23] V. N. Golovach, A. Khaetskii, and D. Loss, Phonon-Induced Decay of the Electron Spin in Quantum Dots, Phys. Rev. Lett. 93, 016601 (2004).

[24] J. R. Petta, Coherent manipulation of coupled electron spins in semiconductor quantum dots, Science 309, 2180 (2005).

[25] L. D. Landau, Zur theorie der energieübertragung. II., Phys. Soc. Union 2, 46 (1932).

[26] C. Zener, Non-adiabatic crossing of energy levels, Proc. R. Soc. A 137, 696 (1932).
[27] S. Shevchenko, S. Ashhab, and F. Nori, Landau–zener–
stückelberg interferometry, Phys. Rep. 492, 1 (2010).
[28] J. R. Petta, H. Lu, and A. C. Gossard, A coherent beam
splitter for electronic spin states, Science 327, 669 (2010).
[29] D. L. Craig et al., Bridging the reality gap in quantum devices
with physics-aware machine learning. arXiv:2111.11285.
[30] R. Winkler, Spin—Orbit Coupling Effects in Two-Dimen-
sional Electron and Hole Systems (Springer, Berlin, Heidel-
berg, 2003).
[31] D. Stepanenko, M. Rudner, B. I. Halperin, and D. Loss,
Singlet-triplet splitting in double quantum dots due to
spin-orbit and hyperfine interactions, Phys. Rev. B 85,
075416 (2012).
[32] J. Fischer, W. A. Coish, D. V. Bulaev, and D. Loss, Spin
decoherence of a heavy hole coupled to nuclear spins in a
quantum dot, Phys. Rev. B 78, 155329 (2008).
[33] P. M. Mutter and G. Burkard, Natural heavy-hole flopping
mode qubit in germanium, Phys. Rev. Research 3, 013194
(2021).
[34] P. M. Mutter and G. Burkard, All-electrical control of hole
singlet-triplet spin qubits at low-leakage points, Phys. Rev.
B 104, 195421 (2021).