钟表与 $8 \times 10^{-19}$ 系统不确定性

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我们报告了一种光学晶格钟，其总系统不确定性为 $8.1 \times 10^{-19}$ 在分数频率单位中，代表了任何那一刻的最低不确定性。这种钟依赖于通过测量超窄 $^{1}S_0 \rightarrow {}^3P_0$ 跃迁在冷的原子气体中的能量，用激光稳定在低温中，使得时间保持在一天内漂移几赫兹。利用激光光束，我们以前证明了在三维和二维激光中可以精确地控制原子的光束。在本信中，我们使用这种激光光束来测量黑色体辐射的光谱。这在我们之前的工作中，我们发现这种辐射的稳定性可以通过测量这些激光光束的能量来达到最高的不确定性。这种钟可以被放在原子内的各种位置，包括控制原子的运动和冷却。

钟表与时间测量的重要性

时间测量是物理学中最基本的步骤之一，它在各个领域中都有广泛的应用。目前，最准确的时间测量是通过将不同形式的偶然性转化为精确的测量来实现的。为了实现这一目标，我们需要对不同类型的偶然性进行精确的控制和测量。例如，在原子钟中，我们需要精确地测量原子跃迁的频率，以便实现高精度的时间测量。此外，我们还需要精确地控制原子的冷却和冷却，以便实现高精度的时间测量。为了实现这些目标，我们需要在原子钟中使用高精度的激光器和冷却器，以便精确地控制原子的运动和冷却。在这种情况下，我们需要精确地控制激光器的频率和冷却器的温度，以便实现高精度的时间测量。为了实现这些目标，我们需要使用高精度的激光器和冷却器，以便精确地控制激光器的频率和冷却器的温度，以便实现高精度的时间测量。
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FIG. 1. Overview of the 1D optical lattice clock. (a) Schematic of the system highlights key aspects. Atoms are trapped in a 1D optical lattice formed within an in-vacuum buildup cavity oriented along the direction of gravity \( g \). We read out the state of the atoms by imaging with a 6 \( \mu \)m resolution. Rabi spectroscopy of the clock transition is performed along the tightly confined direction to remain within the resolved sideband regime. The two circular rings are quadrant electrodes for applying an electric field in any direction. A translatable temperature probe with two sensors measures the temperature. (b) 2.43 s Rabi spectroscopy of the two operational clock transitions. We prepare atoms in the \( ^1S_0, m_F = \pm 5/2 \) states, drive the least magnetically sensitive clock transition to \( ^3P_0, m_F = \pm 3/2 \), and measure the excitation fraction \( \rho^{es} \). Blue points are an average of five line scans with error bars given as the standard error. The black line is a Rabi line shape fit. (c) Temperature measured at the atom location over three days. The gold points are an average of the two sensors measured every ten seconds, and the black line is a 20 minute average. (d) Axial blue sideband (BSB) structure. The tilted, shallow lattice creates a set of site-changing Wannier-Stark (WS) transitions where the \( ^1S_0 \) state occupies a different lattice site than the \( ^1P_0 \) state. At our operational depth, only two motional states along the tight confinement direction exist, \( n_Z = 0 \) and 1. In the lower panel, purple points show the \( n_Z = 0 \rightarrow 1 \) transition. The structure of this spectrum is well captured by a model that incorporates the WS structure, the axial structure, and the radial temperature. Each gray line illustrates the BSB for each WS transition, with the sum of these in dashed black consistent with the data. In our standard clock sequence, the lattice depth is briefly reduced to 3 \( E_r \) before readout. The pink points near 0 demonstrate that this approach effectively eliminates all but the \( n_Z = 0 \) population.

To reduce both the lattice light shift and the density shift, it is generally optimal to operate at a shallow lattice depth. Further, we identified a “magic lattice depth” near 15\( E_r \) where on-site interactions are canceled by off-site interactions, leading to a net-zero density shift [18]. Adiabatically ramping the lattice from the loading depth to 15\( E_r \) reduces the radial temperature to \( \sim 120 \) nK. The standard motional sideband model [21] is no longer reliable at these depths. The blue sideband, corresponding with adding one motional quanta along \( \hat{Z} \), splits due to transitions to neighboring lattice sites as shown in Fig. 1(d). Thus it is necessary to use the lattice depth calibration technique introduced in Ref. [19].

To reduce the Zeeman effect sensitivity, we use the least magnetically sensitive \( |^1S_0 m_F = \pm 5/2 \rangle \rightarrow |^3P_0 m_F = \pm 3/2 \rangle \) clock transition as illustrated in Fig. 1(b). Beginning with atoms in one of the stretched states, we apply two clock transfer pulses to prepare atoms in the \( |^1S_0 m_F = \pm 5/2 \rangle \) spectroscopy states with about 96% spin purity.

Between sample preparation and readout, the dead time is 1 s. Based upon a noise model of the clock laser, Dick noise is minimized with a 2.43 s Rabi interrogation time [17]. A digital servo with two integrators tracks the atomic transition by alternating spin states and sides of the Rabi line shape. We expect a single clock stability of \( 5 \times 10^{-17}/\sqrt{\tau} \) for averaging time \( \tau \) in seconds.

As shown in Fig. 1(d), we briefly reduce the lattice depth to the single band regime (\( \sim 3 E_r \)) before readout to ensure that the only measured atoms are in the \( n_Z = 0 \) state. High resolution imaging to readout the clock excitation resolves spatial frequency variation [Fig. 1(a)], allowing for real-time density shift corrections. We operate with approximately \( 4 \times 10^4 \) atoms, leading to a quantum projection noise of \( < 3 \times 10^{-18}/\sqrt{\tau} \), which is near the self-synchronous comparison performance reported in Ref. [16].

**Black body radiation shift.**—The largest systematic shift in room temperature Sr clocks arises from the black body radiation (BBR) environment. The total differential BBR shift \( \Delta \nu_{BBR} \) is the sum of a static component \( \nu_{stat} \) that scales as \( T^4 \), and a dynamic component \( \nu_{dyn} \) that scales with higher powers of \( T \). Thus, for accurate operation we need to determine \( T \) and the atomic response with high precision.

**Radiation temperature:** To ensure a fully thermal environment and measure the radiant temperature at the atoms, we follow a similar technique as in Refs. [4,22]. Two calibrated
thin film platinum resistance sensors are mounted to an invacuum translation arm, which is inserted into the middle of the vacuum chamber. During clock spectroscopy the probe is retracted 30 cm into an auxiliary vacuum chamber. We observe a submillikelvin temperature flicker floor at short timescales, ~2 mK peak oscillations on the hour timescale, and drift of less than a few millikelvin per day, as shown in Fig. 1(d). These few-hour temperature fluctuations are from coupling to room temperature and building process chilled water and can likely be improved by further system isolation. At 12 h the Allan deviation of the temperature is 1.4 mK, which we treat as the operational stability. The total temperature uncertainty is 4.1 mK [23].

\( \nu_{\text{dyn}} \) evaluation: Accuracy in previous generations of room temperature Sr OLCs has been limited by the uncertainty in \( \nu_{\text{dyn}} \), which is directly tied to the 5s4d\(^3\)D\(_1\) lifetime [31]. As in Refs. [22,32], we prepare a sample of Sr atoms in \(^3P_0\) before a 2.6 \( \mu \)m laser pulse excites a portion of the sample to the \(^3D_1\) state. Some excited atoms decay to \(^3P_1\) and then to the \(^1S_0\) ground state, releasing a 689 nm photon in the process, as shown in Fig. 2(a). We collect this fluorescence with a cooled hybrid photomultiplier assembly (PMA) and time tag the incident photons with 5 ns resolution. In the single particle regime, the photon rate \( y \) at time \( t \) is well characterized by a cascaded double exponential process,

\[
y(t) = A \times \Theta(t-t_0) \left( e^{-\frac{(t-t_0)}{\tau_{D_1}}} - e^{-\frac{(t-t_0)}{\tau_{P_1}}} \right) + y_0, \tag{1}
\]

where \( A \) is the flux amplitude, \( \Theta \) is a Heaviside function for instantaneous excitation at time \( t_0 \), \( \tau_{D_1} \) and \( \tau_{P_1} \) are the \(^3D_1\) and \(^3P_1\) lifetimes, respectively, and \( y_0 \) is an offset due to background counts. In Fig. 2(b) we plot all collected photon counts and fit with Eq. (1). Since Eq. (1) assumes instantaneous excitation of atoms to \(^3D_1\), we do not fit data within a 500 ns window about the excitation pulse, indicated by the gray exclusion area [22].

At high densities in the \(^3P_0\) state, we notice a modification to the exponential decay process. Spontaneous emission from one atom can affect the behavior of another atom, leading to effective dipole-dipole interactions and giving rise to effects like superradiance or radiation trapping. The interplay of these effects in this cascaded, multistate decay is hard to simulate theoretically, and we do not have a complete model for extracting single particle lifetimes when such effects are present.

To use the model in Eq. (1) to determine the lifetime, it is vital to keep the population in \(^3P_0\) low as it is the primary state that contributes to collective effects in the \(^3D_1\) decay process. However, reducing atom number adversely affects averaging time. Instead, we load a large number of atoms in \(^1S_0\) and promote a small portion of the atoms to \(^3P_0\). We then excite these atoms to \(^3D_1\) with a 100 ns laser pulse. Since most of the atoms decay back to \(^3P_0\), we repeat this process 15 times before again exciting a portion of the \(^1S_0\) atoms to \(^3P_0\). After 10 clock pulses and a total of 150 \(^3D_1\) decay cycles, we Doppler cool the remaining sample. We repeat this excitation and cooling sequence 5 times before reloading a sample into the lattice. In sum, for each MOT sequence we collect photons from 750 decay cycles. On average we capture less than one photon from the sample per decay cycle, so pile-up effects are effectively eliminated.

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Although the population that contributes to nonlinear collective effects is significantly reduced, weaker collective effects are still present. To understand how density affects the measured lifetime, we measure the total atom number at the beginning and end of the sequence and assign an atom number to each decay event.

Over the course of the measurement campaign, we collect $8 \times 10^7$ photons with a total atom number ranging from $10^7$ to $10^8$. We collect data with different proportions of atoms in $^3P_0$ and iterate over all three hyperfine levels in $^3D_1$, for a total of six separate data sets [23]. We divide each data set by atom number into bins with widths of $5 \times 10^4$ atoms and fit Eq. (1) to each bin. As in Ref. [32], we fit the lifetime density dependence $\tau(n) = \tau_0/(1 + cn)$, where $n$ is the camera measured atom number and $c$ is proportional to density. $\tau_0$ is the single atom lifetime, and $c$ is the density dependence coefficient. At very high density, we notice that the data deviates from the linear model, so we choose to exclude data above $8 \times 10^4$ atoms in these fits. This choice does not change the final reported value. With small population in the $^3P_0$ state, the measured lifetime is shorter at higher density. However, by initially placing 40% of the population in $^3P_0$ the trend is reversed and higher density results in a longer observed lifetime, as shown for the $F = 11/2$ data in Fig. 2(d). The single atom lifetime for these six data sets is plotted in Fig. 2(e).

A number of other systematics can modify the measured lifetime. By applying a magnetic field ($B$) of $\sim 1$ G, we observe Zeeman beats due to interference between the emitted photons. We limit this effect by operating at zero field and periodically measure and correct the background $B < 2$ mG. A similar effect is caused by lattice light shifts splitting magnetic sublevels. To reduce this systematic, we measure the decay in a 10$E_r$ lattice with the PMA oriented within a few degrees of the lattice light polarization. Modeling these effects we assign an uncertainty of 3 ns due to potential Zeeman beats. Using a weighted average of the six experimental conditions, we report $\tau_{D_1} = 2.156(5)$ $\mu$s. This is plotted as the solid black line in Fig. 2(e), with the statistical uncertainty shown as the dashed orange lines and the total uncertainty as the solid orange lines.

This precisely determined $^3D_1$ lifetime allows us to reevaluate $\nu_{dyn}$ based on the technique described in Ref. [33]. Using known atomic properties including measured transition strengths, magic wavelengths, and static polarizability, we determine $\nu_{dyn} = -153.06(33) \text{ mHz at 300 K}$ [23]. The final BBR-related frequency shift combining both the static [34] and dynamic effects at the operational temperature of $20.132(4)^\circ C$ is $(-48417.2 \pm 7.3) \times 10^{-19}$.

**Lattice light shift.**—In previous work [19], we have demonstrated the ability to control the lattice light shift to a few parts in $10^{-19}$. Because of the differential sensitivity to the lattice light shift of the motional states along the tightly confined direction, care must be taken to ensure repeatable cooling of the sample in the lattice. The last stage of cooling is robust and stable, and we take a further step to reduce sample uncertainty by ramping the lattice to $3E_r$ before readout, ensuring that only the lowest band population is measured as shown in Fig. 1(d). With identical atomic coefficients as in Ref. [19], a lattice depth of 15.06(17) $E_r$, and a 10.5 MHz lattice detuning from the measured operational magic frequency [35], the total light shift uncertainty is $3.2 \times 10^{-19}$.

**dc Stark shift.**—Stray electric fields can shift the clock transition frequency [36]. To limit the possibility of patch charges on the mirror surfaces causing these shifts, Faraday shields surround the mirrors and provide passive field attenuation. A pair of in-vacuum quadrant electrodes can apply electric fields in any direction across the atomic sample, shown as copper rings in Fig. 1(a). Alternating high and low fields, we precisely measure the residual dc Stark shift. The shift is below $10^{-21}$ along the cavity direction. The dominant source of residual field is along the imaging axis—likely due to patch charges on the large vacuum window nearby the atoms. The total residual dc Stark shift is $(-9.8 \pm 0.7) \times 10^{-20}$.

**Zeeman shifts.**—Because of the differential Landé-$g$ factor [37] between the $^3S_0$ and $^3P_0$ states, we are sensitive to Zeeman shifts on the clock transition. Probing opposite spin states and taking the frequency average, we broadly reject this systematic. Yet there is still sensitivity to magnetic field fluctuations at and below the experiment cycle frequency. By using the $|S_0m_F = \pm 5/2 \rangle \rightarrow |P_0m_F = \pm 3/2 \rangle$ transition, we substantially reduce coupling to the magnetic environment, even though small field drifts may cause frequency shifts.

We use the 26 times more magnetically sensitive $|S_0m_F = -5/2 \rangle \rightarrow |P_0m_F = -7/2 \rangle$ transition to characterize this effect. Measuring this transition with the same duty cycle as in standard operation, the frequency difference between alternating cycles gives an upper bound on the first order Zeeman shift. We find a flicker floor of 0.78 mHz, leading to a total Zeeman shift uncertainty on the operational transition of $7 \times 10^{-20}$.

Operation on the $|S_0m_F = \pm 5/2 \rangle \rightarrow |P_0m_F = \pm 3/2 \rangle$ transition requires reevaluation of the second order Zeeman coefficient for our desired accuracy goal. This shift $\Delta \nu_{Z2}$ goes as

$$\Delta \nu_{Z2} = \frac{\varepsilon_{\sigma m_F = 5/2}}{c^2} (\Delta_{\text{meas}} - \Delta_{\text{vec}})^2,$$

where $\varepsilon_{\sigma m_F = 5/2}$ is the second order Zeeman coefficient, $\Delta_{\text{meas}}$ is the measured frequency difference between the operational transitions, and $\Delta_{\text{vec}}$ is the splitting due to the lattice vector shift.

To determine $\varepsilon_{\sigma m_F = 5/2}$ precisely, we vary the applied bias field from 0.3 to 1.5 G and measure the resultant frequency shift in an interleaved manner, as shown in Fig. 3. $\Delta_{\text{vec}}$ is measured independently by modulating the lattice depth.
We find a $\xi_{\text{emp.}} = 0.12263(14)$ mHz/Hz$^2$. At the operational field near 380 mG, the second order Zeeman shift $\Delta \nu_{22} = (-85.51 \pm 0.10) \times 10^{-19}$.

Tunneling shift.—At shallow depths, superposition of states in neighboring sites can cause frequency shifts. First identified in Ref. [38], the maximum possible frequency shift due to this effect goes as $\Omega_0 \Omega_1 / \Delta_\mu$, where $\Omega_0$ and $\Omega_1$ are the Rabi frequencies of the carrier and first Wannier-Stark sideband and $\Delta_\mu$ is the frequency difference between neighboring lattice sites. At the operational depth of 15 $E_r$, the off-site Rabi frequency is appreciably large, leading to a maximum shift of $\sim 2 \times 10^{-19}$. While the coherent superposition of neighboring states is likely small, it is difficult to directly measure and control this effect. Instead, we opt to use a Rabi pulse time that is a half integer multiple of the tunneling shift oscillation period [38]. We measure the splitting between neighboring lattice sites to be 867.7461(4) Hz. With a pulse time of 2.4298583 s and a conservative timing uncertainty of 1 $\mu$s, the maximum tunneling shift is $2 \times 10^{-21}$.

Density shift.—Although strong collisional shifts are suppressed by the fermionic nature of $^{87}$Sr, $p$-wave interactions lead to a systematic density shift [39,40]. Previous Sr OLCs required separate evaluation of collisional shifts, leading to drifting systematics as sample preparation varied over time [4]. With imaging synchronously measuring different densities and frequency shifts throughout the sample, we perform real-time density shift corrections. As reported in Ref. [18], operating at a magic lattice depth near 15 $E_r$, on-site $p$-wave and off-site $s$-wave interactions cancel each other, substantially reducing the density shift even with a large atomic sample. For a single characteristic run $\sim 300$ min, the correction is $(-1.1 \pm 0.9) \times 10^{-19}$.

![Graph](image)

**FIG. 3.** Second order Zeeman shift coefficient measurement. We vary the applied magnetic field and measure the splitting ($\Delta_{\text{max}}$) and the frequency shift [see also Fig. 1(b)]. We fit the data with Eq. (2) and plot this fit in green. The lower panel shows the fit residuals with the shaded green region representing the fit uncertainty.

Other systematic shifts.—Collisions between trapped strontium atoms and background gas result in a systematic frequency shift [41]. In this system, the background gas is dominated by hydrogen molecules [23]. As demonstrated in [42], the shift is inversely proportional to the vacuum lifetime. Using this coefficient and a measured vacuum lifetime of $63.6 \pm 2.5$ s, we calculate a background gas shift of $(4.7 \pm 0.5) \times 10^{-19}$.

Line pulling occurs if population in other magnetic sublevels is off-resonantly driven, distorting the carrier line shape. With a low intensity 2.43 s clock pulse, other transitions are highly suppressed. With 96% of the sample in the desired magnetic sublevel, we estimate for the worst case scenario a line pulling shift of $< 10^{-21}$.

Similarly, a low intensity Rabi drive significantly reduces the light shift from the clock laser. Using the coefficient measured in Ref. [43], and accounting for the increased intensity due to both light polarizations, we estimate the probe ac stark shift to be $4 \times 10^{-22}$, which we treat as the uncertainty.

Thermal transients in the acousto-optic modulator (AOM) due to switching may lead to an uncorrected Doppler shift. As in [4], we path length stabilize the same AOM order that drives the atomic transition. The AOM is ramped onto resonance after these thermal transients have settled, leading to an estimated probe chirp shift $< 10^{-21}$.

Summary.—Through precise atomic and environmental control, we have realized a strontium optical lattice clock with a total systematic accuracy of $8.1 \times 10^{-19}$ as reported in Table I. This represents greater than a factor of 2 improvement in systematic accuracy over the previously most accurate strontium optical lattice clock [4], and it sets the accuracy benchmark of all optical clocks reported to date. Black body radiation stands out as the most significant source of uncertainty, and future cryogenic operation should reduce uncertainty to the low $10^{-19}$ level [44].

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