

Suppression of differential light shifts in ground and metastable trapped-ion qubits

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In the presence of a magnetic field, hyperfine clock qubits can acquire a vector differential light shift that can be tuned via polarization to suppress the total differential light shift of high-power, off-resonant laser light. We experimentally measure this “magic” polarization condition, suppressing differential light shifts in both the $^2S_{1/2}$ ground and $^2F_{7/2}^o$ metastable clock qubits of $^{171}\text{Yb}^+$. We present calculations of the minimum bias magnetic fields required to suppress differential light shifts in the ground-state clock qubits of commonly trapped-ion species, finding that they are below the strengths of fields already typically present in experiments. We further present methods for metastable clock-qubit control in $^{171}\text{Yb}^+$, demonstrating a state preparation and measurement infidelity of $2.9^{+3.0}_{-1.5} \times 10^{-4}$ (-35 ± 4 dB).

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I. INTRODUCTION

In trapped-ion qubit systems, high-power off-resonant laser light is an often-used method for performing single- and two-qubit gates [1,2]. However, this off-resonant laser light can induce an undesirable differential light shift between the qubit states. These light shifts can map gate-laser intensity noise to the qubit frequency, resulting in dephasing and reducing the effective coherence time. Further, in operational schemes that rely on the isolation of qubit subspaces such as the *omg* architecture [3], these differential light shifts may cause detrimental crosstalk errors [4], in which qubits not participating in the interaction experience phase shifts from the perturbing light.

Magnetic-field-induced vector differential light shifts [5] have been shown to give rise to a “magic” polarization condition in neutral atoms that suppresses this source of decoherence [6,7]. Recently, this magic condition was demonstrated in the ground-state qubits of trapped ionic species $^{133}\text{Ba}^+$ [4] and $^{171}\text{Yb}^+$ [8]. In this paper, we explore the theory of magic polarization for ground-state qubits of $^{171}\text{Yb}^+$, $^{173}\text{Yb}^+$, $^{133}\text{Ba}^+$, $^{135}\text{Ba}^+$, $^{137}\text{Ba}^+$, $^{87}\text{Sr}^+$, $^{25}\text{Mg}^+$, $^{43}\text{Ca}^+$, and $^9\text{Be}^+$. We present experimental observations of magic polarization in the ground $m_F = 0$ clock qubit of $^{171}\text{Yb}^+$ and measure the minimum magnetic field B_{crit} required to fully suppress the differential light shift at 532 nm. We discuss operation and state preparation and measurement (SPAM) of the $^2F_{7/2}^o$ metastable $m_F = 0$ clock qubit of $^{171}\text{Yb}^+$, and we again demonstrate magic polarization and measure B_{crit} . Both ground and metastable qubit observations of magic

polarization are operated in magnetic field environments typical of trapped-ion experiments.

II. THEORY

We consider the clock qubit encoded in the $nS_{1/2}$ ground-state hyperfine manifolds characteristic of many commonly used trapped-ion species with nonzero nuclear spin I . This manifold consists of two hyperfine components. We focus on the two B -field-insensitive qubit clock states $|F\rangle \equiv |F, m_F = 0\rangle$ and $|F'\rangle \equiv |F', m'_F = 0\rangle$, where F' is the upper energy state and F is the lower energy state. These are eigenstates for the unperturbed Hamiltonian $\hat{H}_0$. For the positive hyperfine structure constant $A_{nS_{1/2}}$, $F' = I + 1/2$ and $F = I - 1/2$, and the assignment is reversed otherwise. In the absence of applied fields, the qubit angular frequency $\omega_q = E_{F'} - E_F > 0$ ($\hbar \equiv 1$).

Below we review the derivation of the magic conditions, where both qubit levels experience identical perturbations by the laser field. This section is based on the earlier work on magic trapping of neutral atoms [7,9–11], which was recently extended to the laser-noise insensitive operation of the $^{133}\text{Ba}^+$ qubit [4]. We also justify certain approximations made in the earlier work.

Interaction of an ion with a weak laser field of frequency ω_L , amplitude $\mathcal{E}_L$, and polarization $\hat{\mathbf{e}}$ can be parameterized by an optical potential operator,

$$\hat{U} = -\hat{\alpha}(\omega_L, \hat{\mathbf{e}}) \frac{\mathcal{E}_L^2}{4} + \mathcal{O}(\mathcal{E}_L^4), \quad (1)$$

where $\hat{\alpha}(\omega_L, \hat{\mathbf{e}})$ is the operator of dynamic polarizability that can be expanded into irreducible tensor operators of rank $K = 0, 1$, and 2 ,

$$\hat{\alpha}(\omega_L, \hat{\mathbf{e}}) = \hat{\alpha}^{(0)}(\omega_L) + \mathcal{A}\hat{\alpha}^{(1)}(\omega_L) + \hat{\alpha}^{(2)}(\omega_L), \quad (2)$$

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where $\mathcal{A} \equiv -i(\hat{\mathbf{e}}^* \times \hat{\mathbf{e}}) \cdot \hat{\mathbf{z}}$ is the photon's helicity projection on the quantization axis $\hat{\mathbf{z}}$, and ac polarizabilities $\hat{\alpha}^{(K)}$ are proportional to the conventional scalar, vector, and tensor polarizabilities [9].

Importantly, since the laser field breaks the rotational invariance of the atomic Hamiltonian $\hat{H}_0$, $\hat{\alpha}(\omega_L, \hat{\mathbf{e}})$ can have both diagonal and off-diagonal matrix elements in the $\{|F\rangle, |F'\rangle\}$ basis. In an externally applied magnetic field, energies of qubit states can be obtained by diagonalizing the $\hat{H}_0 + \hat{U} + \hat{H}^Z$ matrix, where $\hat{H}^Z$ is the Zeeman Hamiltonian. We will take the quantization axis along the laser propagation direction $\hat{\mathbf{k}}$. In our treatment (and in the experiment), the B field is also directed along $\hat{\mathbf{k}}$. Then $\hat{H}^Z = -\hat{\mu}_z B$, where $\hat{\mu}_z$ is the magnetic dipole operator's z component. While the diagonal matrix elements of $\hat{H}^Z$ vanish for our qubit states (there is no linear Zeeman shift for the $M_F = 0$ states), it admixes the qubit basis states, leading to quadratic Zeeman shifts.

At zero magnetic field, the off-diagonal matrix elements of the optical potential lead to $\mathcal{O}(\mathcal{E}_L^4)$ corrections, which will be ignored for consistency with the omitted hyperpolarizability contributions in Eq. (1). If the laser field is off, eigenenergies repel each other, exhibiting quadratic Zeeman shift. When both fields are on, the correction to the energy splitting can be estimated in the second-order perturbation theory as

$$\Delta\omega_{\text{DLS}} \approx U_{FF} - U_{F'F'} + \frac{4}{\omega_q} \text{Re}(H_{FF'}^Z U_{F'F}), \quad (3)$$

where we omit the quadratic Zeeman shift. This is the differential light shift (DLS) of interest. The DLS expression (3) contains a B -field-free DLS and a Zeeman-optical interference term.

We would like to find the conditions when the DLS vanishes, thereby eliminating the associated laser-induced decoherence. The DLS can be recast in terms of differential polarizability $\Delta\alpha(\omega_L, \hat{\mathbf{e}})$, so that $\Delta\omega_{\text{DLS}} = -\Delta\alpha(\omega_L, \hat{\mathbf{e}}) \mathcal{E}_L^2/4$. Here

$$\Delta\alpha(\omega_L, \hat{\mathbf{e}}) \equiv \alpha_{F'F'} - \alpha_{FF} - \frac{4B\mu_B}{\omega_q} \alpha_{F'F}, \quad (4)$$

where we use $\langle S_{1/2}, F, M_F = 0 | \hat{\mu}_z | S_{1/2}, F', M_{F'} = 0 \rangle = \mu_B$, with μ_B being the Bohr magneton. We suppressed the ω_L and $\hat{\mathbf{e}}$ dependencies of polarizabilities on the right-hand side of this equation for brevity.

Since both qubit states are attached to the same electronic level, their polarizabilities are nearly identical, so that $\alpha_{F'F'} \approx \alpha_{FF}$. The small but important differences arise due to hyperfine interaction (HFI) admixtures to electronic wave functions and energies. The detailed formalism for computing HFI-mediated corrections $\hat{\beta}^{(K)}$ to the ac polarizabilities can be found in Ref. [9], which generalizes the earlier static polarizability derivation [12]. Such HFI-mediated corrections are derived in third order of Floquet perturbation theory and involve matrix elements of HFI and two dipole operators along with energy denominators. Three types of diagrams appear in the calculations depending on the position of the HFI operator in the chain of three operators: top, center, bottom, see Fig. 1 of Ref. [9]. The third-order diagrams also include normalization corrections. We do not reproduce the lengthy formulas from Ref. [9] here, but these are the expressions that

we use in our calculations. Note that simply including hyperfine splittings in the energy denominators while neglecting hyperfine corrections to the electronic wave functions can lead to qualitatively erroneous results [9], because the two effects enter at the same order.

In general, both the scalar ($K = 0$) and tensor ($K = 2$) HFI-mediated polarizabilities $\hat{\beta}^{(K)}$ contribute to the $\alpha_{F'F'} - \alpha_{FF}$ term in Eq. (4). For ground-state alkali-metal atoms, the scalar contribution $\hat{\beta}^{(0)}$ strongly dominates [10]. The same inequality $\beta^{(0)} \gg \beta^{(2)}$ also holds for the alkali-metal-like $S_{1/2}$ ions considered here. A detailed analysis reveals that $\beta^{(2)}$ is dominated by the center diagram which requires HFI matrix elements evaluated between the P_J intermediate states, while $\beta^{(0)}$ is largely determined by the normalization (residual) diagram involving a much larger expectation value of HFI in the $nS_{1/2}$ ground state. The normalization diagram vanishes identically for the $\beta^{(2)}$ tensor polarizability. Thus $\beta^{(0)} \gg \beta^{(2)}$ and we neglect tensor contribution, so that $\alpha_{F'F'} - \alpha_{FF} \approx \beta_{F'F'}^{(0)} - \beta_{FF}^{(0)}$ in Eq. (4). To quantify systematic errors associated with this approximation, we carried out numerical tests for the Yb^+ ion which has a softer core compared to other ions of interest and we found that neglecting tensor polarizability introduces the largest few-percent fractional errors near the $6S_{1/2}$ - $6P_J$ resonances. These two scalar polarizabilities are strictly proportional to each other precluding B -field-free magic wavelengths for the $S_{1/2}$ hyperfine manifolds [10]. This is where the Zeeman-optical interference term in Eq. (3) becomes important as it can cancel out the residual B -field-free DLS.

The Zeeman-optical interference term is dominated by the conventional vector (axial) polarizability $\alpha_{nS_{1/2}}^a(\omega_L)$. The off-diagonal matrix element of the scalar polarizability vanishes by angular selection rules, while the tensor contribution for an $S_{1/2}$ state requires HFI admixtures and is suppressed. Neglecting this tensor contribution, the off-diagonal polarizability is [11]

$$\alpha_{F'F}(\omega_L, \hat{\mathbf{e}}) \simeq \mathcal{A} \alpha_{FF}^{(1)}(\omega_L) = \frac{\mathcal{A}}{2} \alpha_{nS_{1/2}}^a(\omega_L).$$

Consequently, the Zeeman-optical contribution to the differential polarizability in Eq. (4) is explicitly

$$-\frac{4B\mu_B}{\omega_q} \alpha_{F'F} \simeq -\mathcal{A} \frac{2B\mu_B}{\omega_q} \alpha_{nS_{1/2}}^a(\omega_L).$$

As discussed earlier, $\alpha_{F'F'} - \alpha_{FF} \approx \beta_{F'F'}^{(0)} - \beta_{FF}^{(0)}$, which simplifies further using [11] $\beta_{F'F'}^{(0)} = -(I+1)\beta_{FF}^{(0)}/I$. Then setting the right-hand side of Eq. (4) to 0, we find that DLS vanishes at the “magic” helicity projection

$$\mathcal{A}_m(\omega_L) \approx -\frac{2I+1}{2I} \frac{\beta_{FF}^{(0)}(\omega_L)}{\alpha_{nS_{1/2}}^a(\omega_L)} \frac{\hbar\omega_q}{\mu_B B} \equiv -\frac{B_{\text{crit}}(\omega_L)}{B}, \quad (5)$$

where $\beta_{FF}^{(0)}(\omega_L)$ is the HFI-mediated scalar polarizability of the lower qubit clock state. Here we define the critical magnetic field B_{crit} as the minimum field for which the DLS can be canceled by an allowed, $|\mathcal{A}| \leq 1$, polarization. Fields larger than B_{crit} allow cancellation with partially circular polarization, while fields below B_{crit} cannot provide a sufficiently large Zeeman-optical interference contribution to compensate the scalar differential shift in Eq. (4).

The ratio $\beta_{FF}^{(0)}(\omega_L)/\alpha_{nS_{1/2}}^a(\omega_L)$ is on the order of the ratio of the hyperfine splitting ω_q to the fine-structure splitting in the lowest energy P -state manifold. In addition, the vector polarizability $\propto \omega_L$, and we expect $B_{\text{crit}}(\omega_L)$ to scale as $1/\omega_L$ below the lowest-energy $nS_{1/2}$ - $nP_{1/2,3/2}$ resonance.

If the B field and the direction of laser propagation are set at an angle θ , then $\mathcal{A}$ in Eq. (5) is proportional to $\cos \theta$, providing an additional experimental handle. Increasing the angle and reducing the degree of circular polarization raises the critical B field.

The DLS cancellation is independent of the laser intensity: intensity sets the magnitude of the uncompensated differential light shift, while the ratio of scalar to vector polarizability determines the polarization required for cancellation. At the magic polarization, the two clock states experience equal light shifts from the perturbing laser, so fluctuations in the laser intensity no longer map directly onto fluctuations of the qubit state splitting.

Numerical evaluation of critical fields

We carried out critical-field calculations for the $S_{1/2}$ ground-state hyperfine clock manifolds of several commonly used trapped-ion species: $^{171}\text{Yb}^+$, $^{173}\text{Yb}^+$, $^{133}\text{Ba}^+$, $^{135}\text{Ba}^+$, $^{137}\text{Ba}^+$, $^{87}\text{Sr}^+$, $^{25}\text{Mg}^+$, $^{43}\text{Ca}^+$, and $^9\text{Be}^+$. All these isotopes have a nonzero nuclear spin I tabulated in the Supplemental Materials [13] along with the nuclear g -factors g_I .

The critical B field is expressed as [cf. Eq. (5)]

$$B_{\text{crit}}(\omega_L) = \frac{2I + 1}{2I} \frac{\beta_{FF}^{(0)}(\omega_L)}{\alpha_{nS_{1/2}}^a(\omega_L)} \frac{\hbar\omega_q}{\mu_B}. \quad (6)$$

Computations of vector polarizability $\alpha_{nS_{1/2}}^a(\omega_L)$ of the ground $nS_{1/2}$ state involve a summation over intermediate opposite-parity atomic states $|i\rangle = |n_iP_{1/2,3/2}\rangle$, electric-dipole matrix elements $\langle nS_{1/2}|\mathbf{D}|n_iP_{1/2,3/2}\rangle$, and energy denominators of both rotating and counter-rotating form $E_g - E_i \pm \omega_L$ (E_g is the ground-state energy). The energy denominators lead to a resonant behavior of the polarizabilities. Determination of the HFI-induced scalar polarizability $\beta_{FF}^{(0)}(\omega_L)$ introduces an additional summation over intermediate states and requires magnetic-dipole HFI matrix elements, especially of the $S_{1/2}$ - $S_{1/2}$ character for the top and bottom diagrams and P_J - P_J for the center diagrams. The required dipole matrix elements additionally include P_J - D_J channels.

We employed *ab initio* relativistic random-phase approximation and Brueckner orbital (RPA + BO) methods implemented as described in Ref. [14]. Both RPA and BO are all-order (in residual electron-electron Coulomb interaction) many-body methods. To ensure the completeness of the intermediate state basis and to improve numerical accuracy in computing HFI, we employed the dual-kinetic-balance B-spline basis set [15]. These basis sets were generated in the V^{N-1} frozen core Dirac-Hartree-Fock (DHF) potentials with a finite-size nucleus and further rotated by diagonalizing the DHF + BO Hamiltonian to generate numerically complete independent particle DHF + BO basis sets [14].

A typical accuracy of the relativistic RPA + BO method is sub-percent for light monovalent ions such as Be^+ and a few percent for heavy ions. The most demanding is the

TABLE I. Reduced electric-dipole matrix elements in Yb^+ , in atomic units ea_0 . The RPA+BO values are the present *ab initio* relativistic values. MBPT denotes the third-order relativistic many-body perturbation theory values of Ref. [16], and RCC denotes the relativistic coupled-cluster values of Ref. [17]. Experimental entries are direct or lifetime/branching-fraction inferred values from Refs. [18,19]. Signs of reduced matrix elements are phase-convention dependent; magnitudes are listed.

Transition	RPA + BO	MBPT	RCC	Expt.
$6s_{1/2}$ - $6p_{1/2}$	2.667	2.683	2.72(1)	2.470(3)
$6s_{1/2}$ - $6p_{3/2}$	3.759	3.768	3.83(1)	3.36(3)
$5d_{3/2}$ - $6p_{1/2}$	3.047	2.971	3.06(2)	2.998(2)
$5d_{3/2}$ - $6p_{3/2}$	1.343	1.307	1.35(2)	1.15(4)
$5d_{5/2}$ - $6p_{3/2}$	4.216	4.120	4.23(3)	3.95(10)

case of Yb^+ with an easily polarizable core due to the outer f shell. We present a comparison of the RPA + BO dipole matrix elements for Yb^+ with literature values in Table I. Notice that even more complete (than our RPA + BO) relativistic coupled-cluster treatment substantially disagrees with experimental values for the $6s_{1/2}$ - $6p_j$ matrix elements. To improve the accuracy of resonances we replaced our theoretical DHF + BO energies by the NIST recommended values [20] for several lowest-lying energy levels of the $S_{1/2}$, P_J , and D_J symmetry. To further improve the accuracy, we used literature values for dipole matrix elements and hyperfine structure constants for low-lying states. The employed non-DHF + BO values are listed in the Supplemental Material [13]. Off-diagonal matrix elements of the $n's$ - $n''s$ hyperfine interaction were computed using the geometric mean of the $A_{n's}$ and $A_{n''s}$ hyperfine structure constants to substantially improve the accuracy [21] of these HFI matrix elements. The remaining matrix elements and energies were our RPA + BO values.

For each laser frequency, these scalar and axial polarizabilities were inserted into Eq. (6). The resulting values of B_{crit} as a function of the laser frequency are shown in Fig. 1. For most wavelengths relevant to trapped-ion operation, the calculated critical fields are in the range of a few Gauss, comparable to bias fields already used in typical ion-trap experiments to destabilize coherent dark states and resolve the Zeeman structure. These $B_{\text{crit}}(\omega_L)$ curves reflect the ratio of a small HFI-mediated scalar response to the ordinary electronic vector response. Resonant features in Fig. 1 occur at strong $S_{1/2} \rightarrow P_J$ electric-dipole transitions, where both polarizabilities vary rapidly. $B_{\text{crit}}(\omega_L)$ scales roughly as $1/\omega_L$ below the lowest-energy $nS_{1/2}$ - $nP_{1/2}$ resonance as expected.

The isotope dependence of the critical field follows from the HFI-mediated scalar polarizability and the clock frequency in Eq. (6). We neglect hyperfine anomalies and small isotope shifts in the electronic matrix elements. Then at fixed laser frequency and within a given ion species, the electronic vector polarizability and electronic sums entering $\beta^{(0)}$ are approximately isotope independent.

After angular reduction for the two $S_{1/2}$ hyperfine components [11], $\beta_{I+1/2}^{(0)} = -(I+1)\beta_{I-1/2}^{(0)}/I$ and $\beta_{I-1/2}^{(0)} \propto g_I I$, so

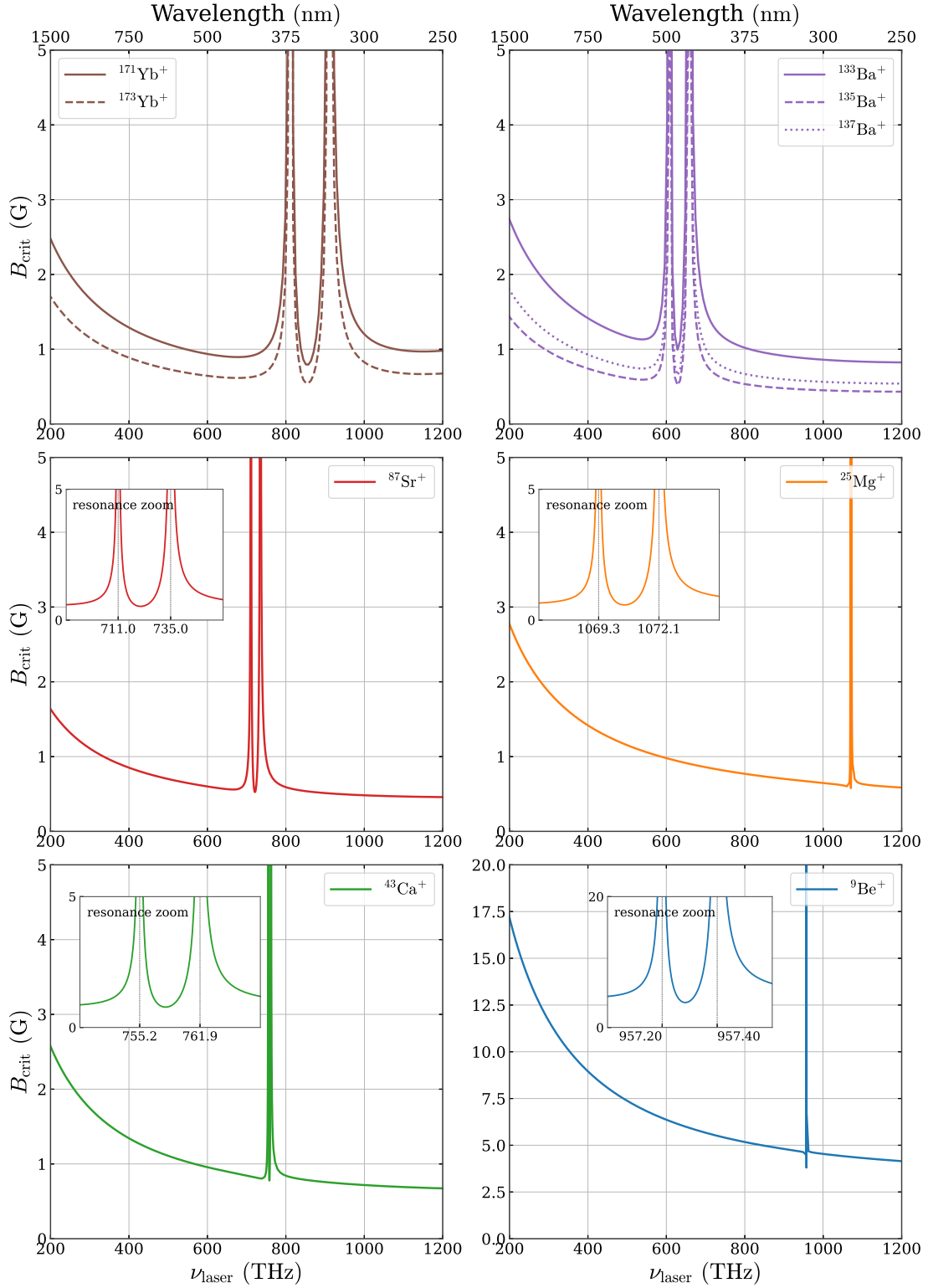


FIG. 1. Full-scale critical magnetic field calculations including the hyperfine interaction for the ground-state qubit of various commonly trapped ions. Note that for most driving frequencies and ions, the critical field is within the range typical to ion trap operation [$O(\text{Gauss})$]. The most-likely exception to this is ${}^9\text{Be}^+$, for which the vertical scale has been expanded to show B_{crit} (bottom right).

the differential scalar response scales as $g_I(2I + 1)$. The clock splitting also scales as $\omega_q \propto |g_I|(2I + 1)$. Therefore,

$$B_{\text{crit}}(\omega_L) \propto g_I^2(2I + 1)^2. \quad (7)$$

This implies that a measurement of the critical field at fixed laser frequency in one isotope can be rescaled to another isotope of the same ion species.

An analogous DLS-cancellation mechanism applies in principle to the ${}^{171}\text{Yb}^+ 4f^{13}6s^2 {}^2F_{7/2}^o$ hyperfine clock qubit.

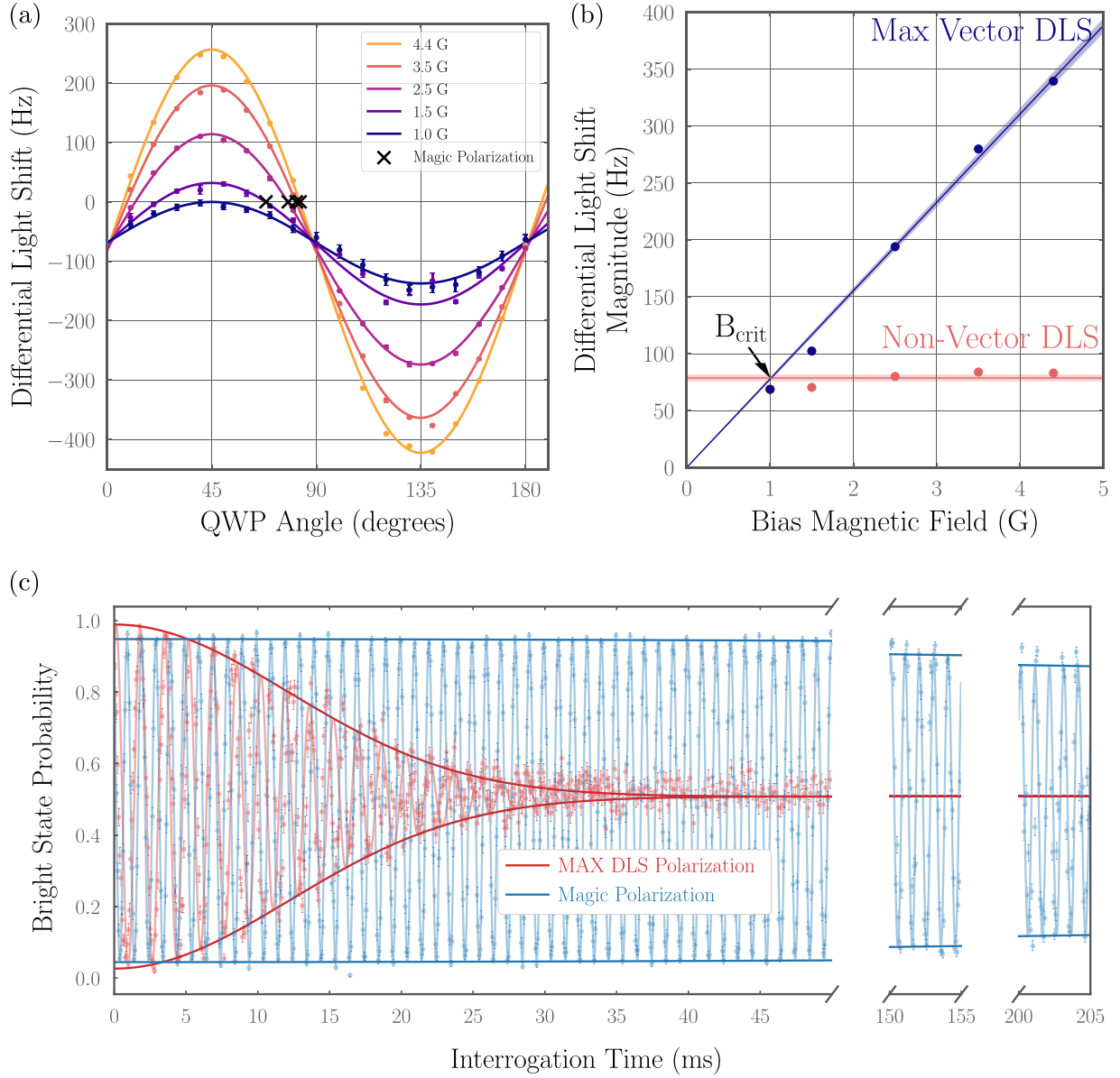


FIG. 2. (a) A polarization-dependent DLS in the $^2S_{1/2}$ $m_F = 0$ clock qubit of $^{171}\text{Yb}^+$ as a function of the QWP angle at different bias magnetic field magnitudes. The $\times$ markers denote magic polarizations for which the DLS is 0. (b) Values of the nonvector and maximum-vector DLS magnitudes with respect to bias magnetic field magnitude. The critical field is calculated from the intersection of these values as $B_{\text{crit},g} = 1.01 \pm 0.04$ G. (c) Detuned microwave Ramsey scans with the perturbing laser on at a maximum DLS inducing polarization vs the magic polarization shows an order of magnitude improvement to coherence time, achieving a value comparable to the control scan.

A quantitative calculation is substantially more demanding than that for the alkali-metal-like $4f^{14}6s$ $^2S_{1/2}$ ground state because it requires an open-shell many-body treatment of the dense excited-state spectrum and a consistent evaluation of HFI corrections to both energies and wave functions, with the scalar, vector, and tensor polarizabilities retained. The single-valence-electron RPA + BO framework used here is not directly applicable to this manifold without substantial extension. As an order-of-magnitude guide, a reduced model that includes hyperfine splittings in the energy denominators but neglects HFI-induced corrections to the wave functions gives [22] $B_{\text{crit},m} \simeq 0.15$ G at 532 nm. Because the omitted contributions enter at the same perturbative order as

those retained, this value is not a precision prediction; the experimental determination reported below, $B_{\text{crit},m} = 1.1 \pm 0.7$ G, provides the relevant practical cancellation condition. In addition, we expect an additional correction due to the nonvanishing tensor contribution to polarizabilities for the $J = 7/2$ state.

III. EXPERIMENTS

A. Magic polarization in the $^{171}\text{Yb}^+$ ground $^2S_{1/2}$ qubit

The ground-state (“g-type”) clock qubit of $^{171}\text{Yb}^+$ is defined in the $^2S_{1/2}$ hyperfine manifold between $|0, 0\rangle$ and $|1, 0\rangle$ (denoted $|F, m_F\rangle$), with frequency splitting

$\omega_g \approx 2\pi \times 12.64$ GHz. Doppler cooling, optical pumping, state readout, and qubit rotations are effected using 369- and 935-nm lasers with fiber electro-optic modulators and an external microwave antenna, operated as in Ref. [23].

Differential light shifts are applied with a continuous-wave laser at $\lambda = 532$ nm with a $1/e^2$ intensity radius of $w_0 \approx 35$ μm at the ion. For the ground-state qubit we apply 1.0 W of this light and infer a standard deviation in power of 27(4) mW at the ion. We measure differential light shifts using a modified detuned microwave Ramsey sequence in which the perturbing laser is on during the Ramsey delay time. We interleave these modified DLS Ramsey scans with control Ramsey scans in which the perturbing laser is off. We fit the frequencies of both Ramsey signals to compute the differential light shift.

The perturbing laser transmits through a polarizing beam splitter for linear input polarization to a zero-order quarter-wave plate (QWP). Rotation of the QWP creates polarization of sinusoidally oscillating helicity. The bias magnetic field points parallel to the perturbing laser $\hat{\mathbf{k}}$ such that π -polarized light is not realizable. We rotate the QWP through 180° and measure the DLS induced by the perturbing laser, repeating these measurements for five different bias magnetic fields, shown in Fig. 2(a). We fit these data [Fig. 2(b)] and find $B_{\text{crit},g} = 1.01 \pm 0.04$ G from the values of the nonvector (scalar and tensor) and maximum-vector DLSs.

To demonstrate preservation of qubit coherence, we compare the coherence times of three different detuned microwave Ramsey scans, shown in Fig. 2(c). One scan was performed with the perturbing laser on at a maximum DLS-inducing polarization, a second with the perturbing laser on at the magic polarization, and a control scan with the perturbing laser off. All three scans were performed with the same bias magnetic field of 5 G. We observe the coherence time with the maximum DLS-inducing polarization to be 17.2 ± 0.2 ms, while the magic polarization and control scan are comparable to one another and an order of magnitude longer. Auxiliary experiments probing the residual light shift between the control and the magic polarization scans found a residual light shift of 0.21(4) Hz, representing a suppression of the DLS by a factor of 2000 ± 400 . From this suppression of the DLS and a decoherence model assuming Gaussian amplitude fluctuations of the perturbing laser, we may infer an extension in the coherence time by the same amount.

B. Metastable qubit operations

The metastable-state (“*m*-type” [3]) clock qubit is defined in the long-lived (3.16 years [24]) $^2F_{7/2}^o$ hyperfine manifold between the $|3, 0\rangle$ and $|4, 0\rangle$ states, with zero-field qubit frequency splitting we measured to be $\omega_m = 2\pi \times 3.62052720(5)$ GHz. Other work in Ref. [25] used 411- and 3432-nm pulses with the intermediate $^2D_{5/2}$ state to achieve coherent microsecond-level one-way conversion rates of 99.5% between ground and metastable qubits, demonstrating a metastable qubit SPAM error of 2.3(2)%. In contrast, we incoherently populated the metastable state from $^2D_{5/2}$ decays, then heralded state preparation with 760-nm light. Our metastable SPAM protocol does not use any sub-Doppler cooling techniques or require motional coupling and is explained in detail below with relevant states shown in Fig. 3.

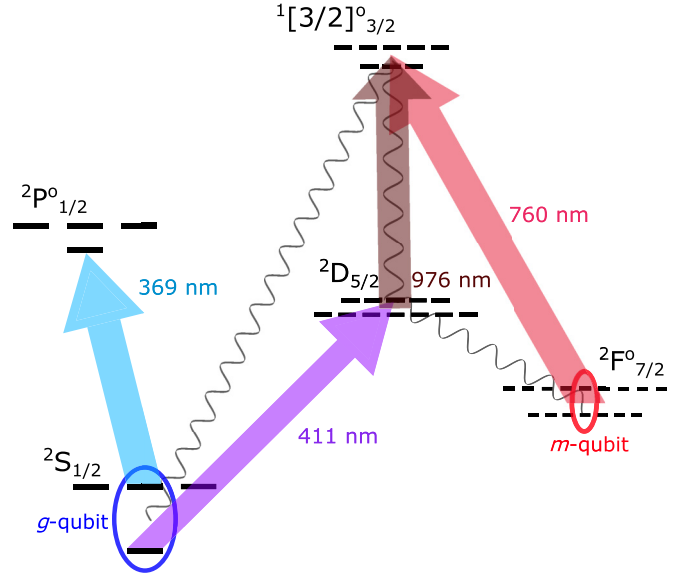


FIG. 3. The 369-nm laser controls Doppler cooling, optical pumping, and fluorescence state detection in the $^2S_{1/2}$ ground qubit with two electro-optic modulator tones. The metastable qubit is prepared by driving an $E2$ ($^2S_{1/2} F = 0$ to $^2D_{5/2} F = 2$) transition at 411 nm, which decays to $^2F_{7/2}^o F = 3$. Readout is performed with a $^2F_{7/2}^o F = 3$ to $^1[3/2]^o_{3/2} F = 1$ 760-nm transition, which decays to the $^2S_{1/2}$ hyperfine manifold for fluorescence detection. The 976-nm laser accompanies the 760 nm laser to repump population from $^2D_{5/2}$ during readout operations. Qubit rotations for both the *g*- and *m*-type qubits are effected via microwaves applied by external horn antennas.

We first Doppler cool and optically pump into $^2S_{1/2} |0, 0\rangle$. A 411-nm laser with polarization optimized for $\Delta m_F = \pm 1$ electric quadrupole ($E2$) transitions is applied for 100 ms to populate $^2D_{5/2} |2, \pm 1\rangle$. This state may decay back to $^2S_{1/2}$, but preferentially decays to $^2F_{7/2}^o F = 3$, populating the $^2F_{7/2}^o |3, 0\rangle$ clock state with $\approx 20\%$ probability.

To prepare the $^2F_{7/2}^o |4, 0\rangle$ state, we drive a microwave π pulse from $^2F_{7/2}^o |3, 0\rangle$ and then drive any remaining population in $^2F_{7/2}^o F = 3$ to $^1[3/2]^o_{3/2} F = 1$ with a 760-nm laser applied for 50 ms. This excited state has a short lifetime (29 ns [26]) and preferentially decays to the $^2S_{1/2}$ ground state [27]. We then use ground-state Doppler cooling light to detect any population in the $^2S_{1/2}$ manifold, which would indicate failed preparation of $^2F_{7/2}^o |4, 0\rangle$. Our pulse sequencer is not capable of changing the pulse sequence depending on photomultiplier tube (PMT) count readings during this heralding step, so we run the full pulse sequence for each experiment and remove the data showing unsuccessful state preparation. In principle, however, this state preparation could be repeated until successful in an algorithmic context.

To prepare the other qubit state $^2F_{7/2}^o |3, 0\rangle$, we perform a second heralding sequence after preparing $^2F_{7/2}^o |4, 0\rangle$. Another *m*-type qubit π -pulse prepares $^2F_{7/2}^o |3, 0\rangle$, and a second 760-nm laser applied for 50 ms drives any population remaining in $^2F_{7/2}^o |4, 0\rangle$ to $^1[3/2]^o_{3/2} F = 2$ where it decays to the ground state. Fluorescence detection in the $^2S_{1/2}$ manifold then indicates failed preparation of $^2F_{7/2}^o |3, 0\rangle$.

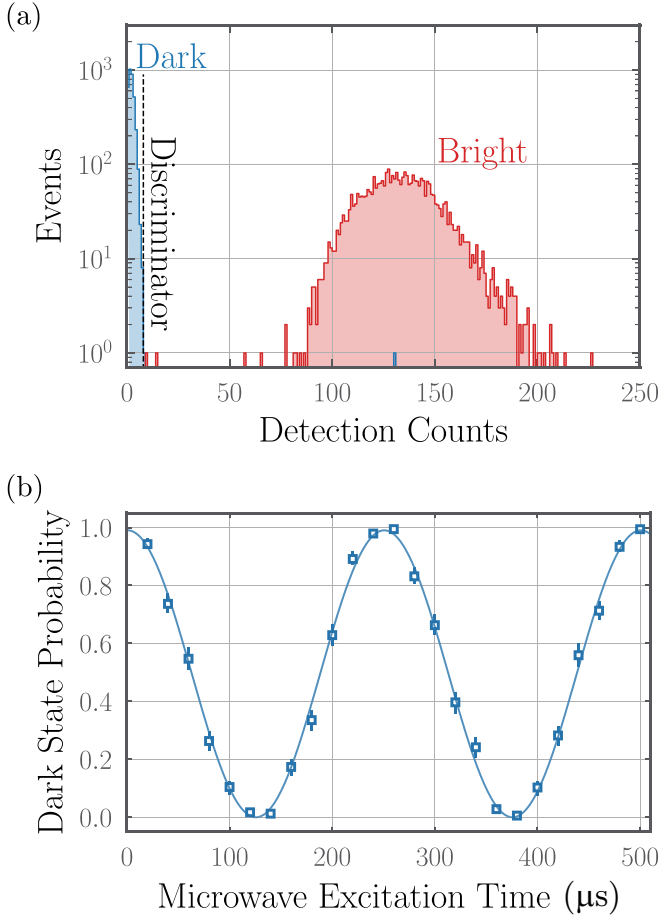


FIG. 4. (a) Detected photon count histograms for prepared metastable bright and dark states during 3.5 ms of state detection time. (b) Metastable-qubit Rabi flops show high contrast even without correcting for finite SPAM fidelity.

Readout at the end of an experiment involves first applying a 760-nm laser pulse for 50 ms to drive population out of $^2F_{7/2}^o |3, 0\rangle$, followed by 369-nm fluorescence detection in the $^2S_{1/2}$ manifold for 3.5 ms. After readout, any population remaining in the $^2F_{7/2}^o$ manifold is returned to the Doppler cooling cycle with both 760-nm lasers applied for 50 ms. To repump the decay from the $^1[3/2]_{3/2}^o$ manifold to $^2D_{5/2}$, 760-nm lasers are accompanied by a 976-nm laser.

State detection histograms for each target state prepared using this procedure are shown in Fig. 4(a), with high-contrast single-qubit rotations on the m -type qubit shown in Fig. 4(b). The average m -type qubit SPAM infidelity using the two-step heralding outlined above is $2.9^{+3.0}_{-1.5} \times 10^{-4} = -35 \pm 4$ dB. This infidelity is limited by off-resonant driving of $^2F_{7/2}^o |4, 0\rangle$ to $^1[3/2]_{3/2}^o F = 2$ during readout, resulting in a dark-state reading as bright.

C. Magic polarization in the $^{171}\text{Yb}^+$ metastable $^2F_{7/2}^o$ qubit

We repeat a similar QWP angle experiment as in the ground state to measure the DLS in the metastable clock qubit (Fig. 5). High-fidelity operations are not required to observe DLS suppression, so we operate the metastable SPAM pro-

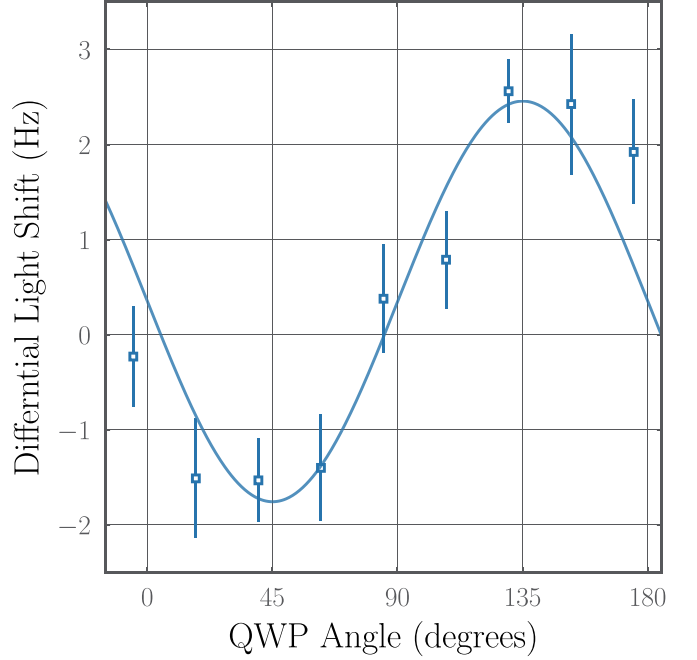


FIG. 5. Polarization-dependent DLS in the $^2F_{7/2}^o m_F = 0$ clock qubit as a QWP is rotated through 180° with a 6.75-G bias magnetic field.

ocol with shorter shelving and deshelling pulse lengths and without the 976-nm laser to relax technical challenges. Due to the small magnitude of the m -type DLS, high measurement sensitivity is required to observe the vector light shift. We therefore increase the perturbing laser power to approximately 1.85 W and increase the magnetic field magnitude to 6.75 G. Performing the experiment, we observe a small polarization-dependent DLS with a zero crossing “magic” polarization, shown in Fig. 5. We estimate $B_{\text{crit},m} = 1.1 \pm 0.7$ G from the values of the nonvector and maximum vector DLSs.

IV. CONCLUSION

We suppress differential light shifts via magic polarization in both the $^2S_{1/2}$ g -type and $^2F_{7/2}^o$ m -type qubits of $^{171}\text{Yb}^+$, and we recover a coherence time in the g -type qubit in the presence of a perturbing laser that is comparable to the coherence time without the perturbing laser. The critical magnetic field is measured as the minimum field for which the DLS may be eliminated through magic conditions with polarization orthogonal to the quantization axis. The observed g -type qubit critical field $B_{\text{crit},g} = 1.01 \pm 0.04$ G is consistent with the theoretical prediction $B_{\text{crit},g} = 0.98$ G. For the m -type qubit, we measure a comparable magnitude of $B_{\text{crit},m} = 1.1 \pm 0.7$ G.

Using Eq. (7) and our measurement in $^{171}\text{Yb}^+$, we estimate the critical magnetic field at 532 nm to be $B_{\text{crit}}^{(173)} = 0.66(3)$ G for the ground-state qubit of $^{173}\text{Yb}^+$. Applying the same scaling to the $^{133}\text{Ba}^+$ critical field of 1.17 G measured in Ref. [4], we infer $B_{\text{crit}}^{(135)} = 0.61$ G and $B_{\text{crit}}^{(137)} = 0.76$ G for the ground states of $^{135}\text{Ba}^+$ and $^{137}\text{Ba}^+$ at 532 nm, respectively.

The critical fields observed and inferred in this work for all isotopes of Ba^+ and Yb^+ are well below the field employed for typical ion trap operation, indicating that this procedure

may be adopted by other researchers with minimal modifications. For all ions considered except ${}^9\text{Be}^+$, the calculated magnetic field required to observe the magic polarization is below 3 G even for $\lambda = 1.5\ \mu\text{m}$ and is below 10 G for ${}^9\text{Be}^+$ for $\lambda < 800\ \text{nm}$.

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DATA AVAILABILITY

There are no publicly available research data or software supporting this manuscript. Requests for further information or data should be sent to the authors.

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