

Observation and measurement of the 1S_0 - 3P_0 transition in an optical lattice clock using the ^{198}Hg bosonic isotope

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We report the observation of the magnetic-field-induced $(5d^{10}6s^2)^1S_0$ - $(5d^{10}6s6p)^3P_0$ transition in a bosonic isotope of mercury, ^{198}Hg , realized in an optical lattice clock. We characterize this ^{198}Hg clock transition at 265.6 nm, determining key features such as the quadratic Zeeman shift, the probe light shift, and the magic frequency. We also report a comparison between the ^{198}Hg optical lattice clock and ^{87}Sr . In this comparison, the ^{198}Hg clock has a relative frequency stability of $6 \times 10^{-16}/\sqrt{\tau}/\text{s}$ and a total relative systematic uncertainty of 6.9×10^{-16} . This comparison yields the first direct determination of the $^{198}\text{Hg}/^{87}\text{Sr}$ optical frequency ratio: $^{198}\text{Hg}/^{87}\text{Sr} = 2.629\,315\,734\,684\,118\,1$, with the same relative uncertainty.

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I. INTRODUCTION

Recent advances in optical frequency metrology, particularly the development of ultrastable lasers with fractional frequency instabilities below 10^{-16} [1–5], have enhanced the performance of optical clocks. Fully leveraging such lasers requires the capability to probe the clock transition for durations exceeding 1 s. Atomic species that exhibit not only minimal sensitivity to external perturbations but also extremely long-lived clock states are the most promising candidates for achieving the highest levels of clock stability and accuracy. This combination of properties is found in several neutral-atom and ion-based clocks presently studied, as well as in proposed future nuclear clocks [6,7]. Optical clocks play a critical role in testing fundamental physical laws and in searches for new physics [8–11]. They are also under active consideration for the redefinition of the second [12].

Among neutral atoms, mercury (Hg) offers promising properties for optical lattice clocks. It is significantly less sensitive to blackbody radiation and stray electric fields, 30 times less so than Sr, for example [13]. Hg is relevant for fundamental physics tests and dark matter searches [14] owing to its sensitivity to putative variations of the fine-structure constant [15] and to its seven natural isotopes enabling searches based on isotope shifts [11,16].

To date, only the 1S_0 - 3P_0 clock transition of fermionic ^{199}Hg has been studied in a lattice clock [13,17–19]. This transition is weakly allowed by hyperfine mixing, which governs the spontaneous emission rate from 3P_0 to 1S_0 , setting the 3P_0 lifetime. In ^{199}Hg , the lifetime is short, ~ 1.5 s [20,21], which

may limit the advantages of next-generation ultrastable lasers [1,2] that enable longer probing times and improved stability [22–24].

Using bosonic mercury isotopes circumvents this limitation. With zero nuclear spin, their 1S_0 - 3P_0 transition is strictly forbidden, so the 3P_0 state does not spontaneously decay. This transition can be weakly induced by applying a strong static magnetic field $\mathbf{B}$, as proposed in Ref. [25] and demonstrated in ^{88}Sr [26–28] and ^{174}Yb [29]. The magnetic field primarily mixes the 3P_0 and 3P_1 states (Fig. 1) via the magnetic-dipole coupling $\hbar\Omega_B = \langle ^3P_0 | \hat{\mu} \cdot \mathbf{B} | ^3P_1 \rangle$, with $\hat{\mu}$ the magnetic-dipole operator, perturbing 3P_0 with a small 3P_1 admixture. As a result, the otherwise forbidden 1S_0 - 3P_0 transition acquires a weak electric-dipole amplitude through the 1S_0 - 3P_1 coupling $\hbar\Omega_L = \langle ^1S_0 | \hat{\mathbf{d}} \cdot \mathbf{E} | ^3P_1 \rangle$, driven by an optical field of amplitude $\mathbf{E}$ and frequency ω . Residual coupling to other states, such as 1P_1 , also exists (Fig. 1). This magnetic-field-induced mixing allows tuning at will of the transition strength, enabling extended probe times to match the coherence time of the probe laser.

However, direct laser excitation of a bosonic mercury clock transition has not yet been observed. Here, we report the observation of the ^{198}Hg bosonic transition in an optical lattice clock and the first measurement of the $^{198}\text{Hg}/^{87}\text{Sr}$ optical frequency ratio, enabled by several key experimental advances and a challenging search for a narrow transition across a wide uncertainty range.

II. SEARCH FOR THE 1S_0 – 3P_0 TRANSITION IN THE BOSONIC ISOTOPE ^{198}Hg

The setup for searching the bosonic clock transition builds on the existing experiment used for probing the ^{199}Hg fermionic isotope, as described in Refs. [19,30–34].

First, a well-controlled and sufficiently strong magnetic field is required to induce the mixing. For this, we generate a pulsed magnetic field by using the 3D MOT coils and switching electronics to rapidly reverse the current in one of

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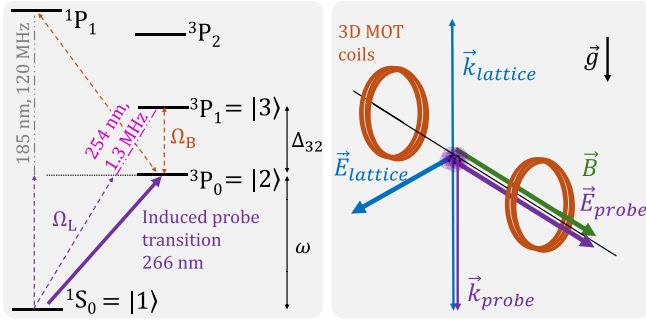


FIG. 1. (Left) The ^{198}Hg bosonic clock transition magnetic-field-induced mixing scheme: The doubly forbidden $1S_0 \rightarrow 3P_0$ transition is enabled by a strong static magnetic field $\mathbf{B}$, coupling the $3P_0$ and $3P_1$ states via the matrix element Ω_B through magnetic dipole interaction (dotted orange lines). The cooling transition $1S_0 \rightarrow 3P_1$ with its electric-dipole coupling Ω_L is indicated by the dotted pink line. (Right) Probing configuration for ^{198}Hg : The probe beam polarization is aligned with the magnetic field from three-dimensional (3D) magneto-optical trap (MOT) coils. The vertical lattice is polarized perpendicular to the probe beam.

the coils during interrogation, following the approach used in Ref. [29]. A maximum field of $B = 18.6$ mT can be reached with 12 A current, with rise/fall times as fast as 3/2.8 ms to minimize unwanted dead times. Three additional pairs of coils compensate residual background magnetic fields. Their current is set by minimizing the Zeeman splitting of the ^{199}Hg clock transition.

Next, we extended our ultrastable light generation to provide a sufficiently broad continuous tuning range to reach the bosonic isotope. We implemented a setup with an intermediate laser [35], offset phase locked to the ultrastable light as detailed in Appendix A. This extends the tuning range at 265.6 nm from a few tens of megahertz [30,31] to over 10 GHz, enabling ultrastable probing of all mercury isotopes, not just ^{199}Hg . We also ensured that spectral purity was maintained at the present 4×10^{-16} flicker frequency noise level for the ultrastable light [36] and that the system remains compatible with future improvements well beyond this stability.

To facilitate the search for the transition, achieving the highest possible coupling in our setup was critical. As it scales with intensity, we first enhanced our deep-UV probing light to a maximum of 400 μW at 265.6 nm. The probe beam was reshaped and its waist was estimated to be 64 ± 5 μm (see Appendix A for the beam characterization), yielding peak intensities as high as $I = 6.3$ W/cm 2 .

Following notations of Ref. [25], and with the probe polarization aligned along the applied magnetic field B , the estimated Rabi frequency is $\Omega_{12} = \frac{\Omega_L \Omega_B}{\Delta_{32}} = \alpha \sqrt{IB}$. Here, α is the magnetic-field-induced coupling coefficient and Δ_{32} the angular frequency detuning between $3P_1$ and $3P_0$ (see Fig. 1). We predict $\Omega_{12}/2\pi \simeq 75$ Hz.

The best-predicted bosonic mercury clock transition corresponded to ^{198}Hg [37,38], motivating our focus on this specific isotope. Since ^{198}Hg atoms are trapped in the optical lattice, the lack of prior knowledge of the exact lattice magic frequency [39] posed a potential challenge because of the broadening and contrast reduction of the clock resonance [40].

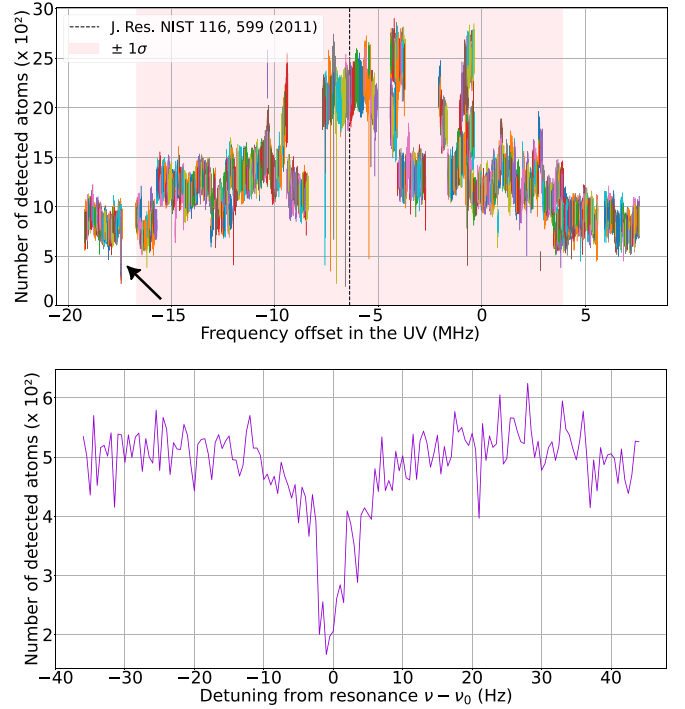


FIG. 2. (Top) All spectra acquired for the search of the ^{198}Hg transition in the optical lattice. Each color corresponds to a scan covering a specific range of 100 kHz by steps of 120 Hz. The step size was carefully selected to avoid the risk of missing the transition without resulting in a prohibitively long overall search duration, considering an expected resonance full width at half maximum (FWHM) of 120 Hz for $\Omega_{12}/2\pi = 75$ Hz. The 100 kHz range was selected for technical reasons related to the need for regular checks of the setup. In total, more than 25 MHz were scanned before successfully identifying the transition (black arrow). (Bottom) Narrowest ^{198}Hg spectroscopy obtained, after optimization: For 200 ms probe time, the FWHM is 4 Hz.

Informed by magic frequency shifts found between Sr and Yb isotopes [27–29,41], ± 2 GHz was considered a conservative upper limit for the expected ^{198}Hg – ^{199}Hg lattice detuning. For our trap depth, the corresponding upper limit for broadening is ~ 26 Hz, which could have hindered the search if the experimental Rabi frequency were significantly lower than expected. To mitigate the risk of unacceptably reducing the transition detection probability, we recorded spectra for the ^{199}Hg fermionic isotope under conditions representative of the search and verified that they showed no significant loss of contrast for lattice detunings up to ± 2 GHz. During the search, the lattice was set to the ^{199}Hg magic frequency. The lattice wavelength is 362 nm and the trap depth was set to $92.4E_r$, with recoil energy $E_r = h \times 7.57$ kHz (setup details in Appendix B). When trapping ^{198}Hg , we measured a MOT loading time constant of 1.2 s and a lattice lifetime of 200 ms, both comparable to those of ^{199}Hg .

Once the experimental setup was ready, we conducted a thorough exploration over the predicted frequency range for the ^{198}Hg bosonic transition [37], scanning a total range of 25 MHz by 120 Hz steps (Fig. 2, top). This meticulous search led to the successful observation of the bosonic transition. Its

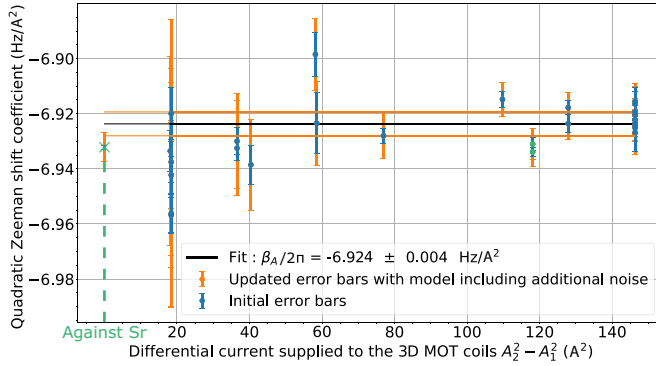


FIG. 3. Measurement of the quadratic Zeeman shift empirical coefficient $\beta_A/2\pi$ with initial error bars (blue) and error bars incorporating a model with additional uncorrelated noise (orange). See text for more details. Blue and green dots represent differential measurements from March–April 2024 and September 2025, respectively. The cross symbol is a value deduced from three direct measurements against a Sr clock at coil currents I_A of 1.1, 3.6, and 10.9 A [42]. The black line represents the estimated coefficient returned by the fit of the measurement values with the modeled error bars. The orange lines represents our final uncertainty range ($\pm 1\sigma$), where σ is defined here as the smallest modeled error bar. The final value $\beta_A/2\pi = -6.924 \pm 0.004$ Hz/A² is specific to our setup and given for reference. Our estimated coil response is 1.53 ± 0.04 mT/A.

frequency lies slightly beyond the $\pm 1\sigma$ uncertainty of Ref. [37] and beyond $\pm 2\sigma$ of Ref. [38].

The measured Rabi frequency of the transition was actually $\Omega_{12}/2\pi = 45$ Hz. Optimizing parameters (magnetic field reduced by a factor 10 and a lower probing power) allowed a probing time of 200 ms, giving spectra with a 4 Hz full width at half maximum (Fig. 2, bottom). Operating a clock sequence with such parameters and comparing against local Sr clocks, we infer stabilities of $\sim 6 \times 10^{-16}/\sqrt{\tau}/s$ for the ^{198}Hg bosonic clock.

III. PROPERTIES OF THE CLOCK TRANSITION

Having identified the transition, we measured key systematic effects to characterize its properties. First, a high magnetic field is required to couple the transition. Although there is no linear Zeeman shift due to zero nuclear spin, a significant quadratic Zeeman shift (QZS) is expected and observed, necessitating accurate measurement. The QZS is given by $\Delta_B = \beta B^2$. In a given setup and within limits detailed next, B can be approximated as proportional to the current I_A flowing through the coils. We can therefore define an empirical coefficient β_A such that $\Delta_B = \beta_A I_A^2$. We measured β_A using an intertwined differential method: Two alternating configurations, differing only in the applied magnetic field, produce different QZS from which we extract a differential shift value.

To achieve the required accuracy in determining β_A , we carefully characterized the current source generating the applied magnetic field. Nonlinearities in current control were accounted for, transients from rapid switching were mitigated via optimized timing, and long-term reproducibility was verified through current monitoring (yielding a 10 μA uncertainty). In Fig. 3, we show the measured QZS coefficient and

TABLE I. Theoretical and experimental (this work) values along with their uncertainties for the coupling coefficient α , the QZS coefficient β , and the probe light shift coefficient κ . These values are discussed in the main text.

	$\alpha/2\pi$ ($\frac{\text{Hz}}{\text{T}\sqrt{\text{mW}/\text{cm}^2}}$)	$\beta/2\pi$ ($\frac{\text{MHz}}{\text{T}^2}$)	$\kappa/2\pi$ ($\frac{\text{mHz}}{\text{mW}/\text{cm}^2}$)
Basic theory	54	−2.5	0.38
Experiments	45 ± 9	-3.0 ± 0.2	2.7 ± 0.6

its statistical uncertainty (blue points and error bars) against the difference of squared currents. Most data points were taken during an initial round, supplemented by three points measured 20 months later. Their average gives the mean value and uncertainty. However, residual inconsistencies in the QZS data were observed in this initial dataset, with the two most inconsistent points differing by 6.3σ . To obtain a reasonable uncertainty despite the unknown origin of these fluctuations, we modeled the dispersion by adding uncorrelated noise and adjusted the error bars accordingly to reflect the total noise [43], as depicted in Fig. 3. We obtain our empirical QZS coefficient β_A with a fractional uncertainty of 5.8×10^{-4} . For the lowest value of the current we used so far (1.13 A), the corresponding fractional uncertainty on the QZS correction is 4×10^{-18} .

Uncertainties arising from magnetic field fluctuations and residual magnetic field also need to be addressed. With a magnetic field probe, we detect fluctuations of 1 μT . The residual field is minimized to less than 1.5 μT . These fluctuations were estimated to have a negligible impact on the measurement of β_A .

Finally, to compare the QZS with its theoretical value (Table I), we convert the empirical coefficient to β using the estimated coil response, whose fractional uncertainty (2.6×10^{-2}) is constrained by our limited knowledge of the coil geometry.

Another essential parameter of the bosonic isotope probing scheme is the magnetic-field-induced coupling coefficient α , describing the coupling between ground and excited states via the external applied magnetic field. To determine α , we measured Rabi frequencies at constant probe intensity I while varying B , yielding a linear relationship whose slope gives α . We accounted for experimental factors affecting the measured Rabi frequencies to ensure accuracy. The 64 ± 5 μm probe waist is comparable to the atomic distribution size, causing the atoms to experience an inhomogeneous probe beam intensity and hence different Rabi frequencies, leading to damped Rabi oscillations. We used a simplified model that averages the Rabi oscillation over the transverse atomic distribution to estimate a factor linking the first maximum of the measured Rabi flopping curve to the Rabi half-period in the ideal case of an atom at the center of the probe beam, thereby seeing the peak intensity. We find a correction factor of 1.3, in line with strong inhomogeneity. The resulting value of α is given in Table I.

Another crucial parameter to measure for an optical lattice clock is the magic frequency, the nonperturbing lattice frequency at which the light shift is canceled at any trapping depth [39]. The magic frequency is an intrinsic atomic property that varies between species and isotopes; therefore,

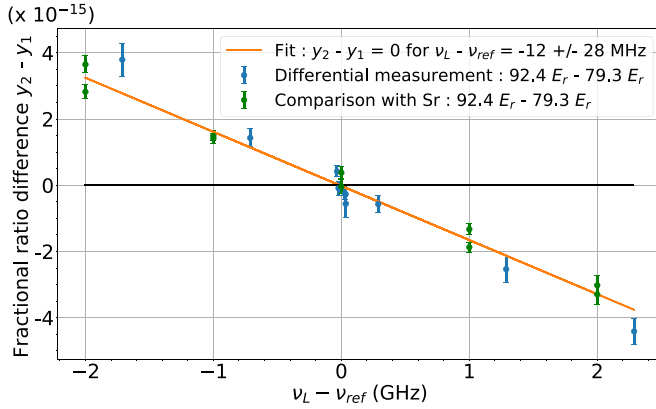


FIG. 4. Magic frequency determination from two series of differential measurements: one with the Hg clock alone and one assisted by a Sr clock [45,46] to improve stability. The reference used for the lattice frequency is $\nu_{\text{ref}} = 826\,855\,240$ MHz. The linear fit (orange) returns the value of the magic frequency $\nu_L^m = 826\,855\,228 \pm 28$ MHz for ^{198}Hg . It also gives the experimental slope of $\eta = (-1.41 \pm 0.04) \times 10^{-4}$ Hz/ E_r /MHz. Both are used to estimate the lattice light shift correction and its uncertainty.

its precise value must be determined for the ^{198}Hg bosonic isotope. To measure the ^{198}Hg magic lattice frequency, we performed differential measurements between two different lattice depths: $92.4E_r$ and $79.3E_r$. Figure 4 shows the differential clock shift as a function of the lattice frequency ν_L . The lattice frequency was precisely known and stabilized using an optical frequency comb with the setup described in Ref. [44].

Given our limited trapping depth range and our frequency resolution, a linear approximation of the lattice light shift is justified to analyze our measurements. In this approximation, the lattice light shift (LLS) is proportional to both the lattice depth (in units of recoil energy E_r) and the detuning of the lattice laser from the magic frequency. For a configuration with lattice depth U_x ($x = 1, 2$), the resulting fractional frequency shift is $y_x = \eta(\nu_L - \nu_L^m)U_x$, where η is the experimental slope of the linear LLS. The magic frequency ν_L^m is found at the detuning where the differential lattice light shift is equal to zero. From this analysis, we extract a magic frequency of $\nu_L^m = 826\,855\,228 \pm 28$ MHz, as shown in Fig. 4. We note that the corresponding uncertainty (Table III) is set by limitations of our setup and not by any fundamental limitation associated with ^{198}Hg .

We also studied the cold collisional shift (CCS) that arises from interactions between cold atoms. To handle this shift, we estimate the atom number during interrogation N_{atoms} by integrating the decay of the atom number within the lattice over the probe pulse duration. This estimation is based on the atom number detected during the subsequent detection phase. The maximum effective atom number during interrogation is $\sim 10^3$, corresponding approximately to one atom per lattice site at the trap center. Atom number was varied by turning on the MOT beams later in the sequence while keeping other key parameters unchanged, such as cycle time and MOT loading time, ensuring a method that is as differential as possible. For differential atom numbers up to 550, we observed an absence of shift well within our typical resolution of 10^{-16} .

In these conditions, it is reasonable to analyze these measurements under the simplifying assumption that the fractional CCS is proportional to N_{atoms} . Doing so, the CCS found is $(-3 \pm 14) \times 10^{-20} \times N_{\text{atoms}}$.

The probe light itself induces an ac Stark shift of the clock levels resulting in a shift of the transition frequency proportional to the probe intensity $\Delta_L(\text{rad/s}) = \kappa I$. To determine κ , we measured frequency shifts against Sr clocks [45,46] while varying the probe power P . We made pairs of measurements at P and $P/2$ while keeping the probe duration, the magnetic field, and all other parameters of the sequence the same. This came at the cost of using a nonoptimal probe pulse in both cases, one stronger and one weaker than the optimum. We also made measurements with optimal pulses. Overall, we explored a power range from 1 to 500 μW , corresponding to a central beam intensity I experienced by the atoms ranging from 15 to 7.9 W/cm^2 .

The probe intensity exhibits significant inhomogeneities across the atomic sample that must be taken into account. To do so, we expand the simple model used for extracting α to include the transverse distribution of detunings caused by the probe light shift. For optimal pulses, the model shows that these inhomogeneities reduce the effective clock shift to about $0.76 \times \kappa I$. For P - $P/2$ measurements, the predicted slope is smaller by an additional factor of 0.9, and their linear extrapolation to $P = 0$ retains a bias of ~ 0.06 times the shift at P . Experimentally, we measure slopes of 33 $\text{mHz}/\mu\text{W}$ for optimal pulses with P in the 10–100 μW range. For the P - $P/2$ method, we obtain slopes from 11 to 29 $\text{mHz}/\mu\text{W}$ for P in the 50–500 μW range, with a mean of 22 $\text{mHz}/\mu\text{W}$. This is not fully consistent with the predicted ratio of slopes between the two methods. We further consider that the beam shape could significantly deviate from a Gaussian profile and that its alignment with the atomic sample may be imperfect or fluctuate over time. As a result, 20% is a reasonable relative uncertainty for our determination of κ given in Table I, which is derived from measurements with the optimal pulses.

Together with experimental values, Table I provides theoretical values of α , β , and κ computed according to Ref. [25]. In addition to the dominant 3P_0 to 3P_1 coupling, we take into account the contribution of 3P_0 to 1P_1 . For the induced coupling coefficient α , 1P_1 accounts for about 10% of the total effect. For the QZS coefficient β , the influence of 1P_1 is found negligible (about 0.1%). The probe light shift coefficient κ is provided for reference only, as the simplified model does not fully account for the polarizabilities of the 1S_0 and 3P_0 clock states. Consequently, it fails to accurately capture the probe light shift. The experiment and theory are in agreement on α , within the rather large experimental uncertainty arising from probe beam inhomogeneities. The experimental value of the QZS coefficient β deviates from theory by 20%, corresponding to a 2.5σ discrepancy. The uncertainty in the applied magnetic field dominates and should be reduced to enable further investigation of this relatively large deviation. As for κ , only our experimental value is relevant, given the previously described weaknesses related to the theoretical value.

Reference [25] defines a figure of merit ξ for the induced mixing scheme. ξ quantifies the trade-off between the coupling strength, which we aim to maximize, and the undesirable

TABLE II. Figure of merit for several bosonic species considered for optical lattice clocks [25], including ^{198}Hg (this work).

	^{198}Hg	Sr	Yb	Ca	Mg
ξ	0.60	0.30	0.60	0.28	0.28

field-induced shifts. Specifically, the Rabi angular frequency can be expressed in terms of the probe light shift Δ_L and the quadratic Zeeman shift Δ_B as $\Omega_{12} = \alpha \sqrt{|\frac{\Delta_L}{\kappa} \frac{\Delta_B}{\beta}|} = \xi \sqrt{|\Delta_L \Delta_B|}$, where $\xi = \alpha / \sqrt{|\kappa \beta|}$ effectively represents the excitation strength relative to the induced shifts. Table II gives our best estimate of ξ for the $^1\text{S}_0 - ^3\text{P}_0$ transition in ^{198}Hg deduced from our work. Hg stands as the most favorable species, together with Yb, among several other considered species.

IV. MEASUREMENT OF THE $^{198}\text{Hg}/^{87}\text{Sr}$ OPTICAL FREQUENCY RATIO

We operated the bosonic mercury lattice clock in two comparison campaigns against the local strontium clocks: a preliminary run in June 2024 and a second after major operational optimization in September 2025. In total, we accumulated about 40 h of data. The systematic uncertainty budget is shown in Table III. The blackbody radiation values rely on theoretical polarizabilities given in Ref. [13]. Background gas collision values rely on Ref. [47] and its practical implementation [48], applied to Hg [35]. The ^{87}Sr counterparts have negligible uncertainties of 4.6×10^{-17} . We obtain a frequency ratio of $^{198}\text{Hg}/^{87}\text{Sr} = 2.629\,315\,734\,684\,118\,1$ with a fractional uncertainty of 6.9×10^{-16} . From this measured ratio and the CCTF 2021 recommended frequency and recommended uncertainty for ^{87}Sr [49], we derive the frequency of the bosonic ^{198}Hg clock transition: $\nu_{^{198}\text{Hg}} = 1\,128\,575\,945\,288\,666.3 \pm 0.78$ Hz.

TABLE III. Typical systematic effect corrections and uncertainties of the ^{198}Hg clock during the campaign for determining the $^{198}\text{Hg}/^{87}\text{Sr}$ optical frequency ratio. The QZS is given for a current of 3.6 A (field ~ 5.5 mT, highest value applied during our frequency ratio determination), and the probe light shift for a power of 100 μW ($I \sim 1.6$ W/cm 2). During the campaign, data were acquired at these values or lower.

Effect	Correction ($\times 10^{-17}$)	Uncertainty ($\times 10^{-17}$)
Quadratic Zeeman effect	7868.7	4.9
Magnetic field fluctuations	0	2.9
Residual magnetic field	0	4.3
Cold collisional shift	4.1	17.3
Linear lattice light shift	0	32.3
Nonlinear lattice light shift	2.1	2.1
Blackbody radiation	14.6	1.5
Background gas collisions	3.0	3.0
Probe light shift	-291.8	58.4
Total	7600.6	69.4

V. CONCLUSIONS

The probe light shift dominates the uncertainty budget (Table III), mainly due to probe beam inhomogeneities and the difficulty of accurately controlling the probe intensity. This uncertainty could be further reduced using a hyper-Ramsey interrogation scheme, which effectively cancels field shifts and their related uncertainties [50,51]. This method has been shown to suppress probe shifts by factors up to 2×10^3 [52] and even 10^4 [53]. Dealing with large and potentially limiting QZS will rely on extending the probe duration to ~ 1 s, reducing the required coupling (and magnetic field) by up to 10 and the QZS by up to 100. Moreover, with probe duration of 1 s, it becomes feasible to implement hyper-Ramsey sequences that shall suppress the QZS together with the probe light shift, reducing their combined uncertainties below 10^{-18} . Reaching this regime requires to use state-of-the-art ultrastable lasers [1–5] and to extend our lattice lifetime [54]. We note that the quenching of the $^3\text{P}_0$ does not impose a limitation [55]. Regarding the CCS, we point out that both fermionic and bosonic clocks have shown comparable nonzero shifts (see, e.g., Refs. [28,56] for Sr). Conditions for fermions to fully benefit from the suppression of the CCS through Pauli blocking are not easily met in practice. CCS in ^{198}Hg should therefore not be considered *a priori* limiting, but rather carefully studied and accounted for. Based on these considerations, ^{198}Hg has the potential to realize an optical lattice clock that fully leverages its low blackbody radiation sensitivity to achieve uncertainties below 10^{-18} and address the objectives outlined in the introduction. This potential extends to other bosonic isotopes of Hg that may be investigated in the future.

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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.

APPENDIX A: ULTRASTABLE LIGHT FREQUENCY SYNTHESIS

We describe the ultrastable light synthesis setup used in our experiments. We emphasize how we modified our original setup [30,31] to obtain a flexible, ultrastable light source capable of probing the clock transition in any of the mercury isotopes, which lie within a tuning range of ~ 10 GHz.

The original setup was initially designed to address the clock transition of ^{199}Hg and its tuning range was tightly limited because of the bandwidth of the acousto-optic modulator (AOM) used to shift the probe light frequency. To bypass this limits, we installed a new fiber laser with a narrow linewidth, operating at 1062 nm. This laser light is then sent through

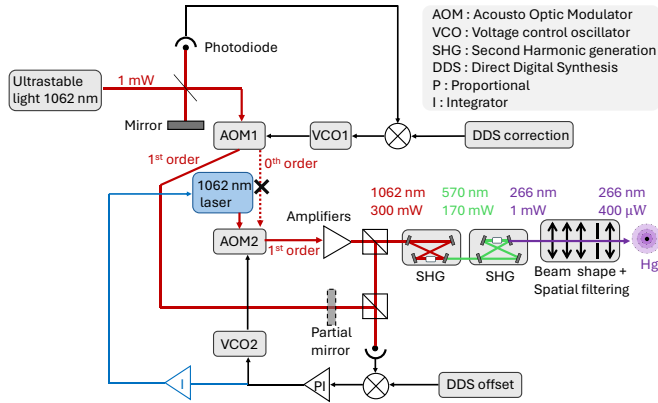


FIG. 5. Simplified schematic of the modified ultrastable light generation. The ultrastable light initially passed through two AOMs (AOM1 and AOM2) before amplification, using a combination of a fibered semiconductor amplifier followed by a Yb fiber amplifier, and then passing through two frequency-doubling cavities. The tuning range was therefore limited to a few tens of megahertz by the AOM2 bandwidth, which is determined by its first-order diffraction efficiency. A 1062 nm laser (blue path) was added and offset phase locked to the ultrastable source using a DDS-controlled offset frequency, bypassing the limitation of AOM2 (dotted arrow with cross: former path). The beam was reshaped after the second second harmonic generation (SHG) stage to improve its Gaussian profile before doing spatial filtering, thus reaching 400 μW at the atoms. This upgrade enables a tuning range of a few nanometers, largely sufficient to access the clock transitions of all mercury isotopes. The figure also shows the phase-locked loop used for fiber noise stabilization, implemented via the DDS correction applied to AOM1.

AOM2 and offset phase locked to the ultrastable light, as illustrated in Fig. 5. Thanks to this new configuration, the frequency of the 1062 nm laser can be tuned as far as we want from the frequency of the ultrastable light, as long as the direct digital synthesis (DDS) offset produces a frequency equal to the detuning between the two lasers.

Also, we optimized the UV beam shape before it reaches the Hg atoms, as the output from the second-harmonic generation cavity was highly non-Gaussian. The beam was reshaped using lenses and spatially filtered through a 75 μm pinhole [35]. We characterized the resulting mode with a CCD camera and estimated the beam waist at the atoms position by reproducing the optical path on a test setup, allowing us to infer the

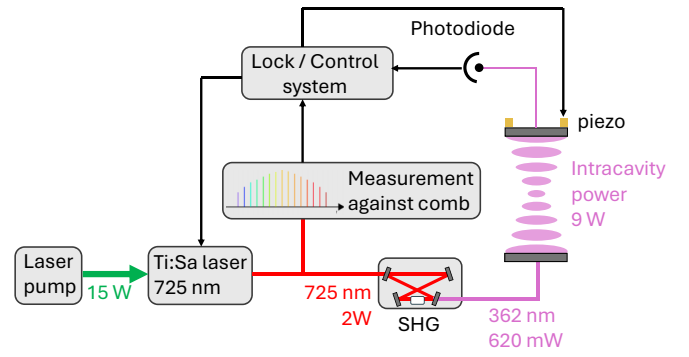


FIG. 6. Simplified schematic of the optical lattice generation. The Ti:sapphire output at 725 nm is frequency doubled in an lithium triborate (LBO) nonlinear crystal and coupled into the lattice cavity, reaching an intracavity power of 9 W. The lattice laser is indirectly locked to the optical frequency comb [44].

beam waist at the atoms as $64 \pm 5 \mu\text{m}$, with a near Gaussian shape. Figure 5 also gives optical powers in our setup.

APPENDIX B: OPTICAL LATTICE SETUP

To trap the atoms, we use a one-dimensional vertical optical lattice, created by a high-intensity laser beam inside an optical Fabry-Perot cavity with a finesse measured to be about 300. The generation of the lattice trap laser light begins with a commercial Ti:sapphire laser pumped with 15 W of single frequency light at 532 nm, producing red light at 725 nm. Then, we sent 2 W of this light through a home-built SHG cavity to generate approximately 620 mW of UV power at 362 nm. The UV light is directed into the vacuum chamber and into the lattice build-up cavity, where the intracavity circulating power reaches about 9 W (see Fig. 6).

We employ the locking scheme described in Ref. [44], where the lattice laser is indirectly locked to an absolutely referenced optical frequency comb via an intermediate laser that is itself locked to the comb and then frequency doubled. This locking method allows us to stabilize the lattice light frequency at any chosen value at or around the magic frequency with an accuracy of better than 1 kHz [44], which is small enough to make the corresponding uncertainty almost negligible (less than 10^{-19}).

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