Nondestructive Cooling of an Atomic Quantum Register via State-Insensitive Rydberg Interactions

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We propose a protocol for sympathetically cooling neutral atoms without destroying the quantum information stored in their internal states. This is achieved by designing state-insensitive Rydberg interactions between the data-carrying atoms and cold auxiliary atoms. The resulting interactions give rise to an effective phonon coupling, which leads to the transfer of heat from the data atoms to the auxiliary atoms, where the latter can be cooled by conventional methods. This can be used to extend the lifetime of quantum storage based on neutral atoms and can have applications for long quantum computations. The protocol can also be modified to realize state-insensitive interactions between the data and the auxiliary atoms but tunable and nontrivial interactions among the data atoms, allowing one to simultaneously cool and simulate a quantum spin model.

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Introduction.—In recent years, neutral atoms stored in individual traps [1–7] have emerged as a powerful resource for quantum information and quantum technologies [8–13]. Considerable effort is currently being invested in developing neutral atom traps that are insensitive to the internal state of the atom [2,14–16]. These so-called magic traps attempt to achieve what is naturally available with trapped ions, since the trapping of the latter relies on the net charge of the ion, and hence is independent of its internal electronic state. The magic trapping of neutral atoms reduces heating and dephasing associated with the fact that different electronic states may have different trapping potentials. Nevertheless, even with such magic trapping conditions, heating of the motional degrees of freedom (d.o.f.) of the atoms can occur because of, for example, the shaking of the atomic array due to laser intensity noise [17], mechanical forces from Rydberg interactions [18–20], or incoherent light scattering [21].

Such heating of the atomic motion, when combined with state-dependent Rydberg mediated gates, generally leads to reduced fidelities and loss of coherence, which is particularly problematic for long quantum simulations or computations [22–24]. It is therefore desirable to develop schemes to cool the atomic motion without destroying the quantum information stored in the internal states. The conventional laser cooling techniques [25–27] are not suitable for this task since they involve optical pumping which, in general, destroys the quantum information.

Several approaches for this problem have already been proposed in the past, from immersing the atomic lattice in a superfluid [28] to using cavity-assisted cooling [29]. It has also been shown that alkaline-earth atoms can be laser-cooled without destroying the quantum information provided it is stored in the nuclear spin [30].

In this Letter, we introduce two schemes for achieving state-insensitive interactions between neutral atoms, another natural and useful tool of trapped ions. We further show how to use these interactions to realize a state-preserving cooling procedure, inspired by sympathetic cooling of trapped ions [31,32]. In contrast to the protocols in Refs. [28,29], ours requires only ingredients and capabilities that are already present in many neutral atom experiments: auxiliary atoms and Rydberg interactions.

The scenario we have in mind is the following: we assume one starts with a quantum data register composed of an array of N atoms, each in an individual trap, cooled to the vibrational ground state and optically pumped to a particular ground state. Each atom encodes a two-level system in its ground states. One then uses Rydberg interactions to perform a quantum computation or simulation, during which the atoms are heated. To cool the data register we introduce N auxiliary atoms, one for each data atom [see Fig. 1(a)], that have been precooled using any of the standard methods. The data and auxiliary atoms can then be coupled via Rydberg interactions, implementing a *phonon-swap* gate—a coherent exchange of vibrational quanta. A key requirement of

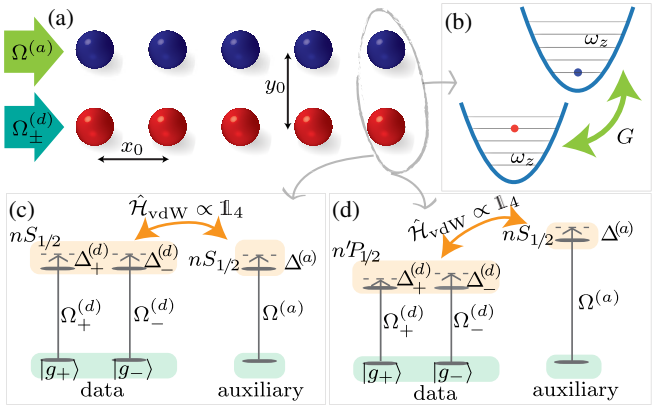

FIG. 1. Schematic of the phonon-swap protocol. (a) For each data atom (red, bottom) we place another auxiliary atom (blue, top) at an equal distance y_0. We assume a 1D chain of data and auxiliary atoms, with a lattice spacing of x_0. (b) The Rydberg interactions give rise to effective coupling G between the vibrational modes of the data and auxiliary atoms. (c),(d) Two schemes that lead to spin-insensitive interactions between the data and the auxiliary atoms are, in (c), the ground states of all atoms are weakly coupled to Rydberg $S_{1/2}$ states. In (d), the data atoms are coupled to $n'P_{1/2}$ states and the auxiliary to $nS_{1/2}$, where $|n-n'| \gg 1$.

this protocol is for the interactions between the auxiliary and data atoms to be insensitive to the internal state of the data atoms. Unlike the Coulomb interaction between ions which naturally satisfies this requirement, the Rydberg interactions between neutral atoms are inherently state dependent. As we show in this Letter, a careful choice of the Rydberg states can nevertheless satisfy this requirement.

Another requirement is that the phonon-swap interactions should not induce unwanted state-dependent couplings between the data atoms. In the first of our two schemes [see Figs. 1(c), 1(d)], the interactions between any pair of atoms (data-data and data-auxiliary) are independent of the internal state. This scheme consists of pausing the quantum computation or simulation, performing the phonon swap, and then resuming the computation or simulation. In the second scheme, the data and auxiliary atoms are addressed separately, which allows one to design state-insensitive data-auxiliary interactions but tunable data-data interactions. As an example, we show how this can be used to implement the phonon-swap while simultaneously performing a quantum simulation of a spin model on the data atoms. Finally, for both of the above schemes, one can laser cool the auxiliary atoms during the phonon swap. Because of the quantum Zeno effect [33], this has the additional advantage of preventing certain coherent heating mechanisms, such as those due to the Rydberg interactions, from taking place at all. We leave the detailed study of such a scheme for future work.

Phonon-swap for two atoms.—Let us first consider the case of two atoms: a two-level data atom "d" and a single-level auxiliary atom "a." The two atoms are each trapped in a three-dimensional harmonic potential separated by a distance $\mathbf{r}$. In recent experiments [4,25–27,34,35], the confinement along two directions (x, y) is often much stronger than along the third (z); i.e., the trap frequencies satisfy $\omega_x, \omega_y \gg \omega_z$. For simplicity, we focus on cooling the weakest direction (z). Cooling the two components perpendicular to the interatomic axis is a trivial generalization of this section. The third component [y direction in Fig. 1(a)] requires more care but can be cooled via an adiabatic protocol [36].

The Hamiltonian consisting of both the vibrational and the internal d.o.f. is ($\hbar = 1$) $\hat{\mathcal{H}} = \omega_z(\hat{d}^\dagger\hat{d} + \hat{a}^\dagger\hat{a}) + \hat{\mathcal{H}}_s + \hat{\mathcal{H}}_{\rm int}(\mathbf{r})$, where $\hat{d}(\hat{a})$ is the phonon annihilation operator of the data (auxiliary) atom along the z direction; $\hat{\mathcal{H}}_s$ acts on the internal (spin) d.o.f. of the data atom, and $\hat{\mathcal{H}}_{\rm int}(\mathbf{r})$ describes the interaction between the two atoms that, in principle, couples motion and spin. Since we want to preserve the spin state of the data atom, the phonon dynamics should be decoupled from the spin; i.e., we want $\hat{\mathcal{H}}_{\rm int} = \mathbb{1}_{\rm internal} \otimes V(r)$ to be an identity operator on the internal states. As we later show, by weakly laser dressing the ground states with Rydberg states, it is possible to obtain effective interactions of such form, where $V(r) = \mathcal{A}/(r^6 + R_c^6)$ for coupling $\mathcal{A}$ and blockade radius R_c.

For now, let us assume these interactions and Taylor expand them to second order in the small quantum fluctuations on top of the macroscopic separation r_0, which we assume to be along one of the strongly confined directions [Fig. 1(a)]. This gives rise to a quadratic Hamiltonian [36],

$$\hat{\mathcal{H}}_{\rm ph,2} = \omega_z(\hat{d}^\dagger\hat{d} + \hat{a}^\dagger\hat{a}) - \frac{G}{2}[(\hat{d} + \hat{d}^\dagger)^2 + (\hat{a} + \hat{a}^\dagger)^2] + G(\hat{d} + \hat{d}^\dagger)(\hat{a} + \hat{a}^\dagger), \quad (1)$$

where $G = (3\mathcal{A}/M\omega_z r_0^8)\{1/[1 + (R_c/r_0)^6]^2\}$ is the phonon coupling strength and M is the atomic mass. In the regime where $\omega_z \gg G$, only the number-conserving terms are relevant, giving a "beam splitter" interaction (in the rotating frame) $\hat{\mathcal{H}}_{\rm ph,2} = G(\hat{a}^\dagger\hat{d} + \hat{d}^\dagger\hat{a})$. This Hamiltonian effectuates a state-transfer between the two vibrational modes in a time of $t_s = \pi/2G$, swapping the phonons of the data atom with those of the auxiliary atom. This cools the data atom down to the initial phonon occupancy of the auxiliary atom.

Phonon-swap for 1D chain.—The discussed protocol can be easily generalized for an ensemble of atoms. We simply associate each data atom with a cold auxiliary atom. For concreteness, we consider a chain of data atoms with a lattice constant x_0, separated by a distance of y_0 from a chain of cold auxiliary atoms [Fig. 1(a)]. The many-body Hamiltonian is quadratic with approximate power-law decaying hopping between the sites [36]

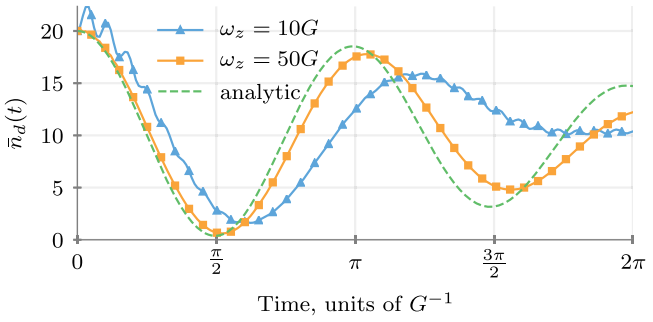

FIG. 2. The average number of phonons in the data atoms as a function of time computed numerically (solid lines) for different values of ω_z for two chains of 50 atoms, including the counter-rotating terms in Eq. (1), and using Eq. (3) (dashed line). Here, $\eta = 1$, $\bar{n}_d(0) = 20$, $\bar{n}_a(0) = 0$.

$$\hat{\mathcal{H}}_{\mathrm{ph},1D} = \sum_{i\neq j} \frac{G}{\eta^8 |i-j|^8} (\hat{a}_i^\dagger \hat{a}_j + \hat{d}_i^\dagger \hat{d}_j) + \sum_{ij} \frac{G}{[\eta^2 (i-j)^2 + 1]^4} (\hat{a}_i \hat{d}_j^\dagger + \hat{a}_i^\dagger \hat{d}_j), \quad (2)$$

where $\eta \equiv x_0/y_0$. Here we defined G in terms of the smallest distance between a data atom and its auxiliary, i.e., y_0 (see Fig. 1). Clearly, as $\eta \to \infty$, it is sufficient to consider only the nearest-neighbor data-auxiliary interactions. In such a case, we recover the situation in the previous section: each data-auxiliary pair perfectly swaps their phonons after a time of $t_s = \pi/2G$. If we also take into account next-nearest-neighbor data-auxiliary interactions, we find [36] that the average phonon occupancy of the data atoms is

$$\bar{n}_d(t) = \frac{\bar{n}_a(0) + \bar{n}_d(0)}{2} - \frac{\bar{n}_a(0) - \bar{n}_d(0)}{2} \times J_0\left(\frac{4Gt}{(1+\eta^2)^4}\right) \cos(2Gt), \quad (3)$$

where $\bar{n}_d(t)$ $[\bar{n}_a(t)]$ is the average occupancy of data (auxiliary) atoms at time t and $J_0(z)$ is a Bessel function of the first kind. Equation (3) is quantitatively accurate (see Fig. 2) at short timescales, when the effects of the long-range interactions are less important. As $\eta \to \infty$ we have $J_0 \to 1$ which reproduces the case of independent pairwise phonon swaps. Moreover, $t_s = \pi/2G$ is still the nearly optimal swap time (see Fig. 2) and even with $\eta = 1$ we can still achieve a high-efficiency swap. Assuming for simplicity that the auxiliary atoms are initially in the vibrational ground state, we obtain a swap efficiency of $1 - [\bar{n}_d(t_s)/\bar{n}_d(0)] = \frac{1}{2} + \frac{1}{2} J_0(\pi/8) \approx 98\%$. Furthermore, Eq. (3) remains qualitatively accurate even at longer timescales. As $t \to \infty$ we have $J_0 \to 0$ and we see that the mean phonon occupancy of all atoms is the average of the total initial number of phonons, as expected.

The above discussion concludes that to cool an atomic register consisting of many atoms in arbitrary geometries and dimensions, we simply perform the phonon swap as if all the data-auxiliary pairs are independent. The many-body interactions only lead to a small degradation in the swap efficiency.

State-insensitive Rydberg interactions.—We now turn to discuss how to obtain the spin-independent interactions by utilizing the van der Waals (vdW) couplings between Rydberg states. Specifically, we concentrate on alkali atoms and consider weakly laser admixing two hyperfine ground states (see Supplemental Material for an explicit example [36]) representing the spin-1/2, $|g_+\rangle, |g_-\rangle$, to Rydberg states $|r_+\rangle, |r_-\rangle$, depicting the magnetic sublevels of either $S_{1/2}$ or $P_{1/2}$ manifolds, as shown in Figs. 1(c), 1(d). The vdW couplings $\hat{\mathcal{H}}_{\mathrm{vdW}}$ between the Rydberg states then give rise to effective interactions between the dressed ground states. The relevant Hamiltonian is $\hat{\mathcal{H}} = \sum_{i=1,2} (\hat{\mathcal{H}}_A^{(i)} + \hat{\mathcal{H}}_L^{(i)}) + \hat{\mathcal{H}}_{\mathrm{vdW}}$ where $\hat{\mathcal{H}}_A^{(i)} = -\Delta_+^{(i)} |r_+^{(i)}\rangle\langle r_+^{(i)}| - \Delta_-^{(i)} |r_-^{(i)}\rangle\langle r_-^{(i)}|$ and $\hat{\mathcal{H}}_L^{(i)} = (\Omega_+^{(i)}/2)|g_+^{(i)}\rangle\langle r_+^{(i)}| + (\Omega_-^{(i)}/2)|g_-^{(i)}\rangle\langle r_-^{(i)}| + \mathrm{H.c.}$ are the atomic and laser Hamiltonians, respectively, in the rotating frame within the rotating wave approximation. Here, $\Omega_\pm^{(i)}$ are Rabi frequencies and $\Delta_\pm^{(i)} \gg \Omega_\pm^{(i)}$ the laser detunings. Note that for the auxiliary atoms, it is sufficient to consider a single ground state and hence a single laser. However, we must take into account all the states in the Rydberg manifold because, in general, $\hat{\mathcal{H}}_{\mathrm{vdW}}$ may contain both diagonal and off-diagonal matrix elements. This fact has been used previously to construct tunable spin-spin interactions [37,38]. A sufficient condition to obtain spin-independent interactions is for $\hat{\mathcal{H}}_{\mathrm{vdW}}$ to be proportional to an identity, together with a suitable choice of the laser parameters. We show below two simple schemes using $S_{1/2}$ and $P_{1/2}$ states that satisfy well this requirement.

The vdW Hamiltonian between two atoms in either $S_{1/2} + S_{1/2}$, $S_{1/2} + P_{1/2}$ or $P_{1/2} + P_{1/2}$ in the Zeeman basis has the following form [36]:

$$\hat{\mathcal{H}}_{\mathrm{vdW}} = \frac{C_6}{r^6} \mathbb{1}_4 - \frac{C_6^{(a)} + C_6^{(b)} - C_6^{(c)} - C_6^{(d)}}{r^6} \mathcal{D}_0(\theta, \phi), \quad (4)$$

$$C_6 \equiv \frac{2}{27} [C_6^{(a)} + 4C_6^{(b)} + 2(C_6^{(c)} + C_6^{(d)})], \quad (5)$$

where the $C_6^{(p)}$ coefficients correspond to the four different channels describing the possible (L, J) quantum numbers of the intermediate states and $\mathcal{D}_0(\theta, \phi)$ is a traceless matrix that depends on the angles between the interatomic and quantization axes. The channels for $S_{1/2} + S_{1/2}$ and $S_{1/2} + P_{1/2}$ are shown in Table I.

Phonon-swap with $S + S$ states.—The first scheme uses the fact that for a pair of atoms in $nS_{1/2}$ states, the second term in Eq. (4) approximately vanishes [37,39]. This can be

TABLE I. The four channels describing the dipole-allowed virtual processes $(L_1, J_1) + (L_2, J_2) \to (L_1', J_1') + (L_2', J_2')$ that lead to vdW interactions.

	$S_{1/2} + S_{1/2}$	$S_{1/2} + P_{1/2}$
(a)	$S_{1/2} + S_{1/2} \to P_{1/2} + P_{1/2}$	$S_{1/2} + P_{1/2} \to P_{1/2} + S_{1/2}$
(b)	$S_{1/2} + S_{1/2} \to P_{3/2} + P_{3/2}$	$S_{1/2} + P_{1/2} \to P_{3/2} + D_{3/2}$
(c)	$S_{1/2} + S_{1/2} \to P_{3/2} + P_{1/2}$	$S_{1/2} + P_{1/2} \to P_{3/2} + S_{1/2}$
(d)	$S_{1/2} + S_{1/2} \to P_{1/2} + P_{3/2}$	$S_{1/2} + P_{1/2} \to P_{1/2} + D_{3/2}$

seen from Table I, which shows that the difference between the four channels is only in the fine structure of the intermediate states. In the limit of vanishing fine structure, we have $C_6^{(a)} = C_6^{(b)} = C_6^{(c)} = C_6^{(d)}$. This can also be understood intuitively: the vdW interactions arise from second-order perturbation theory, where the two electrons undergo virtual transitions to intermediate states allowed by the selection rules. If we neglect the fine structure, we are free to use the uncoupled basis $(|L, m_L\rangle \otimes |S, m_S\rangle)$ for the intermediate levels. Since $S_{1/2}$ states are proportional to electronic spin states with definite m_S, i.e., $|S_{1/2}, m_J = \pm\frac{1}{2}\rangle = |L = 0, m_L = 0\rangle \otimes |S = \frac{1}{2}, m_S = \pm\frac{1}{2}\rangle$ and because the dipole-dipole interactions do not act on the electronic spin, the vdW couplings cannot mix states with different m_J. The correction to this scales as $\Delta_{\mathrm{FS}}/\delta$, where Δ_{FS} is the fine structure splitting and δ the energy difference to the intermediate states.

Neglecting these small corrections, and to fourth order in the small parameter $\epsilon = \Omega/2\Delta$, the effective spin-spin interactions between any two data atoms are given by

$$\hat{\mathcal{H}}_{\mathrm{int}}(r) = \mathrm{diag}(\tilde{V}_{++}\tilde{V}_{+-}\tilde{V}_{-+}\tilde{V}_{--}). \tag{6}$$

In the case of data-auxiliary interactions, we have a 2×2 version of Eq. (6). In both cases, the matrix elements are

$$\tilde{V}_{\mu\nu} = \left(\frac{\Omega_\mu^{(1)}\Omega_\nu^{(2)}}{4\Delta_\mu^{(1)}\Delta_\nu^{(2)}}\right)^2 \frac{C_6}{r^6 - \frac{C_6}{\Delta_\mu^{(1)} + \Delta_\nu^{(1)}}}, \tag{7}$$

which are spin independent (i.e., $\tilde{V}_{++} = \tilde{V}_{+-} = \tilde{V}_{-+} = \tilde{V}_{--}$) for a suitable choice of the laser parameters. A trivial example consists of the two laser fields being identical. The cooling protocol with this scheme would thus consist of stopping the quantum simulation or computation, weakly coupling the ground states of both the data and auxiliary atoms to $nS_{1/2}$ states, and waiting for a time of t_s. As an example, Rb atoms separated by 2.36 μm, and weakly coupled to $60S_{1/2}$ ($C_6/2\pi \approx 138.5$ GHz μm^6) with $\Omega/2\pi = 100$ MHz [40] and $\Delta/2\pi = 200$ MHz would experience a phonon coupling of $G/2\pi \approx 1.48$ kHz assuming a trap frequency of $\omega_z/2\pi = 15$ kHz. G is about an order of magnitude smaller than the trap frequency and about 2 orders of magnitude larger than the effective decay rate $\epsilon^2\Gamma_{60S}/2\pi \approx 0.043$ kHz, where Γ_{60S} is the decay rate of $60S_{1/2}$ states. The deviation of $\hat{\mathcal{H}}_{\mathrm{vdW}}$ from identity, which we define by the ratio of the operator norms of the two terms in Eq. (4), is, in this case, $\approx$0.027. This error can be reduced by driving the two atoms to different principal quantum numbers [Fig. 3(b)]. This generally reduces the C_6 coefficient [Fig. 3(a)], but it can nevertheless be sufficiently strong. For instance, $74S_{1/2} + 64S_{1/2}$ yields $C_6/2\pi \approx 29$ GHz μm^6 (a factor of 5 smaller than for $60S_{1/2} + 60S_{1/2}$) with an error of $\approx$0.003 (an order of magnitude smaller).

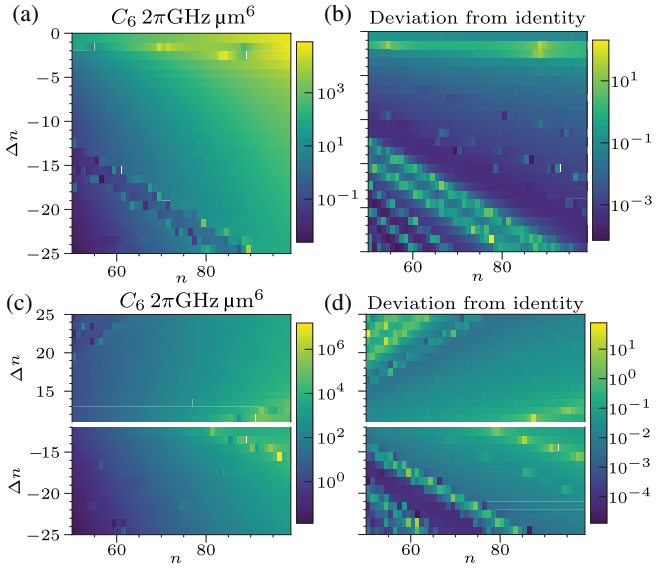

FIG. 3. (a),(c) The spin-insensitive interaction strength C_6 and (b),(d) deviation from identity for (top) $nS_{1/2} + n'S_{1/2}$ and (bottom) $nS_{1/2} + n'P_{1/2}$ as a function of n and $\Delta n = n' - n$ for Rb atoms. In the case of $nS_{1/2} + n'P_{1/2}$, we take $\min|\Delta n| = 10$.

Phonon-swap with $S + P$ states.—This brings us to the second scheme, in which the auxiliary atoms are coupled to $nS_{1/2}$ states, while the data atoms to $n'P_{1/2}$ states, where $|\Delta n| = |n' - n| \gg 1$ in order to ensure that the dipolar interactions between them can be ignored [41]. Such a configuration not only gives spin-independent data-auxiliary interactions, as we will explain below, but also gives rise to tunable spin-spin interactions between the data atoms [37]. To see why $S_{1/2} + P_{1/2}$ gives rise to $\hat{\mathcal{H}}_{\mathrm{vdW}} \propto \mathbb{1}$, note that channels ($a$, c) as well as (b, d) in Table I only differ by the fine structure in one of the terms. In the limit of vanishing fine structure, the four channels cancel each other pairwise, eliminating $\mathcal{D}_0(\theta, \phi)$ in Eq. (4).

Intuitively, the same argument as in the $S_{1/2} + S_{1/2}$ case shows that there cannot be any mixing between states involving different m_J of the $S_{1/2}$ atom. Hence, in the absence of fine structure in the intermediate manifold, the $S_{1/2}$ atom is effectively decoupled and $\hat{\mathcal{H}}_{\mathrm{vdW}}$ must at least

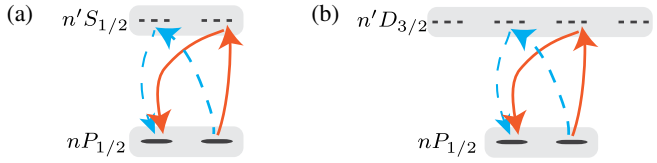

FIG. 4. The four virtual transitions that can couple the $m_J = 1/2$ magnetic state to $m_J = -1/2$, in the $P_{1/2}$ manifold.

be block diagonal. Within this approximation, we can understand why the remaining off-diagonal matrix elements also vanish by focusing solely on the $P_{1/2}$ atom. For each possible sub-channel of the $P_{1/2}$ atom [Fig. 4], there are exactly two processes that can couple its $m_J = +\frac{1}{2}$ and $m_J = -\frac{1}{2}$ states. These two processes, however, precisely destructively interfere. The resulting spin interactions between the data and auxiliary atoms have the same form as in Eqs. (6) and (7). The corresponding C_6, and the error due to the spin-dependent couplings, are shown in Figs. 3(c), 3(d), respectively.

The data atoms, on the other hand, experience nontrivial spin-spin interactions due to the $P_{1/2} + P_{1/2}$ vdW couplings. For the configuration in Fig. 1 (quantization axis parallel to interatomic axis), the $\mathcal{D}_0(\theta, \phi)$ matrix reads

$$\mathcal{D}_0(0, \phi) = \frac{2}{81}\begin{pmatrix} 1 & 0 & 0 & 0 \\ 0 & -1 & -4 & 0 \\ 0 & -4 & -1 & 0 \\ 0 & 0 & 0 & 1 \end{pmatrix}, \tag{8}$$

which gives rise to the following spin-1/2 Hamiltonian for the data atoms:

$$\hat{\mathcal{H}} = \sum_{ij} J_z^{ij} \hat{S}_z^{(i)} + J_{zz}^{ij} \hat{S}_z \hat{S}_z + (J_{+-}^{ij} \hat{S}_+^{(i)} \hat{S}_-^{(j)} + \text{H.c.}), \tag{9}$$

where $\hat{S}_\alpha^{(i)}$ are the spin-1/2 operators of atom i and $J_{\mu\nu}^{ij}$ are coefficients that depend on the geometry, laser parameters, and Rydberg interactions [37]. This approach can be extended to generate other spin-1/2 models, for instance, in two dimensions [37], with simultaneous cooling.

Summary and outlook.—We have presented a protocol for sympathetically cooling Rydberg atoms without destroying the quantum information stored in their internal states. This can have applications for future Rydberg-based quantum computers and simulators as well as other quantum technologies. Note that while we focused here on the weak coupling regime ($G \ll \omega_z$), which inevitably limits the phonon-swap time to $\sim 1/G \gg 1/\omega_z$, it is possible to speed it up by working in the strong coupling regime $G \sim \omega_z$ and employing optimal control techniques [42–44]. Furthermore, while we used vdW interactions, state-insensitive interactions can also be realized with dipole-dipole interactions and microwave dressing of Rydberg states [45]. Finally, our state-insensitive interaction schemes could potentially be used in other contexts, such as generating nonclassical states [46] and novel phases of matter [47] combining motional and electronic d.o.f.

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