# Differential Light-Shift Cancellation in a Magnetic-Field-Insensitive Transition of $^{87}$Rb

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The precise measurement of transition frequencies of trapped atomic samples is susceptible to inaccuracy arising from the inhomogeneous differential shift of the relevant energy levels in the presence of the trapping fields. We demonstrate near-complete cancellation of the differential ac Stark shift (''light shift'') of a two-photon magnetic-field-insensitive microwave hyperfine (clock) transition in $^{87}$Rb atoms trapped in an optical lattice. Up to 95(2)% of the differential light shift is cancelled while maintaining magnetic-field insensitivity. This technique should have applications in quantum information and frequency metrology.

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Optical trapping of atoms is an indispensable tool for the coherent quantum control of atomic spins, with a range of applications including frequency metrology [1], coherent light storage [2], and quantum information processing in general [3]. However, in many cases coherent quantum control is limited by the presence of differential light shifts (DLS), whereby the relevant transition frequency will vary as a function of the local intensity. The inhomogeneity of the DLS constitutes a major limitation on spin-coherence times, particularly in ensemble samples where an inhomogeneous (and fluctuating) DLS will result in a dephasing of spin coherence. Recent work [4–7] has shown that in some cases the detrimental DLS in ground-state hyperfine levels of alkali-metal atoms can be compensated through the use of local differential Zeeman shifts. This approach has been used to increase spin-coherence times [4,7] and could impact quantum memories and computation, as well as atom-based metrology. In these demonstrations, the DLS cancellation has been obtained at the expense of Zeeman sensitivity to magnetic fields. In this Letter, we experimentally explore the possibility of obtaining simultaneous DLS and magnetic-field insensitivity. This idea has also been investigated theoretically in a recent proposal [8], which suggests that simultaneous cancellation of DLS and Zeeman shifts is possible for certain spin transitions.

The shift $\delta\nu$ of the transition frequency of an optically trapped atomic sample from its free-space, field-free value is determined by a combination of the electronic interaction with light and the electronic and nuclear Zeeman interaction with an external magnetic field, modified by the hyperfine interaction that couples the electronic and nuclear degrees of freedom. The resulting sensitivities $\partial\nu/\partial B$ and $\partial\nu/\partial I$ are functions of the magnetic field $\mathbf{B} = B\mathbf{e}_B$ and the intensity $I$ and polarization $\vec{\epsilon}$ of the trapping light. For ground-state hyperfine transitions of alkali-metal atoms in the absence of light, $\partial\nu/\partial B$ is given by the Breit-Rabi formula [9]. At low field, magnetic-field-insensitive transitions ($\partial\nu/\partial B = 0$) occur for the well-known (single-photon) ''clock'' transitions $|F, m_F = 0\rangle \leftrightarrow |F', m_{F'} = 0\rangle$ at $B = 0$, as well as for multiple-photon transitions $|F, m_F\rangle \leftrightarrow |F', m_{F'} = -m_F\rangle$ at certain nonzero ''magic'' $B = B_m$ (see, e.g., [10]). Here, $F$ denotes the atomic total angular momentum quantum number, and $m_F$ its projection on the quantization axis.

In the presence of light, the energy shifts $\Delta U$ giving rise to $\partial\nu/\partial I$ can be expressed as a sum of a scalar ($\Delta U^s$) and a vector ($\Delta U^v$) component [11]. $\Delta U^s$ is rotationally invariant, depending only on $F$ and $F'$. $\Delta U^v$ depends on $m_F$ and $m_{F'}$ and to lowest order can be formally expressed as a differential Zeeman shift produced by a light-induced effective magnetic field $\mathbf{B}_{\text{eff}} \equiv (E_0/2)^2(\alpha_F^v/2F\mu_F)(i\vec{\epsilon}^* \times \vec{\epsilon})$ [12]. Here $E_0$ and $\vec{\epsilon}$ are the amplitude and (possibly complex) polarization unit vector of the electric field, such that $(i\vec{\epsilon}^* \times \vec{\epsilon})$ represents the direction and relative magnitude of any circular polarization that may be present, $\mu_F$ is the magnetic moment and $\alpha_F^v$ is the (possibly $B$-dependent) vector polarizability. In the limit where $B_{\text{eff}} \ll B$, as is the case in our current experiment where $B_{\text{eff}}/B \approx 0.01$, $\mathbf{B}$ defines the quantization axis and only the component of $\mathbf{B}_{\text{eff}}$ along $\mathbf{B}$ contributes to the shift, so that

$$\frac{\partial\nu}{\partial I} \propto (\alpha_{F'}^s - \alpha_F^s) + \mathcal{A}\left(m_{F'}\frac{\alpha_{F'}^v}{2F'} - m_F\frac{\alpha_F^v}{2F}\right), \tag{1}$$

where $\alpha_F^s$ is the scalar component of the atomic polarizability, and $\mathcal{A} = (i\vec{\epsilon}^* \times \vec{\epsilon}) \cdot \mathbf{e}_B$ represents the degree of circularity of the polarization projected along $\mathbf{B}$.

For alkali ground-state hyperfine transitions, there is no magic condition where $(\alpha_{F'}^s - \alpha_F^s)$ vanishes, implying that the scalar DLS cannot be cancelled with linearly polarized light, where $\mathcal{A} = \mathbf{0}$ [13]. The technique from [4–7] uses the *vector* light shift on a magnetic-field-sensitive transition to cancel the scalar DLS, requiring some component of circular polarization such that $\mathcal{A} \neq \mathbf{0}$.

 

It might appear impossible to use this technique to simultaneously cancel differential Zeeman and light shifts, as canceling the latter requires a nonzero vector DLS. Since the vector light shift acts as an effective magnetic field $\mathbf{B}_{\rm eff}$, we seemingly require a Zeeman sensitivity to $\mathbf{B}$. The key insight is that (to lowest order) light couples only to the *electronic* degree of freedom, whereas the Zeeman interaction includes a contribution from the nuclear magnetic moment. Therefore, the effective magnetic moments for the two interactions differ by roughly the nuclear magnetic moment, so at $B = B_m$, where the differential Zeeman shift cancels, the vector DLS does *not* vanish. The vector DLS can still be used to counteract the scalar DLS.

We can derive approximate expressions for $\Delta U^s$ and $\Delta U^v$ to determine if cancellation is possible. For alkali-metal atoms (with total angular momentum $F$ and $F' = F + 1$ for the two hyperfine ground states), the Zeeman interaction $\hat{H}_Z = -\mu_B(g_J\hat{\mathbf{J}} + g_I\hat{\mathbf{I}}) \cdot \mathbf{B}$, combined with the hyperfine interaction $\hat{H}_{\rm HF} = A_{\rm HF}\hat{\mathbf{J}} \cdot \hat{\mathbf{I}}$, leads to slightly different total magnetic moments for the $F$ and $F'$ hyperfine states, and to field insensitive transitions at $B_m$ (for $m_{F+1} = -m_F$). Within the electric dipole approximation, however, the light interacts only with the electron angular momentum: $\hat{H}_v = -\mu_B g_J \hat{\mathbf{J}} \cdot \mathbf{B}_{\rm eff}$, and the magnitude of the total magnetic moment is the same for $F$ and $F'$. $\hat{H}_v$ differs from $\hat{H}_Z$ by the residual term $\mu_B g_I \hat{\mathbf{I}} \cdot \mathbf{B}_{\rm eff}$, where $g_I$ is the nuclear Landé factor. At $B = B_m$, the vector DLS between states $|F', -m_F\rangle$ and $|F, m_F\rangle$ is $\Delta U^v = -2m_F\mu_N g_I B_{\rm eff}$ [14]. In terms of the total vector light shift $U^v$ of an individual hyperfine level, the differential vector light shift is $\Delta U^v \simeq -2g_I\frac{\mu_N}{\mu_B}U^v$. The differential scalar shift relative to $U^v$ can be estimated as $\Delta U^s \simeq (3\Delta_{\rm HF}/\Delta_F)U^v$ [6]. For $^{87}$Rb, $-2g_I\frac{\mu_N}{\mu_B} = 0.0020$ and $3\Delta_{\rm HF}/\Delta_F = 0.0029$, indicating $\Delta U^v/\Delta U^s$ is on the order of unity, and cancellation is possible. A more detailed calculation, including hyperfine corrections to the electronic wave functions, is needed to accurately determine the degree of cancellation. An accurate calculation for several alkali-metal atoms predicts that full cancellation for two-photon transitions is not possible, but in $^{87}$Rb the DLS sensitivity can still be decreased by a factor of 20 [8].

In the experiment described here, we study DLS of the $|F = 1, m_F = -1\rangle \leftrightarrow |F' = 2, m_{F'} = +1\rangle$ transition of $^{87}$Rb [see Fig. 1(a)] near $B_m = 0.322\,891\,7(3)$ mT [10], by performing high-precision microwave ($\mu$w) spectroscopy of ultracold atoms trapped in an optical lattice. Measuring the transition frequency as a function of light intensity provides the sensitivity $\partial\nu/\partial I$. While it is generally relatively easy to measure the depth of an optical lattice (see [15]), it is more difficult to accurately measure the total light intensity at the position of the atoms. The main reason is the intensity imbalance between the counterpropagating lattice beams due to optical losses and beam area mismatch. To accurately quantify the expected reduction of the DLS, we therefore compare the two-photon $\mu$w transition with the single-photon $\mu$w transition between the $|F = 1, m_F = 0\rangle$ and $|F = 2, m_F = 0\rangle$ states. This transition has a zero vector light shift at $B = 0$.

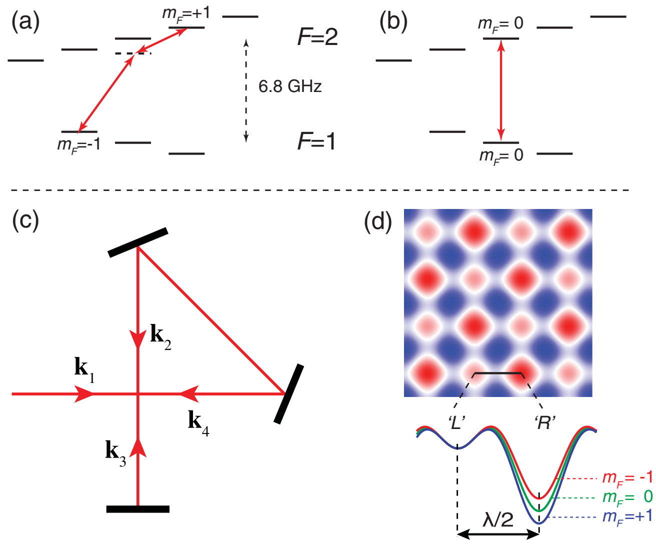


FIG. 1 (color). Schematic of the $^{87}$Rb $\mu$w transitions studied here: (a) the two-photon transition $|F = 1, m_F = -1\rangle \leftrightarrow |F' = 2, m_{F'} = +1\rangle$, used to demonstrate the DLS reduction, and (b) the $|F = 1, m_F = 0\rangle \leftrightarrow |F' = 2, m_{F'} = 0\rangle$ clock transition, used for comparison. Atoms are trapped in an optical lattice formed at the intersection of four laser beams ($\mathbf{k}_1$ to $\mathbf{k}_4$), obtained by a folded, retroreflected single laser beam (c). The intensity pattern used for DLS measurements is shown in (d), along with sections through the unit cell for three different $m_F$ states, and was experimentally optimized in order to maximize the amount of circular polarization on the right ($R$) sites. Atoms are trapped only in the $R$ sites (see text).

Our atomic sample is prepared by loading a $^{87}$Rb Bose-Einstein condensate (BEC) into a 3D optical lattice. We produce BECs of typically $10^5$ atoms, spin polarized in the $|F = 1, m_F = -1\rangle$ state. The final stage of the evaporative cooling is performed in a hybrid trap, created by a focused dipole laser beam at 1550 nm and a quadrupole magnetic field slightly offset vertically with respect to the center of the dipole beam, which provides longitudinal confinement and also compensates gravity (similar to [16]). Subsequently, the atoms are adiabatically ($\approx 200$ ms) loaded into the optical lattice, with a typical depth of about $30E_R$, undergoing the Mott insulator transition [17]. The recoil energy $E_R = \hbar^2k^2/2m_{\rm Rb}$, where $k = 2\pi/\lambda$ is the lattice wave vector, and $m_{\rm Rb}$ is the rubidium mass.

To avoid collisional shifts of the clock transition due to state-dependent on-site interaction (which are of the same order of magnitude as the DLS), it is important to prepare only singly occupied sites. We measured the site occupancy distribution by performing Rabi spectroscopy on the two-photon transition [18]. To control the number of atoms in the BEC (before loading the lattice), we apply a nonadiabatic $\mu$w frequency sweep, removing a controlled fraction of the atoms by transferring them into the untrapped state $|F = 2, m_F = -2\rangle$. We obtain single occupancy when the BEC atom number is $<4 \times 10^4$.

The setup of our optical lattice has been described in detail elsewhere [19,20]. It consists of a 2D lattice in the horizontal ($\hat{x}$-$\hat{y}$) plane and an independent, linearly polarized vertical lattice (along $\hat{z}$). The 2D lattice is obtained from a single, retroreflected laser beam [Fig. 1(c)], and is adiabatically transformed during the experiment from a $\lambda/2$-period lattice with purely linear polarization (used during the loading stage) into the configuration shown in Fig. 1(d) (used for spectroscopy), where the right ($R$) sites have an adjustable circular polarization component. $\mathbf{B}$ was aligned along the resulting $\mathbf{B}_{\rm eff}$ (ideally in the $\hat{x}$-$\hat{y}$ plane), minimizing the DLS sensitivity on the $R$ sites. Based on [19], we developed a procedure to trap atoms only in the $R$ sites: after loading the lattice, we spectroscopically address the atoms on the left ($L$) sites and transfer them to the $F = 2$ manifold, before removing them with a resonant 20 $\mu$s light pulse, which does not affect the $F = 1$ atoms in $R$.

To measure the transition frequencies, we use the detuned Ramsey method, consisting of two $\pi/2$ pulses, separated by a variable hold time $\tau$, with a typical detuning close to 1 kHz (large compared to the DLS, but otherwise arbitrary). We probe the two-photon transition using $\mu$w and rf fields, each detuned by 90 kHz from the intermediate $|F = 2, m_F = 0\rangle$ state [see Fig. 1(a)], resulting in a two-photon Rabi frequency of about 1 kHz. The single-photon transition [Fig. 1(b)] is driven using a single $\mu$w field with a Rabi frequency of 9 kHz. After the Ramsey interrogation, state detection is performed by transferring the atoms between $|F = 1, m_F = -1\rangle$ and $|F = 2, m_F = 0\rangle$ with a $\mu$w $\pi$ pulse, switching off the lattice in $\approx 600$ $\mu$s and absorption imaging the cloud after 18 ms of time of flight and Stern-Gerlach separation.

Figure 2 shows the DLS for both transitions near $\lambda = 806$ nm, as a function of the lattice intensity, expressed as the lattice depth in the corresponding $\lambda/2$ configuration. For circularly polarized 2D lattice light at the $R$ sites, we observe a significantly reduced sensitivity to the lattice intensity of the two-photon transition at $B_m$, compared to the single-photon transition near $B = 0$. The scalar DLS is the same in both cases, and the ratio of the two- and single-photon transition sensitivities quantifies the reduction of the DLS.

The circularity of the lattice light along $\mathbf{B}$ was optimized by minimizing $\partial\nu/\partial I$. From an independent measurement of the optical losses, we estimate that the lattice light has a projected circularity $\mathcal{A} \simeq 0.99$. Moreover, by reversing the direction of $B$, we confirmed an increased DLS sensitivity of the two-photon transition, as in this case the scalar and vector components add together.

To preserve a well-defined quantization axis for the single-photon transition, we maintained a small (20 $\mu$T) magnetic field at which the residual vector DLS is calculated to be $<0.6$ Hz. In a lattice the dependence of $\delta\nu$ on $I$ is not strictly linear, due to a zero-point energy offset [21], but we estimate that for our parameter range the deviation from linearity is smaller than the measurement uncertainties. The zero-point energy contributes to a slight shift in the extracted slopes, of $<4\%$ compared to a traveling wave of equal intensity, but this effect does not contribute to the ratio of the single- and two-photon slopes.

Figure 3 shows the wavelength dependence of the $\partial\nu/\partial I$ ratio between 802 and 815 nm. We observe a local minimum near 806 nm, corresponding to a $\approx 95(2)\%$ DLS reduction, in agreement with theory (Fig. 3). Away from 806 nm, the theory and experiment have the same general trend. The range of wavelengths accessible for this study was limited by two factors: at shorter $\lambda$ the measurement

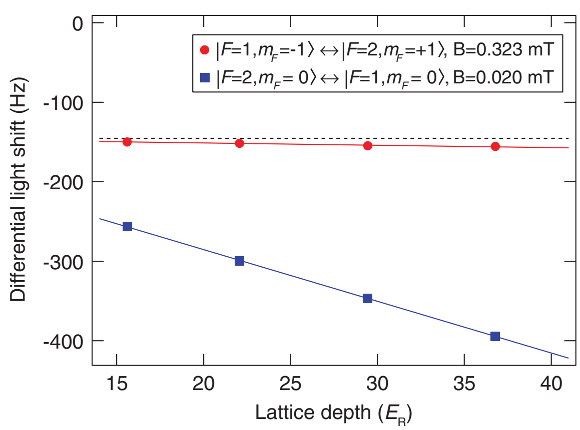


FIG. 2 (color). DLS dependence on intensity, expressed as the lattice depth in the corresponding $\lambda/2$ configuration. We observe a significant reduction of the *total* DLS of the two-photon clock transition (circles), compared to the single-photon clock transition (squares). The full lines represent linear fits to the data. Their extrapolations to zero intensity agree within 5%, and represent the (uncompensated) DLS from the vertical lattice, which was kept at a constant depth of $\approx 30$ $E_R$. The $1\sigma$ uncertainty of each data point is $\approx 1$ Hz.

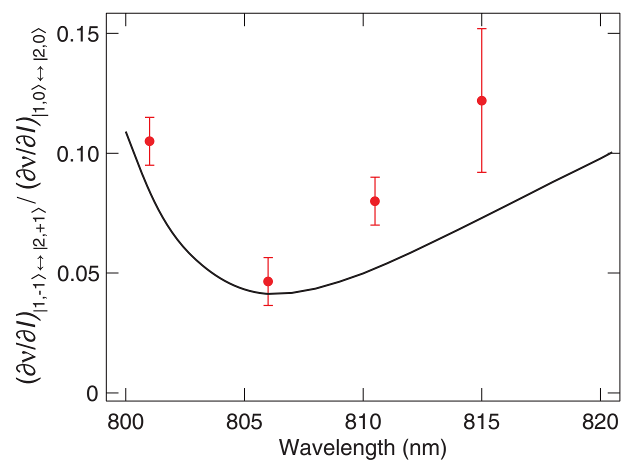


FIG. 3 (color). Ratio between the DLS sensitivities ($\partial\nu/\partial I$) for the two- and single-photon transitions as a function of the lattice wavelength, with a minimum value of 4.5% at 806 nm. The magnetic field was kept near [0.323(3) mT] for the two-photon transition, whereas a small, 20 $\mu$T, bias was used for the single-photon transition. The solid line represents a calculation with no adjustable parameters, following [8]. Error bars represent the $1\sigma$ uncertainty, dominated by statistical errors.

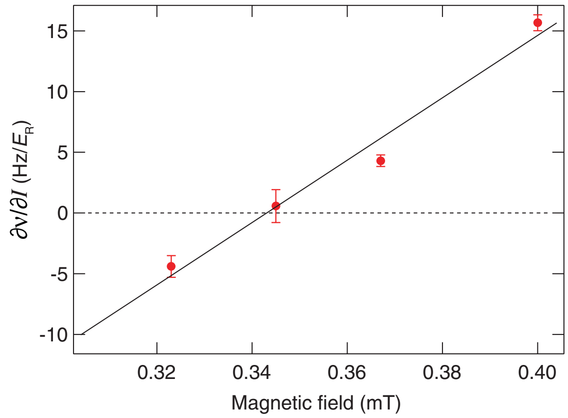


FIG. 4 (color). Sensitivity of the transition frequency to the lattice laser intensity ($\partial\nu/\partial I$), as a function of the bias magnetic field, at $\lambda = 806$ nm. At $B = 0.343(3)$ mT (i.e., 20 $\mu$T away from the magic field), $\partial\nu/\partial I = 0$, and the transition becomes insensitive, at first order, to fluctuations of the light intensity, with a residual field sensitivity $\partial\nu/\partial B$ of 1.7 Hz/$\mu$T.

precision is limited by an enhanced rate of photon scattering from the lattice beams, while at longer $\lambda$ the DLS becomes comparable to our measurement uncertainty.

We also investigated the possibility of achieving $\partial\nu/\partial I = 0$ by slightly shifting the magnetic field away from $B_m$. By measuring the light shift sensitivity of the two-photon transition as a function of the magnetic field at $\lambda = 806$ nm, we observe the linear dependence shown in Fig. 4. The DLS cancels completely at 0.343(3) mT ($\approx 20$ $\mu$T away from $B_m$), in agreement with theory [8]. At this field, the residual magnetic-field sensitivity is $\approx 1.7$ Hz/$\mu$T, about 2 orders of magnitude larger than what is typically used in atomic fountain clocks [22].

We used the detuned Ramsey method, presented above, to compare the sensitivities of the two clock transitions to lattice inhomogeneities. We observed an increase in the coherence time between the two levels of the two-photon transition compared to the single-photon transition. At $\tau = 0$ the Ramsey contrast is 98(2)%; at $\tau = 200$ ms the single-photon contrast has decayed to zero, whereas the two-photon contrast is 10%. This value is likely limited by the inhomogeneous, uncompensated DLS of the vertical lattice, used to support the atoms against gravity.

In summary, we demonstrated a scheme to significantly reduce the light shift sensitivity of an atomic $\mu$w ground-state transition while retaining insensitivity to magnetic-field fluctuations, using a subtle effect which originates in the small difference between the *total* and *electronic* magnetic moments. While simultaneous full cancellation of both differential light and magnetic-field shifts cannot be achieved for $^{87}$Rb, tuning experimental parameters between the differential light and Zeeman shift insensitive points may allow for minimizing the effect of the external field inhomogeneities and fluctuations on the coherence time of trapped atomic samples. The reduced sensitivity, demonstrated here in a 3D lattice, would be even more effective for applications using dipole traps and optical lattices in 1D and 2D geometries [4,7,23].

Our experiments confirm theoretical calculations of DLS [8]. These calculations also predict perfect DLS cancellation at the magnetic-field-insensitive point for four-photon transitions in other alkali-metal atoms ($^{87}$Rb and $^{133}$Cs). In these cases, the difference between the vector and scalar DLS is large enough that cancellation may be possible for arbitrary lattice geometries.

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