Ultracold molecules: vehicles to scalable quantum information processing

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Abstract. In this paper, we describe a novel scheme to implement scalable quantum information processing using Li–Cs molecular states to entangle $^6$Li
and $^{133}$Cs ultracold atoms held in independent optical lattices. The $^6$Li atoms will act as quantum bits to store information and $^{133}$Cs atoms will serve as messenger bits that aid in quantum gate operations and mediate entanglement between distant qubit atoms. Each atomic species is held in a separate optical lattice and the atoms can be overlapped by translating the lattices with respect to each other. When the messenger and qubit atoms are overlapped, targeted single-spin operations and entangling operations can be performed by coupling the atomic states to a molecular state with radio-frequency pulses. By controlling the frequency and duration of the radio-frequency pulses, entanglement can be either created or swapped between a qubit messenger pair. We estimate operation fidelities for entangling two distant qubits and discuss scalability of this scheme and constraints on the optical lattice lasers. Finally we demonstrate experimental control of the optical potentials sufficient to translate atoms in the lattice.

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1. Introduction

The production of scalable, controlled quantum entanglement between many particles would represent a revolutionary breakthrough for information processing. Shortly after Shor’s famous algorithm [1] proved, in principle, that a quantum computer could factor large numbers exponentially faster than any current classical algorithm, there was an exponential growth in the number of proposals for how to implement the essential elements of quantum computation. Since then, many systems have made great strides toward realizing such a computer [2]–[11]; however, truly scalable information processing remains an elusive goal. This is due in part to the stringent requirements on long coherence times, the technical difficulties in implementing high fidelity entangling operations, and the challenge to store and control interactions between many quantum bits (qubits). While neutral atoms provide a natural advantage in coupling weakly to their environment and to other atoms at long distance, atomic interactions at short-range, well described by contact interactions, can be strong, coherent and their effect can be controlled by overlapping the atomic wavefunctions. In particular, the strength of this contact interaction is highly sensitive to underlying molecular structure, and can be precisely manipulated by introducing direct coupling mechanisms between free atoms and molecules.

A system using both ultracold molecules and atoms held in an optical lattice may be a promising system to realize a scalable quantum computer due to the high degree of control available in these systems [12, 13]. Many recent theoretical proposals present schemes to implement entangling operations with neutral atoms in optical lattices [14]–[16] and several experimental groups have demonstrated key steps towards the goal of quantum information processing [17]–[26]. Atoms trapped in optical lattices in particular lend themselves to scalability because thousands of atoms can be isolated in a regular array of micron-sized volumes. Atoms localized in the ground state of each site in the tight-binding regime provide an excellent environment to store quantum information with long coherence times $T_{\text{coh}} > 1 \text{ s}$ [27] and can be spatially transported by controlling the optical phases of the lattice beams [18, 28].

The proposal presented here is a novel approach to use two atomic species, each manipulated by a separate optical lattice potential. Highlighted is the fabrication of lattice structure independent of optical wavelength, use of molecular states to induce entanglement between atoms and introduction of single site addressability without the need for spatially resolved manipulations.

A key aspect of this approach is the introduction of auxiliary messenger atoms used both to probe and to manipulate quantum states and entanglement in an array of qubit atoms. By utilizing two separate species of atom for these two roles and carrying information in their
internal states, it becomes technically feasible to manipulate spatial overlap of atoms and thereby their interactions, without disrupting the sensitive quantum coherences. We propose to use fermionic $^6$Li atoms as qubits, prepared in the lattice with ideally one atom per site. Bosonic $^{133}$Cs will act as messenger atoms to aid in the gate operations and mediate entanglement among the qubits, and will be less densely populated, on average, by one atom per 100 sites of a separate lattice potential of identical structure to the first. By shifting the relative alignment of the lattices through optical phases, each $^{133}$Cs atom can, in principle, be transported to any distant $^6$Li atom; similar schemes can be found in [29]–[33]. Since there may be many $^{133}$Cs atoms, multiple copies of the same computation can proceed in parallel.

2. Scalable quantum information processing with atoms and molecules in optical lattices

The necessary requirements to implement a scalable quantum computer include the ability to initialize the qubit register, fabricate a universal gate set, to have long decoherence times and to read out the information [34]. This section will outline our proposal to meet these requirements.

In recent years, researchers working on neutral atom optical lattice experiments have made great progress obtaining complete quantum control over atoms in a lattice [20]. An optical lattice is the intensity pattern of several interfering laser beams; the resulting periodic pattern can be shaped by varying the intensity, propagation directions, optical phases and polarization of the laser beams. For an effectively two-level atom with far-detuned laser beams, the potential $V(\vec{x})$ is given by $V(\vec{x}) = (\hbar \Gamma/8)[I(\vec{x})/I_{\text{sat}}]/(\Delta/\Gamma)$, where $\Gamma$ is the natural linewidth of the atomic transition, $I_{\text{sat}}$ is the saturation intensity and $\Delta = \omega - \omega_0$ is the laser detuning from resonance at $\omega_0$. $I(\vec{x})$ is the intensity of the optical lattice.

Our scheme to implement a scalable quantum information processor is sketched in figure 1. Two sets of three-dimensional lattices will confine each atomic species independently. One

Figure 1. Scheme for scalable quantum information processing in optical lattices. (a) Qubit atoms (blue dots) form a band insulator in the optical lattices with unity occupancy. Entanglement of two distant qubits can be mediated by the messenger atom (open red circle), which is controlled by a second set of optical lattices (not shown, see the text). (b) Top shows offset intensity profiles for 681 nm light (blue) and 1064 nm (red) light, here $I_{\text{Li}} = 0.24 I_{\text{Cs}}$. Bottom shows resulting potentials for $^6$Li (blue) and $^{133}$Cs (red). (c) Potential energy in the center-of-mass coordinate, including the $^6$Li–$^{133}$Cs interatomic interaction. $|a\rangle$ represents an atomic state, and $|M\rangle$ and $|M'\rangle$ represent molecular states.
lattice at a wavelength $\lambda_1 = 681$ nm will primarily affect confinement of $^6$Li qubit atoms, whose filling ratio will be near unity. A second, less densely populated lattice, at $\lambda_2 = 1064$ nm, will hold $^{133}$Cs messenger atoms to serve as auxiliary quantum bits that allow single site addressing of the qubits, carry entanglement between qubit atoms, and enable readout operations. A sparsely filled Cs lattice with one atom per 100 sites can be prepared by randomly removing the excess atoms via radio-frequency (rf) transitions. The exact position of Cs atoms is unimportant to this scheme, it is only necessary that the Cs lattice has a low filling ratio. The lattices will be fabricated by tuning intersection angles to have identical lattice potential spacings for each species, achieved by using diffractive optical techniques described in section 4. Additionally, one lattice will be physically translatable, to allow controlled contact between qubit and messenger atoms.

The choice of fermionic $^6$Li permits high fidelity initialization of the lattice with one atom per site, achieved by increasing the lattice depth to induce a band insulator state in the atoms [35]. The energy levels for $^6$Li and $^{133}$Cs are shown in figure 2. We propose to use the ground state hyperfine levels $|F = 3/2, m_F = -1/2 \rangle \equiv |1\rangle_{Li}$ and $|F = 1/2, m_F = 1/2 \rangle \equiv |0\rangle_{Li}$, they are denoted by blue circles. In $^{133}$Cs, the qubit levels are the ground state hyperfine clock states, $|F = 4, m_F = 0 \rangle \equiv |1\rangle_{Cs}$ and $|F = 3, m_F = 0 \rangle \equiv |0\rangle_{Cs}$ and are denoted by the red circles. The orange arrows denote the lattice laser wavelengths of $\lambda_1 = 681$ nm and $\lambda_2 = 1064$ nm.

**Figure 2.** Energy level diagrams of $^6$Li and $^{133}$Cs showing the relevant transitions and qubit levels. The qubit levels in $^6$Li are the hyperfine ground states $|F = 3/2, m_F = -1/2 \rangle \equiv |1\rangle_{Li}$ and $|F = 1/2, m_F = 1/2 \rangle \equiv |0\rangle_{Li}$, they are denoted by blue circles. In $^{133}$Cs, the qubit levels are the ground state hyperfine clock states, $|F = 4, m_F = 0 \rangle \equiv |1\rangle_{Cs}$ and $|F = 3, m_F = 0 \rangle \equiv |0\rangle_{Cs}$ and are denoted by the red circles. The orange arrows denote the lattice laser wavelengths of $\lambda_1 = 681$ nm and $\lambda_2 = 1064$ nm.
Figure 3. Plot of limits on the intensities of the 681 nm optical lattice $L_1$ versus the 1064 nm optical lattice $L_2$. The diagonal black lines show the bounds imposed by requiring independent control of $L_1$ over $^6$Li and $L_2$ over $^{133}$Cs. Also shown are the tunneling rate limit and off-resonant scattering rate limit for both $^6$Li (blue lines) and $^{133}$Cs (red lines) for a decoherence rate of $2 \text{s}^{-1}$. The green shaded box shows the available parameter space satisfying all of the above conditions. The black dot corresponds to conditions assumed for calculations in the text.

These states are chosen because they have the same magnetic moment, to minimize decoherence due to external fields.

In $^{133}$Cs, the ground state hyperfine ‘clock’ states $|F = 4, m_F = 0 \rangle \equiv |1 \rangle_{^133}\text{Cs}$ and $|F = 3, m_F = 0 \rangle \equiv |0 \rangle_{^133}\text{Cs}$ will be used (see figure 2). These states are magnetic field insensitive to first order.

Independent control of the qubit and messenger atoms is essential in this setup and can be realized by a careful choice of the lattice laser detunings and intensities. The choice of $^6$Li and $^{133}$Cs is favorable in this sense due to their very different dominant atomic transition lines at $\lambda = 671$ and 852 nm, respectively. This opens up the possibility to independently confine $^6$Li and $^{133}$Cs atoms with two sets of moderately detuned optical lattices $L_1$ and $L_2$.

Figure 3 illustrates the constraints imposed on lattice intensities to maintain independent control of $^6$Li and $^{133}$Cs, and bit lifetime of 500 ms due to off-resonant scattering and tunneling. The lattice spacing was chosen to be 1.5 $\mu$m. By evaluating the maximum force each lattice can exert on each of the two atomic species, we show that the condition for independent control of the atoms by the associated lattices can be expressed in terms of the lattice laser intensities, which should satisfy $0.04 < I_1/I_2 < 1.45$. With the optimal choice of $I_1/I_2 = 0.24$, the dipole force from the farther detuned lattice does not exceed $\alpha = 16\%$ of the force from the nearer detuned lattice. In addition, this plot shows the limits for two primary decoherence sources, off-resonant scattering of lattice light and atomic tunneling. In both cases, tighter constraints are set by the $^6$Li atoms due to their lighter mass and to the smaller detuning of $L_1$ to the $^6$Li transition.

The capability to independently control the two atomic species allows us to have single site addressability of the qubit atoms. This is accomplished by shifting the optical phases of the messenger lattice, allowing the $^{133}$Cs messenger atom to be translated to any $^6$Li qubit atom.
Figure 4. Implementation of the necessary gates via coupling to a molecular state $|M\rangle$. Not only does the molecular state allow entangling operations to be carried out between the atoms, but it also allows single qubit addressing. Part (a) shows how to execute a targeted single qubit rotation. When the atoms are overlapped, rf pulses allow the qubit atom to be rotated into a superposition state. Part (b) shows the sequence to entangle two distant qubits. After entangling messenger and qubit (step 1), the messenger can be transported to a second qubit and subsequently entanglement can be exchanged (step 2).

This is a necessary step for many operations in this proposal, including detection and creation of a universal gate set (see section 3).

There are several possibilities for reading out the quantum information from the qubits. One approach is to use the auxiliary messenger $^{133}$Cs atoms to determine the state of the $^6$Li atoms, allowing information to be obtained without disturbing the qubit lattice. A second approach is to readout the state of the qubit lattice directly. This could be done by using high numerical aperture, state-selective imaging of the qubit lattice directly, for example, see [22].

The most fundamental requirement for realizing a quantum information processor is the need for a universal gate set. This includes both single qubit rotation gates and multi-qubit entangling operations. One possible way to achieve a universal gate set in this system is shown in figure 4. In both cases, the logic states of the atoms can be coupled to a $^6$Li–$^{133}$Cs molecular state through the use of rf fields.

To perform targeted qubit rotations, where only a single qubit is rotated and the neighboring qubits remain unaffected, we will translate the messenger atoms by shifting the $L_2$ lattice in
order to overlap the messenger with the target qubit. When the two atoms are overlapped, they can be coupled to a molecular state using rf transitions and, depending on the frequency and duration of the rf pulses, we can perform any arbitrary Bloch rotation (see figure 4(a)). Global rotations of all qubits can be realized with microwave pulses.

Figure 4(b) shows a possible protocol to entangle two distant qubits. In the first step, the messenger, prepared in a superposition state, is brought to the first qubit and entangled by rf pulses as shown in step 1 in figure 4(b). Next, the messenger atom is translated to the second qubit and the quantum entanglement is swapped between the messenger and qubit, as shown in step 2 in figure 4. This leaves the qubit atoms entangled with each other and the messenger atom disentangled from the qubits. The overall evolution of the quantum states for entangling two distant qubits, Li\textsuperscript{a} and Li\textsuperscript{b}, via a messenger Cs atom is given as

\[
|\text{Cs : 0 + 1} \rangle \otimes |\text{Li}^a : 0 \rangle \otimes |\text{Li}^b : 0 \rangle \quad \text{Step 1} \quad (-|01\rangle + |10\rangle) \otimes |0\rangle
\]

\[
\text{Step 2} \quad -|010\rangle - |001\rangle = -|0\rangle \otimes (|10\rangle + |01\rangle).
\]

(1)

3. Entanglement via quantum states of ultracold \textsuperscript{6}Li–\textsuperscript{133}Cs molecules

Molecular states are excellent candidates to induce entanglement of atoms because the molecular potential, in general, depends on atomic spin. In this section, we evaluate the times and fidelities to induce single-spin rotations and entanglement of qubits via radiative transitions to molecular states as shown in figure 4(b).

There are two distinct ways in which two free atoms can couple to the molecular states. First, coupling to deeply bound \textsuperscript{6}Li–\textsuperscript{133}Cs molecules can be induced by direct radiative (electric dipole D1) transitions. Alternatively, coupling to weakly bound molecules near the continuum can be realized using rf or microwave transitions (magnetic dipole M1). In the following paragraphs, we will focus on the magnetic dipole transitions.

In the small binding energy limit, the Rabi frequency for magnetic dipole transitions between a free-atom state and a molecular state can be estimated as \( \Omega = \Omega_0 C \), where \( \Omega_0 \) is the typical Rabi frequency for the transition for free atoms, and the Franck–Condon factor \( C = \int \psi_a(r) \psi_m(r) \, dr \) is given by the wavefunction overlap of the atomic state \( \psi_a(r) \) and the molecular state \( \psi_m(r) \).

To address a \textsuperscript{6}Li atom in the lattice, the \( L_2 \) is translated to bring a \textsuperscript{133}Cs atom into wavefunction overlap with \textsuperscript{6}Li. The Hamiltonian of the two atoms in a lattice site is characterized by one \textsuperscript{6}Li with mass \( m_1 \), momentum \( p_1 \), position \( r_1 \), trap frequency \( \omega_1 \) and one \textsuperscript{133}Cs atom with \( m_2, p_2, r_2 \) and \( \omega_2 \). The Hamiltonian in the lab coordinate and in the center-of-mass frame are given by

\[
H = \frac{p_1^2}{2m_1} + \frac{p_2^2}{2m_2} + \frac{m_1 \omega_1^2 r_1^2}{2} + \frac{m_2 \omega_2^2 r_2^2}{2} + V(|r_1 - r_2|)
\]

(2)

\[
= \frac{P^2}{2M} + \frac{P_2^2}{2\mu} + \frac{M \omega_1^2 R_1^2}{2} + \frac{\mu \omega_2^2 R_2^2}{2} + V(r),
\]

(3)

respectively, where \( M = m_1 + m_2 \) is the total mass, \( \mu = m_1 m_2 / (m_1 + m_2) \) is the reduced mass, \( P = p_1 + p_2 \) is the total momentum, \( p = \mu (p_1/m_1 - p_2/m_2) \) is the relative momentum, \( r = |r_1 - r_2| \) is the atomic separation, \( \omega_1^2 = (m_1 \omega_1^2 + m_2 \omega_2^2) / M \) is the center-of-mass trap frequency and \( \omega_r = \omega_1 \omega_2 / \omega_b \) is the relative motion trap frequency. Potential energy in the center-of-mass frame is shown in figure 1(c).
In the relative coordinate, the ground state wavefunction of a weakly interacting atom pair is given by \( \psi_s(r) = (r_0^2 \pi)^{-3/4} \exp(-r^2/2r_0^2) \) and \( r_0 = (\hbar/\mu \omega_0)^{1/2} \) is the oscillator length. The wavefunction of a weakly bound molecular state is \( \psi_m(r) = (2\pi ar)^{-1/2} \exp(-r/a) [36] \), where we assume that the \(^6\text{Li}^{133}\text{Cs} \) scattering length \( a \) is larger than the interaction range, but small compared with the trap oscillator length. We can then evaluate the integral and obtain the total operation time of \( \tau \) operations require two \( \pi \) pulses on the atom–molecule transition, and consequently require a total operation time of \( \tau = \pi/\Omega = 2.5 \text{ ms} \).

The fidelity of the above operations can be estimated from the uncertainty of the coupling strength and off-resonant population transfer to other states. One major source of the coupling strength variation comes from the imperfect overlap of the \(^{133}\text{Cs} \) and \(^6\text{Li} \) ground state wavefunctions, particularly when the lattice site positions are not perfectly controlled. In section 4, we show that the relative lattice positions can likely be controlled to about \( \delta = 10 \text{ nm} \) in the near future. By evaluating the wavefunction overlap with such an offset, we derive the resulting fidelity \( \mathcal{F} \), defined as the square of the overlap between target and actual output states, to be \( \mathcal{F} = \exp(-\delta^2/r_0^2) = 99.5\% \) per operation. For both spin rotations and \(^6\text{Li}^{133}\text{Cs} \) entanglement operations, which require two atom–molecule transitions (see section 2), we expect the overall fidelity to be 99%. The dominant off-resonant population transfer will occur when the atoms are excited to unintended molecular states, or the molecule is converted into atoms in other vibrational states. From a two-level model, the off-resonant population transfer is given as \( \delta p = (1 + \Delta^2/4\Omega^2)^{-1/2} \) per \( \pi \)-pulse, where the smallest detuning is determined by the lattice vibrational energy of \( \Delta = \omega_i r_0^2 \) kHz, which suggests \( \delta p = 0.2\% \). The loss in fidelity from off-resonant excitation is thus likely smaller than that from the lattice misalignment.

Entanglement of two qubits requires time not only to entangle \(^6\text{Li} \) and \(^{133}\text{Cs} \) at a lattice site, but also to transport the cesium atoms to a third, distant, \(^6\text{Li} \) atom and to transfer entanglement (see section 2). The entanglement transfer process requires two atom–molecule \( \pi \)-pulses, which take 2.5 ms with a fidelity of 99%, similar to the \(^6\text{Li}^{133}\text{Cs} \) entanglement gate. In the following, we estimate the time required to adiabatically transport a cesium atom over \( N \) lattice sites with fidelity \( \mathcal{F} = 99\% \).

We will adopt an adiabatic transportation process to keep messenger atoms in the vibrational ground state while moving them. The leading order loss of quantum information comes from the population transfer from the ground state \( |0\rangle \) to an excited state \( |i\rangle \). Based on an adiabatic approximation [37], we estimate

\[
p_1 = \frac{1}{\hbar^2 \omega_0^2} \left| \frac{\partial H(t)}{\partial t} |0\rangle e^{i\omega_0 t} \right|^2 \rho(E)\omega_0 t, \tag{4}
\]

where \( \rho(E) = 1/\hbar \omega \) is the density of states in the direction of motion, and the time-dependent \( H(t) = U^\ast \sin^2 k(x - vt) \) in the frame moving with the cesium atoms is mostly due to the \( L_1 \) lattice potential, \( U^\ast \) is the maximum \( L_1 \) lattice depth experienced by the cesium atoms and \( v \) is the velocity of \( L_1 \) relative to \( L_2 \). When a cesium atom moves over \( N \) lattice sites, we can express the population transfer to the lowest excited state \( i = 1 \) as

\[
p_1 = N \frac{\pi}{2} v \frac{U^\ast}{\hbar \omega^2} k^2 x_0^2 e^{-k^2 x_0^2}, \tag{5}
\]

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where we have introduced the reduced velocity $v = v \pi \omega / k$. Using the parameters described in section 2, $U^e = \alpha \hbar \times 760 \text{kHz}$ with the cross-talk suppression factor $\alpha = 0.16$, $k = \pi / d$ and $d = 1.5 \mu\text{m}$, we derive, for a fidelity of 99%, $v = 0.03/N$, which corresponds to an upper bound of the cesium atomic velocity of $v \approx (4/N) \mu\text{m ms}^{-1}$ when it is required to move over $N$ sites.

To summarize, entanglement of two distant $^6\text{Li}$ atoms separated by $N$ lattice sites involves four atom–molecule transitions and one lattice transport. We conclude that the total time needed to perform the whole sequence is $\tau_c = (5 + 0.4N^2) \text{ms}$ with an overall fidelity of 97%. Since the total number of accessible qubits within a travel distance equivalent of $N$ lattice sites is $N_q = (4\pi / \sqrt{3})N^2$, the time needed to realize a single pairwise entanglement gate in a system of $N_q$ qubits is $\tau_c = (5 + N_q/20) \text{ms}$ on average. The weak dependence on $N_q$ confirms that our entanglement scheme is indeed scalable to many qubits.

4. Controlled overlap of bichromatic optical lattices

In order to perform many identical computations simultaneously, it is necessary to match the lattice constants of the messenger and qubit lattices to commensurate values. The constraints imposed by necessary lattice depth to achieve sufficiently low bit migration, off-resonant scattering rate and independent control of each atomic species, preclude the use of lasers with commensurate wavelengths and the formation of lattices by counterpropagating beam pairs. For this reason, we have tuned the intersection angles and relative beam phases of the beams to match the relative propagation vectors. We have chosen to work with the most simple two-dimensional potential which is topologically stable against changes in relative phase of the constituent beams, consisting of three beams at each wavelength whose $k$-vectors projected onto the plane $\vec{k}_{m\perp}$, where $m$ represents the wavelength, are equal in magnitude $k_\perp$ and distributed evenly on a unit circle (see figure 5). The angle of each wavevector to the normal from the plane is chosen to be $\theta_m = \sin^{-1}(2\lambda_m/3d)$, where $d$ is the common lattice constant. Each wavelength then creates a two-dimensional intensity pattern of the form $I(x, y) \propto 6 - \sum_j \cos^2(\sqrt{3}k_\perp r_j / 2 + \phi_j)$, and $r_j = x \cos(2j\pi / 3) + y \sin(2j\pi / 3)$ for $j = 1, 2$ and 3 are determined by the relative wavevectors of the beams, and $\phi_j$ are determined by the relative optical phase of the beams. Finally, in the third direction, a single standing wave is applied by the intersection of two beams at a small angle, producing a lattice constant similar to that in the plane.

Precise and stable tuning of intersection angles and relative beam phases $\phi_j$ can be achieved by the use of a combination of diffractive and refractive optics. In this scheme, a two-dimensional diffraction grating can be employed as a three-way beam splitter whose output beam angles are dependent on wavelength, with precisely the relation necessary to generate matched lattice constants at arbitrary input wavelengths. These diffracted beams can be mapped onto the location of the atoms using refractive imaging techniques, as shown in figure 5; we note that only three selected spatial frequencies are allowed to propagate through the imaging lenses. By employing only ‘common-mode’ optics, through which all beams at each wavelength pass, a highly phase-stable optical setup can be constructed, largely insensitive to mount vibration and drift. A time series of the minimum location for a two-color lattice is shown in figure 6, as recorded by imaging the lattice onto a CCD with a microscope objective, demonstrating a stability of 92 nm over 3000 s. The differential translational stability is measured to be 26 nm over 3000 s. This is to be compared with the site spacing $d = 1.5 \mu\text{m}$ and anticipated oscillator length of the cesium atom in the lattice of 82 nm, and for lithium 165 nm.

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Figure 5. Apparatus for generating a two-color optical lattice. Co-propagating beams at both wavelengths are incident on a diffractive optical element (DOE) shown in (a), formed by a photolithographed gold-coated fused silica surface consisting of a regular array of raised equilateral triangles. The image shown was obtained with an atomic force microscope. In (b), reflected light is diffracted, primarily into three first-order beams at each wavelength in a triangular pattern. These beams are then routed by a pair of lenses to form a pair of triangular optical lattices shown in (c), on the image plane of the DOE. The relative position of the two lattices is controlled with a set of electro-optic phase modulators, formed by patterned deposition of mirror/electrodes onto the rear surface of a single lithium-niobate crystal. The lattice structure shown was imaged with a microscope objective onto a CCD camera.

To control the relative position of cesium and lithium atoms in the lattices, we insert optical phase modulators to control the relative phases of the beams for at least one wavelength (see figure 5). For this purpose, we have chosen electro-optic phase modulators for their high bandwidth and relative precision. In order to retain as much as possible a common-mode optical setup, we integrate several longitudinal electro-optic modulators into a single, large diameter lithium–niobate crystal wafer by patterning multiple electrodes onto its surfaces (see figure 5). To lower the electrical potential difference necessary to effect a change in optical path length, the modulators are used in double-pass, with the rear electrodes serving also as mirror coatings, realized by deposition of a patterned layer of silver onto the back surface of the lithium–niobate crystal. The front electrodes consist of a single indium–tin-oxide coating held at a common potential. While this type of modulator requires substantial electric fields of order 2 kV mm$^{-1}$ to be generated across the crystal in order to provide displacements of a single lattice constant, this is readily obtainable with existing electronics. At this drive level, one already has access with a single messenger atom to seven lattice sites. With simple modifications to the optics, multipass geometries can be generated to extend the range further.

A major concern in preserving the coherence of atomic internal states is to provide a potential that is independent of the internal state of the atom. However, for detunings not large compared with the fine structure splitting, electron spin is not decoupled from the
Figure 6. Translational stability of the two-color optical lattice potential. The intensity distributions for both color lattices were recorded simultaneously by imaging onto a CCD camera with a microscope objective. In (a), a time series of motion for each color optical lattice is presented, demonstrating a single-color root-mean-square displacement of 92 nm over a 3000 s measurement time. The relative motion is substantially smaller due to the use of common-mode optics, with a rms displacement of 26 nm over the same time period. A typical experiment cycle is expected to require on the order of 10 s. In (b), a power spectrum is presented, demonstrating cancellation of motional noise over a wide range of frequencies.

The effect of the driving optical field, and one must account for the internal state-dependent light-shift introduced by polarization gradients. In the context of this experiment, this has two potentially important consequences. First, it leads to a potential dephasing mechanism in the presence of inhomogeneous light fields. In the limit of low magnetic field, the internal states of free atoms are eigenstates $|F, m_F\rangle$ of the total angular momentum $F$, and the vector light shift takes the form of an effective magnetic field $B_{\text{eff}}$ [38], proportional to the constant $D_{\text{FS}} = (\Delta_{3/2} - \Delta_{1/2})/((\Delta_{3/2}/2 + \Delta_{1/2})$, determined by the detunings $\Delta_{3/2}, \Delta_{1/2}$ of the lattice light from the two excited state fine structure components; for lithium and lattice light at $\lambda = 681 \text{ nm}$, $D_{\text{FS}}^{681 \text{ nm}} = 1.4 \times 10^{-4}$, however, for cesium $D_{\text{FS}}^{681 \text{ nm}} = -0.11$ and $D_{\text{FS}}^{1064 \text{ nm}} = 0.19$. We note that a lattice formed by beams with parallel polarizations will exhibit no such state-dependent light shift. While this is possible in a planar geometry with intersection angles $\theta_m = \pi/2$, smaller intersections lead to a nonzero effective magnetic field. For optimally chosen linear polarizations, $B_{\text{eff}}$ varies in space, exhibiting zero magnitude but nonzero gradient at the location of the scalar potential minimum. Assuming application of a substantially larger external magnetic field in the $z$-direction, only the gradient of the component in this direction is relevant. This is of order the of $1 \text{ kG cm}^{-1}$ for cesium due to each lattice, which leads to a spatial separation of the minima for different internal states orthogonal to the direction of optical polarization, but introduces no shift in ground state energy to lowest order. The peak value of $|B_{\text{eff}}|$ reaches a maximum at the scalar potential minimum of order 200 mG for cesium and 100 $\mu$G for lithium.

The amount of decoherence resulting from the polarization gradients above depends critically on the chosen internal states, as well as on the magnitude of applied magnetic field.
At low fields, it is possible to place cesium atoms only in superpositions of the clock states $|F = 3, m_F = 0\rangle$ and $|F = 4, m_F = 0\rangle$, and lithium atoms in states of equal magnetic moment, in which case we expect to be largely insensitive to deleterious magnetic field and polarization gradient inhomogeneities. It is important to note that the cesium clock states remain good quantum states to relatively high magnetic fields of order 100 G, whereas the chosen lithium states enter the high-field regime at smaller fields of order 10 G. However, the relatively small fine-structure splitting of lithium assures a modest influence of polarization gradients in all cases.

5. Conclusion

We have presented a scheme for scalable quantum information processing based on two species of ultracold atoms held in controlled bichromatic optical lattice potentials, including methods to entangle $^6\text{Li}$ and $^{133}\text{Cs}$ atoms locally through coupling to bound $^6\text{Li}–^{133}\text{Cs}$ molecules, and methods to transport entanglement to distant atoms through multiple quantum manipulations. We have identified simple quantum logic gate operations possible in this scenario. Methods are based on the production of translatable optical lattices at two wavelengths with identical structure, for which we have demonstrated a novel realization utilizing diffractive optics and electro-optic modulation. We have discussed gate operations in detail, identifying necessary timescales for entangling via a molecular state and transporting atoms adiabatically. This compares favorably with the expected coherence time, including the effects of off-resonant scattering, qubit tunneling, external field instabilities and state-dependent light shifts. Finally, we have analyzed the effects of realistic experimental uncertainties to ascertain expected fidelities, and compared this with measured errors in lattice construction; with incremental improvement in passive stability, fidelities of $>97\%$ are achievable in entangling distant qubits.

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References

[1] Shor P W 1997 Algorithm for fast factoring SIAM J. Comput. 26 1484
[2] Monroe C 2002 Quantum information processing with atoms and photons Nature 416 238–46
[3] Blatt R and Wineland D 2008 Entangled states of trapped atomic ions Nature 453 1008–15
[4] Vandersypen L M K and Chuang I L 2005 NMR techniques for quantum control and computation Rev. Mod. Phys. 76 1037–69
[5] You J Q and Nori F 2005 Superconducting circuits and quantum information Phys. Today 58 42–7
[6] Hanson R, Kouwenhoven L P, Petta J R, Tarucha S and Vandersypen L M K 2007 Spins in few-electron quantum dots Rev. Mod. Phys. 79 1217
[7] Raimond J M, Brune M and Haroche S 2001 Manipulating quantum entanglement with atoms and photons in a cavity Rev. Mod. Phys. 73 565–82
[8] Pryde G J 2008 Optical quantum information: the quantum information cocoon Nat. Photonics 2 461–2
[9] Kok P, Munro W J, Nemoto K, Ralph T C, Dowling J P and Milburn G J 2007 Linear optical quantum computing with photonic qubits Rev. Mod. Phys. 79 135
[10] Jessen R, Deutsch I H and Stock R 2004 Quantum information processing with trapped neutral atoms Quantum Inf. Process. 3 91–103
[11] Prevedel R, Walther P, Tiefenbacher F, Bohi P, Kaltenbaek R, Jennewein T and Zeilinger A 2007 High-speed linear optics quantum computing using active feed-forward Nature 445 65–9
[12] Daley A J, Boyd M M, Ye J and Zoller P 2008 Quantum computing with alkaline earth atoms Phys. Rev. Lett. 101 170504
[13] Rabi P, DeMille D, Doyle J M, Lukin M D, Schoelkopf R J and Zoller P 2006 Hybrid quantum processors: molecular ensembles as quantum memory for solid state circuits Phys. Rev. Lett. 97 033003
[14] Jaksch D, Briegel H J, Cirac J I, Gardiner C W and Zoller P 1999 Entanglement of atoms via cold controlled reflections Phys. Rev. Lett. 82 1975
[15] Brennen G K, Caves C M, Jessen R and Deutsch I H 1999 Quantum logic gates in optical lattices Phys. Rev. Lett. 82 1060
[16] Jaksch D, Cirac J I, Zoller P, Rolston S L, Cote R and Lukin M D 2000 Fast quantum gates for neutral atoms Phys. Rev. Lett. 85 2208
[17] Greiner M, Mandel O, Esslinger T, Hansch T W and Bloch I 2002 Quantum phase transition from a superfluid to a Mott insulator in a gas of ultracold atoms Nature 415 39
[18] Mandel O, Greiner M, Widera A, Rom T, Hansch T W and Bloch I 2003 Controlled collisions for multi-particle entanglement of optically trapped atoms Nature 425 937
[19] Peil S, Porto J V, Tolra B L, Obrecht J M, King B E, Subbotin M, Rolston S L and Phillips W D 2003 Patterned loading of a Bose–Einstein condensate into an optical lattice Phys. Rev. A 67 051603
[20] Bloch I 2008 Quantum coherence and entanglement with ultracold atoms in optical lattices Nature 453 1016
[21] Scheunemann R, Cataliotti F S, Hänsch T W and Weitz M 2000 Resolving and addressing atoms in individual sites of a CO$_2$-laser optical lattice Phys. Rev. A 62 051801
[22] Nelson K D, Li X and Weiss D S 2007 Imaging single atoms in a three-dimensional array Nat. Phys. 3 556
[23] Johnson T A, Urban E, Henage T, Isenhower L, Yavuz D D, Walker T G and Saffman M 2008 Rabi oscillations between ground and Rydberg states with dipole–dipole atomic interactions Phys. Rev. Lett. 100 113003
[24] Urban E, Johnson T A, Isenhower L, Yavuz D D, Walker T G and Saffman M 2009 Observation of Rydberg blockade between two atoms Nat. Phys. 5 100–4
[25] Anderlini M, Lee P J, Brown B L, Sebby-Strabley J, Phillips W D and Porto J V 2007 Controlled exchange interaction between pairs of neutral atoms in an optical lattice Nature 448 452
[26] Trotzky S, Cheinet P, Folling S, Feld M, Schnorrberger U, Rey A M, Polkovnikov A, Demler E A, Lukin M D and Bloch I 2008 Time-resolved observation and control of superexchange interactions with ultracold atoms in optical lattices Science 319 295
[27] Boyd M M, Zelevinsky T, Ludlow A D, Foreman S M, Blatt S, Ido T and Ye J 2006 Optical atomic coherence at the 1-second time scale Science 314 1430
[28] Miroshnychenko Y, Alt W, Dotsenko I, Förster L, Khudaverdyan M, Meschede D, Schrader D and Rauchschewitz A 2008 Quantum engineering: an atom-sorting machine Nature 442 151
[29] You L and Chapman M S 2000 Quantum entanglement using trapped atomic spins Phys. Rev. A 62 052302
[30] Cirac J I and Zoller P 2000 A scalable quantum computer with ions in an array of microtraps Nature 404 579
[31] Chin C, Vuletić V, Kerman A J and Chu S 2001 Measurement of collision shifts on clock transition and quantum computation in optical lattices Proc. XV Int. Conf. on Laser Spectroscopy
[32] Calarco T, Dorner U, Julienne P S, Williams C J and Zoller P 2004 Quantum computations with atoms in optical lattices: marker qubits and molecular interactions Phys. Rev. A 70 012306
[33] Shimizu F 2004 Scalable quantum computer with optical lattices Japan. J. Appl. Phys. 43 8376–82
[34] DiVincenzo D P 2000 The physical implementation of quantum computation Fortschr. Phys. 48 771
[35] Kohl M, Moritz H, Stoferle T, Gunter K and Esslinger T 2005 Fermionic atoms in a three dimensional optical lattice: observing fermi surfaces, dynamics and interactions Phys. Rev. Lett. 94 080403
[36] Kohler T, Goral K and Julienne P S 2006 Production of cold molecules via magnetically tunable feshbach resonances Rev. Mod. Phys. 78 1311
[37] Schiff L 1972 Quantum Mechanics (New York: McGraw-Hill)
[38] Deutsch I H and Jessen P S 1998 Quantum-state control in optical lattices Phys. Rev. A 57 1972–86