Introduction

Future applications of quantum physics for quantum simulations, computation, communication, and metrology will require an extremely high degree of control of preparation, manipulation, and, last but not least, detection of strongly correlated states of quantum many body systems. Ultracold atoms offer an unprecedented playground for realization of these goals. Several paradigm examples of strongly correlated states have been successfully realized, such as the Mott insulator, the Tonks gas, and the Bose glass (for a review cf. [1]). A standard way of analyzing such systems is by releasing the atoms from the trap and performing destructive absorption spectroscopy, which only allows to measure the column density of the expanded cloud. Considerable attention has thus been devoted recently to novel methods of detection, that allow for measuring (spin) density-density and other higher order correlation functions. One of those methods is atomic noise interferometry [2], whose power is well illustrated in the recent observation of the bosonic and the fermionic Hanbury Brown-Twiss effect [3, 4]. Direct atom counting is another way to measure this effect, and to even go beyond it [5]; it can be realized directly with metastable Helium atoms [6, 7], or by using methods of cavity quantum electrodynamics (QED) [8].

Cavity QED is also essential in the recent proposals of Ref. [9, 10], while Ref. [11] proposes how to prepare and detect magnetic quantum phases using superlattices. All of the above approaches are, at least in some respects, destructive and frequently suffer from undesired atom number fluctuations inevitable in the preparation of the quantum states.

Here we propose a unique method that allows for spatially resolved quantum non-demolition (QND) [12] detection of quantum states of ultracold atoms with internal (pseudo-)spin degrees of freedom. Our approach is based on the idea of quantum noise limited polarization spectroscopy, demonstrated in Ref. [13]. Polarization spectroscopy itself consists in shining a polarized probe laser beam through the atomic spin system, and has been a subject of studies for years [14]. What is novel in quantum noise limited schemes is that quantum polarization memory of light is achieved, this approach has proved very successful for QND-based quantum interfaces between light and atoms. In particular, atomic squeezing [15], atomic entanglement, quantum memory, and teleportation have been achieved in this context (for a review see [16] and references therein). We have recently proposed to apply this approach to detect magnetic order of weakly correlated ultracold atoms [17]. Unfortunately, such proposal does not provide spatial resolution, and thus it cannot, e.g., discriminate between different antiferromagnetic quantum phases present in lattice models. Here we show that this important limitation can be overcome by using a standing wave probe laser configuration. Such a modification is crucial: it allows for spatially periodic QND coupling of light to atomic spins, and thus can reveal spin correlations with the period of coupling. The quantum noise of the transmitted light then carries information on the Fourier components of the spin density. Controlling the parameters of the optical lattice and of the probe light standing wave permits discrimination and characterization of various standard and "exotic" magnetic orderings.

In Fig. 1 we schematically depict various important lattice states, all of them having zero average of any component of the total spin. The paramagnetic state, Fig. 1(a), exhibits large (proportional to the total number of atoms, $N_{at}$) fluctuations of the total spin, whereas Fig. 1(b-d) show global singlets, i.e. states with total spin zero, which do not show any fluctuations. Obviously,
those most interesting strongly correlated global singlet states of Fig. 1(b-d) cannot be distinguished from each other by spatially homogeneous probing. Probing, however, only every second atom in, e.g., the dimerized state of Fig. 1(b), i.e., measuring the total staggered magnetization of the state, will result in its correlations imprinted on the light. As we shall demonstrate below, such type of measurement offers the possibility to distinguish between various global singlets, with minimal disturbance to the system.

Detection scheme The detection mechanism proposed here uses the off-resonant interaction of spin-$F$ atoms (i.e., atoms having a $2F + 1$ dimensional ground state manifold) with a polarized light beam propagating in $z$-direction (for a theoretical description of light-atom coupling cf. [16] [18]). The light pulse is characterized by the Stokes operators $\hat{s}_1$, $\hat{s}_2$, $\hat{s}_3$ corresponding to the difference of the number of photons in the $x$ and $y$ linear polarizations, in the $\pm 45^\circ$ linear polarizations, and in the two circular polarizations, respectively. Let us first consider a one-dimensional (1D) sample of $N_{\text{at}}$ atoms trapped in a 1D optical lattice of period $\pi/k$, oriented in $z$ direction, with the $n$-th atom located at position $z_n = (n - 1)\pi/k$, having the spin operator $\hat{J}(z_n)$. The relevant dispersive part of the interaction Hamiltonian between light and atoms reads [16] [18] (see methods)

$$\hat{H} = -\kappa \hat{s}_3 \hat{J}_{\text{eff}},$$

(1)

Here $\kappa$ is the coupling constant and $\hat{J}_{\text{eff}}$ is the $z$-component of the effective collective atomic spin $\hat{J} = \sum_n c_n \hat{J}(z_n)$. The coefficients $c_n$ account for the modification of the atom-light coupling due to a spatial modulation of the probe beam intensity (in a running wave configuration $c_1 \equiv 1$) and are the key parameters allowing for spatial resolution. As a particular case, we consider a standing wave configuration, as it is described in the caption of Fig. 2. Then,

$$c_n \equiv c_n(k_P,a) = 2 \int dz \cos^2(k_P(z-a))|w(z-z_n)|^2,$$

(2)

with $k_P$ as the wavevector of the standing wave probe, and $a$ its shift with respect to the optical lattice; $w(z-z_n)$ is the usual Wannier-type wave function of a single atom confined at $z_n$ in a deep optical lattice.

We take the probe beam to be strongly polarized in the $x$ direction so that $\hat{S}_1 (\hat{S}_1 = \int \hat{s}_1 dt) \text{ fulfills } \langle \hat{S}_1 \rangle = N_{\text{ph}}/2 \gg 1$. This fact permits to introduce canonical quadrature operators $\hat{X} = \hat{S}_2/\sqrt{N_{\text{ph}}}$, $\hat{P} = \hat{S}_3/\sqrt{N_{\text{ph}}}$, which satisfy $[\hat{X}, \hat{P}] \approx i$. Integrating the Heisenberg equation of motion for $\hat{s}_2$ shows that the effective collective spin $\hat{J} = \hat{J}_{\text{eff}}$ is imprinted on the $\hat{X}$ quadrature of light:

$$\hat{X}_{\text{out}} = \hat{X}_{\text{in}} - \frac{\kappa}{\sqrt{F_{N_{\text{at}}}}} \hat{J}_{\text{eff}}.$$

(3)

Since $\langle \hat{X}_{\text{in}} \rangle = 0$, the mean of this quadrature after passing through the sample is directly proportional to the mean of the $z$-component of the effective atomic spin. Its variance contains the shot noise of the incoming probe pulse and the variance of $\hat{J}_{\text{eff}}$. Assuming a coherent input beam:

$$\langle (\Delta \hat{X}_{\text{out}})^2 \rangle = \frac{1}{2} + \frac{\kappa^2}{2F_{N_{\text{at}}}} \langle (\hat{J}_{\text{eff}} - \langle \hat{J}_{\text{eff}} \rangle)^2 \rangle.$$  

(4)

The variance of the collective atomic spin can be efficiently determined from the measurement on light if $\kappa = \sqrt{d} \geq 1$ [16]. Here $d = N_{\text{at}}/A$ denotes the resonant column optical depth of the atomic sample and $\eta = (N_{\text{ph}} \sigma)^2/(\Delta^2 A)$ is the probability of resonant excitation per atom by the probe, where $\sigma$ is the cross section on resonance for the probe transition, $\Gamma$ the spontaneous decay rate, $\Delta$ the detuning from resonance and $A$ the cross section of the atomic ensemble illuminated by the probe. Since spontaneous emission destroys the spin state and the QND character of the coupling, the condition $\eta \ll 1$ has to be fulfilled. To evaluate the decoherence (and heating) caused by the probe, we need to estimate the effect of spontaneous emission on the variance of the light quadrature. The induced probe decoherence is given by $\langle (\Delta \hat{J}_{\text{coh}})^2 \rangle / F_{N_{\text{at}}} \approx \eta$ (in units of the coherent spin state noise). Therefore, the total measured noise reads

$$\langle (\Delta \hat{X}_{\text{out}})^2 \rangle = \frac{1}{2} + \kappa^2 \left[ \frac{\langle (\Delta \hat{J}_{\text{coh}})^2 \rangle}{F_{N_{\text{at}}} \eta} + \eta \right].$$

(5)

In order to measure the atomic spin fluctuations with the best possible accuracy, we need to maximize the signal to noise ratio $[\kappa^2 (\langle \Delta \hat{J}_{\text{coh}})^2 \rangle) / F_{N_{\text{at}}}]/[1/2 + \kappa^2 \eta]$. For a coherent input beam, this contribution is maximized for $\eta_{\text{opt}} \approx 1/\sqrt{2d}$. Notice that $\eta$ can be adjusted to $\eta_{\text{opt}}$, by choosing the appropriate detuning, intensity and duration of the laser probe. On the other hand, the optical depth of a sample of ultracold atoms, including e.g. a cubic lattice with $100 \times 100 \times 100$ atoms can easily reach values of few hundreds. Thus, the value of $\eta$ can be made much smaller than one [16]. Since the spin noise of interest (as shown later in Fig. 3) is of the order of unity, the additional noise provided by spontaneous emission does not, therefore, significantly modify the measured variance of the light quadrature.

Furthermore, a squeezing of the $\hat{X}$ quadrature of the incoming probe before passing it through the sample such that $\langle (\Delta \hat{X}_{\text{in}})^2 \rangle < 1/2$, allows to improve the signal-to-noise ratio even further. Finally, let us mention that from the data recorded at the homodyne detector also higher order terms $\langle (\Delta \hat{X}_{\text{out}})^m \rangle$ can be extracted.

In order to achieve a spatial modulation of the atom light coupling, $k_p/k \neq 1$ is necessary. The wavelength of the probe laser and thus $k_p$ is constrained by the near resonant condition for an efficient QND probing. Also the choice of the wavelength of the light forming the optical lattice is limited. Thus, tuning the wavelengths of the lasers only gives a restricted control over the ratio $k_p/k$. An appropriate choice of the ratio can, however,
be achieved in various ways: (i) trivially by modifying $k$ by changing the angle between the beams forming the optical lattice; (ii) more surprisingly, by shining the probe light at an angle $\theta$ to the lattice. The effective Hamiltonian then reads

$$H = -\kappa \hat{s}_3 \left( \cos \theta \hat{J}_z^{\text{eff}} - \sin \theta \hat{J}_y^{\text{eff}} \right),$$

and the effective wavevector for a 1D sample oriented in $z$-direction is $k_P = 2\pi \cos \theta / \lambda_P$, with $\lambda_P$ the wavelength of the probe laser. Finally, (iii) a probe standing wave with variable $k_P$ can be obtained from crossing the probe laser beams at a variable angle. This last approach leads, however, to a different interacting Hamiltonian.

**Quantum antiferromagnets in 1D** Let us now illustrate the power of our proposal, by applying it to concrete examples of strongly correlated states that can be realized with ultracold atoms. Particularly challenging are in this context various possible quantum antiferromagnetic states that lie in the center of interest of condensed matter and even high energy physics [21, 22]. We will analyze some of the states that appear as ground states (or idealizations thereof) of the generalized Heisenberg spin-1/2 atomic chain [23-25] (see Methods). For some choice of parameters such systems are in the, so called, Haldane phase. This phase is gapped, and the ground state is well described by a matrix product state (MPS), that is both translationally and rotationally invariant. A prominent example of such states is the AKLT state [20] (see Fig. [1](d)). In another regime of parameters the ground state exhibits the dimer, or Peierls, order and it represents an example of the Valence Bond Solid (VBS) state. To a good approximation the ground state with an even number of sites is formed by placing a dimer on every odd bond (i.e., pairing the atoms on these bonds into singlets, see Fig. [1](b)). Such a state breaks the translational invariance, but not the rotational symmetry. Yet for another choice of parameters, the system is in the critical phase, but exhibits some trimerization. The ideal trimer state corresponds to concatenations of triples of neighboring atoms forming a singlet, i.e. the state of zero total spin of the triple (Fig. [1](c)). The result of applying our detection method to the paramagnetic and these three types of *global singlet* states are shown in Fig. [3].

While the presence of fluctuations at $k_P/k = 0$ signals unambiguously the unpolarized paramagnetic state, the three distinct global singlets can be perfectly distinguished either by combining results for variable $k_P/k$ and $a$, or by fixing $k_P/k$ and varying solely $a$. For instance, at $k_P/k = 1/2$ and $a = 0$, the added noise (in units of the shot noise of the probe light) for the dimmerized, trimerized, and AKLT states is respectively $4\kappa^2/3$, $8\kappa^2/9$, and $2\kappa^2$. The above quoted values correspond the ideal case in which the atoms are $\delta$-localized on the minima of the optical lattice, while Fig. [3] shows the results for extended atomic Wannier functions. This fact yields an overall decrease of the noise contribution towards larger $k_P/k$ ratios. Except for this decrease, the pattern is repeated for larger values of $k_P/k$. The fluctuations on the light quadrature using the atomic ground states of the spin-1 lattice systems obtained numerically [26], present a very close resemblance to the ones depicted in Fig. [3].

**Quantum antiferromagnets in 2D** Our scheme is not limited to arrays of 1D systems. It can also detect and distinguish different antiferromagnetic ground states in lattice models in higher dimensions [21, 22]. As an illustration, we consider a Heisenberg spin-1/2 model on a square lattice with antiferromagnetic coupling. In this case it may happen that translational symmetry is broken and the ground state becomes the dimer VBS state (see Fig. [4](a)), which exhibits long range dimer order. Also exotic spin liquid states which do not break any symmetry and do not possess long range order can appear in an extended Heisenberg model. These states can be visualized as coherent superpositions of random dimer coverings, termed as Resonating Valence Bond (RVB) states (see Fig. [4](b)). We consider a 2D lattice in the $y-z$ plane, and the probe light strongly polarized along the $z$-direction propagating in $z$ direction. Assuming all atoms paired into dimers, we obtain that the additional noise contribution to the outgoing $x$ quadrature is proportional to the number of dimers oriented in the $z$-direction. A valence bond solid with all dimers oriented in the z-direction and a single valence bond state with randomly oriented dimers can then be easily distinguished, since for $k_P/k = 1/2$ and $a = 0$ the additional noise is $2\kappa^2$ and $\kappa^2$, respectively. The antiferromagnetic Néel state appearing in the 2D Heisenberg model can be indeed detected with the method of Ref. [17].

**Fermionic superfluidity** Another application of our scheme is to detect superfluid phases in fermionic ultracold gases in a QND way. For the probing standing wave set in the $z$-direction the transmitted light carries information on the fluctuations of $\hat{J}_z$ in a 2D or 3D system one can measure the fluctuations in $\hat{J}_y$ and $\hat{J}_x$ in the $y$ and $x$ direction, respectively, by changing the standing wave orientation. Similar quantities have been proposed and proved to detect superfluidity in ultracold Fermi gases [27-31]. Our scheme as described here can be specially interesting for a 1D system as one can change the probing momentum $k_P$ with respect to the Fermi momentum by tilting the probing standing wave with respect to the system long axis as discussed above. Comparison of these fluctuations for a paired and a non paired system reveals the existence of a pairing gap. Already in a homogeneous set-up one could measure the fluctuations in $\hat{J}_x$ that, in a similar way as the structure factor, change dramatically across the BCS to BEC crossover [32].

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### Methods

#### Atom-light Hamiltonian

The Hamiltonian of equation (1) is obtained from the off-resonant coupling to an atomic dipole transition. The excited atomic states can be adiabatically eliminated from the dipole-interaction Hamiltonian under the conditions described in the main text. The dispersive effects arising from the Stark shift of the atomic levels are described via an effective Hamiltonian

\[
\hat{H} = -\int_0^L dz \rho A \left( a_0 \hat{\phi} + a_1 \hat{s}_z \hat{j}_z + a_2 \left( \hat{j}_z^2 - \hat{s}_- \hat{j}_+ - \hat{s}_+ \hat{j}_- \right) \right)
\]

(7)

Here \( L \) is the length of the atomic sample and \( A \) its cross section overlapping with the probe light propagating in the \( z \)-direction, and \( \rho \equiv \rho(z) \) is the atomic density. The light is described via Stokes operators \( \hat{s}_i \equiv \hat{s}_i(z,t) \) \((\hat{s}_\pm = \hat{s}_1 \pm i \hat{s}_2)\) and \( \hat{\phi} \) is the total photon density. The coefficients \( a_i \) depend on the laser wavelength and the characteristics of the atomic transition. The first term, proportional to \( a_0 \), gives the (polarization independent) ac Stark shift. In the limit where the detuning is large compared to the hyperfine splitting of the excited state \( a_2 \approx 0 \), and only the linear QND coupling between the Stokes operator and atomic spin remains. This is equivalently written in equation (1) as a coupling between the Stokes operator and the effective atomic spin component \( \hat{j}_z^{\text{eff}} \). To proceed, we assume that \( \hat{j}_z^{\text{eff}} \) is time independent. This is valid since: (i) it is conserved by the Hamiltonian equation (1), and (ii) the measurement time for a pulsed probe is much shorter than the atomic spin diffusion time. The outgoing \( \hat{X} \) quadrature of light is then given by Eq. (3) provided that \( a_1^2 F^2 N_{\text{at}}^2 / 2 \ll 1 \).

Due to the interplay between the bilinear and the bi-quadratic interaction (parametrized by \( \beta \), this Hamiltonian presents three antiferromagnetic quantum phases: dimerized \( \beta \in (-3\pi/4, -\pi/4) \), Haldane \( \beta \in (-\pi/4, \pi/4) \), and critical \( \beta \in [\pi/4, \pi/2] \). The representative point of the dimerized phase is found at \( \beta = -\pi/2 \) where the two-site ground state is a singlet (dimer. \( \beta \in [\pi/4, \pi/2] \)): 

\[
|s\rangle_{1,2} = (|+1\rangle_1 |−1\rangle_2 + |−1\rangle_1 |+1\rangle_2 - |0\rangle_1 |0\rangle_2)/\sqrt{3}
\]

Here we denote by \( |m_z = \pm 1, 0\rangle_n \) the eigenstates of \( \hat{j}_z(z_n) \). This ground state is analytical only in the thermodynamic limit or for periodic boundary conditions. In a finite chain with an even number of sites, however, the ground state is close to a concatenation of spin-1 singlets (see Fig. 1(b)). The representative state of the Haldane phase is found at \( \beta = \arctan(1/3) \), where the exact ground state corresponds to the AKLT state (see Fig. 1(d)). The critical phase is less well understood, though known to have three correlations. At \( \beta = \pi/4 \) the three-site ground state is a singlet (trimer), and a caricature of the ground state for a larger number of sites is given by a concatenation of trimers (cf. Fig. 1(c)). These three isotropic ground states are eigenstates of any component of the total spin. Thus for \( k_F = k \) where all atoms couple equally to the light, no additional noise is imprinted on the \( \hat{X} \) quadrature. For \( k_F \neq k \), in general these states are not eigenstates of \( \hat{j}_z^{\text{eff}} \) (although \( \langle \hat{j}_z^{\text{eff}} \rangle = 0 \)) and as we demonstrate in Fig. 1(c), the fluctuations permit to discriminate them unambiguously.

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FIG. 1: Antiferromagnetic states of spin-1 lattice systems. (a) 3D cubic lattice with a paramagnetic state of unpolarized atoms; (b) dimerized state with pairs of neighbouring atoms forming singlets; (c) trimerized state with triples of neighbouring atoms forming singlets; (d) AKLT (Affleck-Kennedy-Lieb-Tasaki) state obtained from a concatenation of spin-1/2 singlets by projecting pairs of spins from different bonds into the subspace of total spin-1 [20].

FIG. 2: Schematic experimental setup. A strong laser beam initially polarized in $x$-direction is impinging on a 99/1 beamsplitter. The transmitted attenuated part of this probe is propagating through the sample and reflected off a mirror, such that a standing wave with wavevector $k_P$ in $z$ direction is formed at the position of the atomic sample. After the second pass, the laser beam is outcoupled to a homodyne detector, where $\langle \hat{S}_z \rangle$ is recorded. The atoms are trapped in an optical lattice with wavevector $k$. Its relative displacement with respect to the probe standing wave can be changed by either moving the end mirror, or by introducing a phase shift between the counterpropagating laser beams.
FIG. 3: Detection of antiferromagnetic states of spin-1 lattice systems. Fluctuations $\epsilon \equiv \langle (\Delta \hat{X}_{\text{out}})^2 \rangle - 1/2$ imprinted on the $X$ quadrature of the probe light beam transmitted through a sample of spin-1 atoms in a one-dimensional lattice for (a) unpolarized paramagnetic, (b) dimerized, (c) trimerized, and (d) AKLT states. While the presence of fluctuations at $k_P/k = 0$ signals the unpolarized paramagnetic state, the other three global singlets phases can be unambiguously distinguished at different values of the ratio $k_P/k$ and/or the shift $a$ (see text for details).
FIG. 4: Two-dimensional states of dimers. (a) Valence Bond Solid arranged in a regular crystal, (b) Valence Bond Solid consisting in a random covering of dimers.