Tunable Superfluidity and Quantum Magnetism with Ultracold Polar Molecules

Citation

Published Version
doi:10.1103/PhysRevLett.107.115301

Permanent link
http://nrs.harvard.edu/urn-3:HUL.InstRepos:8148862

Terms of Use
This article was downloaded from Harvard University’s DASH repository, and is made available under the terms and conditions applicable to Open Access Policy Articles, as set forth at http://nrs.harvard.edu/urn-3:HUL.InstRepos:dash.current.terms-of-use#OAP

Share Your Story
The Harvard community has made this article openly available. Please share how this access benefits you. Submit a story.

Accessibility
Tunable Superfluidity and Quantum Magnetism with Ultracold Polar Molecules

Alexey V. Gorshkov,1,* Salvatore R. Manmana,2,* Gang Chen,2
Jun Ye,2 Eugene Demler,3 Mikhail D. Lukin,3 and Ana Maria Rey2

1Institute for Quantum Information, California Institute of Technology, Pasadena, California 91125, USA
2JILA, University of Colorado and NIST, and Department of Physics, CU Boulder, Colorado 80309-0440, USA
3Physics Department, Harvard University, Cambridge, Massachusetts 02138, USA

(Dated: June 10, 2011)

By selecting two dressed rotational states of ultracold polar molecules in an optical lattice, we obtain a highly tunable generalization of the $t$-$J$ model, which we refer to as the $t$-$J$-$V$-$W$ model. In addition to XXZ spin exchange, the model features density-density interactions and novel density-spin interactions; all interactions are dipolar. We show that full control of all interaction parameters in both magnitude and sign can be achieved independently of each other and of the tunneling. As a first step towards demonstrating the potential of the system, we apply the density matrix renormalization group method (DMRG) to obtain the 1D phase diagram of the simplest experimentally realizable case. Specifically, we show that the tunability and the long-range nature of the interactions in the $t$-$J$-$V$-$W$ model enable enhanced superfluidity. Finally, we show that Bloch oscillations in a tilted lattice can be used to probe the phase diagram experimentally.

PACS numbers: 67.85.-d, 71.10.Fd, 33.80.-b, 71.10.Pm

Experiments with ultracold atoms have recently extended the range of candidate systems for realizing unconventional states of matter and enabled the simulation of models describing condensed matter phenomena [1]. One major goal of current research at this interface between condensed matter and atomic physics is to emulate models describing condensed matter phenomena [1]. The under-
An important difference of Eq. (1) from the Hamiltonian studied in Ref. [17], which is most closely related to our work, is the presence of the $J_z$ term [11].

Eq. (1) possesses several aspects that can lead to unconventional many-body phases. Most notably, all the interactions are long-range; in particular, repulsive long-range density-density interactions are expected to stabilize superfluid correlations [23]. Furthermore, the novel $W$ term can break SU(2) symmetry of the system even when $J_z = J_z$. It is also crucial that $J_z$ and $J_z$ can be tuned in both sign and magnitude (up to $\sim 100$ kHz in LiCs) independently from each other and from $t$. This contrasts with the cold atom realization of the $t$-$J$ model, where $J < t$. Finally, we note that, in the limit of unit filling ($n = 1$), only the terms with $J_z$ and $J_z$ survive, realizing an XXZ model with long-range dipolar interactions, and, for $J_z = J_z$, the Heisenberg model.

**Tuning $J_z$, $J_z$, $V$, and $W$.**—We now show how to use DC electric and microwave fields [12, 14–18, 22, 24–28] to control $J_z$, $J_z$, $V$, and $W$, leaving a detailed derivation of Eq. (1) to Ref. [29]. A polar molecule in a DC electric field is described by the rigid-rotor Hamiltonian $H_0 = BN^2 - d_0E$, with the rotational constant $B$, the angular momentum operator $N$, and the dipole moment operator $d$. We set $d = e_\phi \cdot d$, where $e_\phi = z$ and $e_\phi = \mp (x \pm iy)/\sqrt{2}$ [25]. At $E = 0$, $H_0$ has eigenstates $|N, M\rangle$ obeying $N^2|N, M\rangle = N(N + 1)|N, M\rangle$ and $N_z|N, M\rangle = M|N, M\rangle$. As $E$ is increased, eigenstates with the same $M$ mix, forming superpositions of $|N, M\rangle$ with different $N$. We denote the state at finite $E$ adiabatically connected to $|N, M\rangle$ as $|\phi_{N,M}\rangle$ and define, for notational convenience, $|N\rangle = |\phi_{N,0}\rangle$, $|N\rangle = |\phi_{N,1}\rangle$, and $|N\rangle = |\phi_{N,2}\rangle$ [see Fig. 1(a)]. For appropriate choices of levels, energy conservation ensures that dipole-dipole interactions between two molecules preserve their total $N_2$, resulting in the many-body interaction Hamiltonian

$$H_{dd} = \frac{1}{2} \sum_{i \neq j} |R_i - R_j|^{-3} (d_0^{(i)} d_0^{(j)} + d_+^{(i)} d_-^{(j)} + d_-^{(i)} d_+^{(j)}).$$

We now focus on the two configurations of levels and microwave fields shown in Figs. 1(b,c), which allow, respectively, for the realization of the simple case $J_z = V = W = 0$ and for the full tunability of $J_z$, $J_z$, $V$, and $W$.

In the first configuration, states $|m_0\rangle = \sqrt{\alpha}(0) + \sqrt{1 - \alpha}|\bar{T}\rangle$ and $|m_1\rangle = |1\rangle$ [Fig. 1(b)] are chosen, where the coefficient $\alpha$ in the dressed state $|m_0\rangle$ is controlled by the ratio between the Rabi frequency and the detuning of a $\sigma_+$-polarized microwave field acting on the $|0\rangle - |\bar{T}\rangle$ transition [26]. In $^{40}$K$^{87}$Rb, microwave intensity of a few W/cm$^2$ is sufficient to address all hyperfine levels equally. Projecting on states $|0\rangle$, $|1\rangle$, and $|\bar{T}\rangle$ and keeping only energy conserving terms, we obtain

$$d_0^{(1)} d_0^{(j)} = \mu_0^2 |0\rangle\langle 1| + \text{h.c.}. \quad d_+^{(1)} d_-^{(j)} + d_-^{(1)} d_+^{(j)} = -\mu_2^2 |0\rangle\langle 1| + \text{h.c.},$$

where $k \in \{i, j\}$, $s \in \{0, 1, \bar{T}\}$, $\mu_s = \langle s|d_0 |s\rangle$, and $\mu_{s'} = \langle s|d_0 |s'\rangle$ for the appropriate $p$. The minus sign in front of $\mu_2^2$, and $\mu_2^2$ is crucial to the tunability of $V$, $W$, and $J_z$, and $J_z$ appears because two dipoles rotating in the $x$-$y$ plane interact on average attractively [26]. Projecting on $|m_0\rangle$ and $|m_1\rangle$ and again keeping only energy conserving terms, we obtain

$$d_0^{(1)} d_0^{(j)} + \frac{1}{2} (d_+^{(1)} d_-^{(j)} + d_-^{(1)} d_+^{(j)}) = \sum_{p} B_p |m_p m_p| \langle m_p m_p |ij + \frac{1}{2} \sum_{p,q} B_p A_q |m_p m_q| \langle m_p m_q |ij + \frac{1}{2} \langle m_0 m_1 |m_1 m_0 |ij + \text{h.c.},$$

where $p, q \in \{0, 1\}$ and $A_0 = \mu_0^2 (1 - a)\mu_T$, $A_1 = \mu_1$, $B_0 = -\mu_2^2 a (1 - a)$, $B_1 = 0$, $J_z = 2\mu_2^2 a - \mu_2^2 (1 - a)$. $A_0$ can be thought of as the effective dipole moment of $|m_0\rangle$, while $B_0$ is the contribution to $H_{dd}$ from transition dipole moments between rotor states making up $|m_p\rangle$. From the comparison of this equation and Eq. (1), we can read off $V = (A_0 + A_1)^2 + B_0 + B_1)/4$, $W = (A_0^2 + B_0 - A_1^2 - B_1)/2$, $J_z = (A_0 - A_1)^2 + B_0 + B_1$. A simple case $J_z = W = 0$ and $J_z > 0$ studied below can be approximately implemented using $a = 1$, which does not require a microwave field. A small $E$ field ($dE/B > 0.1$ in $^{40}$K$^{87}$Rb) is needed to prevent dipole-dipole and hyperfine interactions from populating $|\phi_{1,\mp}\rangle$. At $dE/B = 0.1$, $V$, $W$, and $J_z$ are two orders of magnitude smaller than $J_z$, and can be neglected.

The second configuration we consider is $\{m_0\rangle, |m_1\rangle = \{\sqrt{\alpha}(2) + \sqrt{1 - \alpha}|\bar{T}\rangle, \sqrt{\alpha}(1) + \sqrt{1 - \alpha}|\bar{T}\rangle, \sqrt{1 - \alpha}|\bar{T}\rangle\}$ [Fig. 1(c)]. The three microwave fields shown in the figure, which allow to control the coefficients $a$, $b$, and $c$, together with the DC electric field constitute four knobs that allow for the full control over the four coefficients $V$, $W$, $J_z$, and $J_z$ [30]. We find that in a small sphere in the 4-dimensional $(dE/B, a, b, c)$ space around the point $(dE/B, a, b, c) = (2.97, 0.059, 0.56, 0.38)$, where $J_z = J_z = V = W = 0$, one can achieve any value of $V$, $W$, $J_z$, and $J_z$ up to an overall positive prefactor. Similarly, in the special case where a single microwave couples $|\bar{T}\rangle$ and $|1\rangle$, we have a two-dimensional $(dE/B, b)$ subspace (with $a = 0$, $c = 1 - b$), in which $J_z$ and $J_z$ can be fully controlled -- for simulations of the XXZ model.
unit filling — around \((dE/B, b) = (4.36, 0.56)\), where they both vanish. While these examples prove full controllability in their respective cases, for any desired relationship between \(V, W, J_z\), and \(J_{\perp}\), there is likely a different level configuration that gives stronger interactions and uses weaker \(E\), lower microwave intensity, and/or more convenient microwave frequencies.

**Phase diagrams of the nearest-neighbor and dipolar \(t-J_{\perp}\) chains.**—The full tunability of the \(t-J-V-W\) model provides access to a great variety of models with potentially exotic physics. As the simplest example of this physics, we present in Fig. 2 the 1D phase diagram in the limit \(V = W = J_z = 0\), which is one of the simplest experimentally achievable cases (see above). Before analyzing dipolar interactions, we present in Fig. 2(a) the phase diagram of the nearest-neighbor \(t-J_{\perp}\) chain as obtained using DMRG and following the analysis of Ref. [9]. The diagram is qualitatively similar to that of the standard \(t-J\) chain [9]: At fillings \(n < 0.65\), we identify a repulsive Luttinger liquid (Luttinger parameter \(K_\rho < 1\)) with dominant spin-density-wave (SDW) correlations, an attractive Luttinger liquid (\(K_\rho > 1\)) with dominant singlet and triplet superfluid correlations, a singlet superfluid with a spin gap, and phase separation. At larger fillings, the spin gap is always zero, but the other phases remain.

In Fig. 2(b), we present the phase diagram of this system in the presence of dipolar interactions. At low fillings, SDW, the gapped singlet superfluid, and phase separation are obtained. This suggests that experiments with ultracold molecules can be a useful tool for exploring the phase diagram of the standard \(t-J\) model for arbitrary values of \(J\) in contrast to \(J \ll t\) in ultracold atom realizations. At the same time, crucially, both diagrams in Fig. 2 feature a significant enhancement of the superfluid region compared to the original \(t-J\) model because the absence of attractive density-density interactions suppresses phase separation. Furthermore, the maximum value of the spin gap in the dipolar \(t-J_{\perp}\) chain is twice that in the nearest-neighbor \(t-J_{\perp}\) chain, which is, in turn, twice that in the original \(t-J\) model. The larger spin gap should facilitate its experimental observation. Note that, as a conservative estimate of the numerical errors in computing \(K_\rho\) in the presence of long-range interactions with up to 100 sites, we estimate the true line \(K_\rho = 1\) to lie in the shaded region \(0.85 \leq K_\rho \leq 1.15\). Thus, the gapless superfluid cannot be identified in Fig. 2(b) within our numerical precision. Furthermore, while the line \(K_\rho = 1\) and the line where the spin gap closes may coincide, our analysis of the correlation functions cannot rule out the existence of an exotic intermediate phase with a spin gap, \(K_\rho < 1\), and dominant superfluid correlations.

**Preparation and detection.**—Ground-states at specific points in the phase diagram can be prepared, e.g., by applying an additional microwave field coupling \(|m_0\rangle\) and \(|m_1\rangle\) and performing an adiabatic passage from an easily accessible state to the desired ground state by tuning the Rabi frequency and the detuning of the microwave field [18]. Direct probing of molecules [32] and conversion of molecules back to atoms [5] can in principle both be used for detection via noise-correlations in the time-of-flight [1] or via in-situ single-site imaging [33, 34].

As a specific example of a detection technique avail-
able in current experiments, we propose to use Bloch oscillations [35, 36]. In Fig. 3, we present our results as obtained via the Krylov-space variant of the adaptive t-DMRG [37, 38] for 20 sites when adding a linear field along the chain at filling \( n = 0.1 \) and \( n = 0.2 \) for different values of \( J_\perp/t \). In the singlet superfluid, the amplitude of the oscillations drops and the frequency of the oscillations doubles (see \( J_\perp/t = 7 \) in the inset) relative to the gapless phase due to the presence of bound pairs. Bloch oscillations should be observable in direct absorption imaging [32]. While neither the frequency nor the amplitude of the oscillations show any sharp features at this small system size, the fit of experimental data to numerical results should allow for the location of the phase transition (vertical line in Fig. 3) even for small system sizes. As a complementary method for identifying the transition, we propose spectroscopic measurement of the spin gap [39].

Outlook.—We have presented a toolbox for simulating a highly tunable anisotropic \( t-J-V-W \) model with polar molecules. The advantages of this molecular toolbox over its atomic counterpart are higher energy scales and independent tunability of interactions and tunneling. This toolbox should enable the simulation of condensed matter phenomena, as well as the stabilization and controllable preparation of unconventional phases, such as \( d \)-wave superfluids [3]. The phase diagram of the experimentally simple case of a \( t-J_\perp \) chain shows an enhanced superfluid region, which we propose to probe via Bloch oscillations. We expect that the 2D \( t-J-V-W \) model can similarly be tuned into exhibiting enhanced superfluidity [40].

The present Letter also opens other exciting research avenues. In particular, natural extensions of the model include [29]: spatially anisotropic interactions produced by a tilt in the DC electric field, spin-dependent tunneling obtained by adjusting lattice beams, \( S > 1/2 \) models realized by choosing more than two dressed states, and systems with an orbital degree of freedom encoded in the nuclear spin. Moreover, by considering molecular Wigner crystals [24, 41], where the intermolecular distances are smaller than in an optical lattice, one can further increase the interaction strength. Furthermore, by analogy with Ref. [18], we expect our ideas to be extendable to Rydberg atoms. Finally, one can envision applications of the present system to quantum computation (especially if one uses nuclear spin to store information), precision measurements, and controlled quantum chemistry [4].

We thank A. Muramatsu, M. Troyer, A. Moreno, D. Jin, G. Refael, J. Aldeguende, P. Julieanne, M. Babadi, I. Bloch, P. Rabl, G. Quéméner, A. Potter, B. Wunsch, P. Zoller, S. Pielawa, E. Berg, A. Daley, F. Mila, R. Noack, and H. Weimer for discussions. This work was supported by the NSF, NIST, the Lee A. DuBridge Foundation, the ARO with funding from DARPA-OLE, CUA, and AFOSR MURI. We acknowledge CPU time at ARSC.

* Authors contributed equally.

[29] A. V. Gorshkov et al., same-day arxiv submission.
[30] In this configuration, \( I_\perp = 2\mu_2^2 + (1 - a) - \mu_2^2 (1 - a)b - \mu_2^2 (1 - a)(1 - b - c) - \mu_3^2 a c, A_0 = \mu_2 a + \mu_2 (1 - a), A_1 = \mu_2 b + \mu_2 c + \mu_2 (1 - b - c), B_0 = -\mu_3^2 a (1 - a), \) and \( B_1 = 2\mu_2^2 b (1 - b - c) - \mu_3^2 b c - \mu_3^2 c (1 - b - c) \).
