Letter

Two-atom interferometer

M. F. Locke [1,2] and Y. H. Shih [2]

[1] *Advanced Quantum Optic Sensors Laboratory, Naval Air Warfare Center, Patuxent River, Maryland 20670, USA*

[2] *Department of Physics, University of Maryland, Baltimore County, Baltimore, Maryland 21250, USA*

(Received 14 September 2020; revised 19 February 2021; accepted 22 February 2021; published 18 March 2021)

This Letter presents a nonclassical mechanism of two-atom interferometry. The two-atom interferometer introduces two different yet indistinguishable alternatives for randomly paired atoms to create a joint atom-detection event. The superposition of two-atom amplitudes yields two-atom interference with peculiar features which outshine the classic atom interferometer: (1) two-atom interference is still observable when the time delay of the interferometer is greater than the coherence time of the atom beam and (2) two-atom interference may eliminate phase noises, including background, variation, turbulence, and Raman laser induced phase noises, thereby allowing for higher sensitivity and stability sensing than is achievable by classic atom interferometers. These features are crucial in high sensitivity-accuracy acceleration and rotation measurements. The presented concept and mechanism of two-atom interference can be adapted to other matter-wave interferometers, such as two-neutron interferometer and two-electron interferometer.

DOI: 10.1103/PhysRevA.103.L031303

A classic atom interferometer introduces two different yet indistinguishable paths for an atom to create an atom detection event at space-time $(\mathbf{r}, t)$. The superposition of the two quantum amplitudes yields an interference pattern in space-time which is interpreted as the probability of observing an atom at coordinate $(\mathbf{r}, t)$. In Dirac's language, the atomic interference phenomenon is the result of an atom interfering with the atom itself. A two-atom interferometer introduces two different yet indistinguishable paths for a pair of atoms to create a joint atom-detection event at space-time $(\mathbf{r}_1, t_1; \mathbf{r}_2, t_2)$. According to quantum mechanics, the superposition of two-atom amplitudes yields a two-atom interference pattern in space-time, which is interpreted as the probability of jointly observing two atoms at coordinate $(\mathbf{r}_1, t_1; \mathbf{r}_2, t_2)$. It can be said this interference phenomenon is the result of a pair of atoms interfering with the pair itself. In this Letter, we propose, study, and analyze the mechanism of two-atom interferometry, which can be adapted to other matter-wave interferometers, such as two-neutron interferometer and two-electron interferometer: a pair of matter-wave interference with the pair itself.

Two-photon interference phenomenon has been studied since the time of Hanbury-Brown and Twiss [1]. The formalism for second-order coherence of radiation developed by Glauber [2] and the research of entangled two-photon systems significantly advanced the development of two-photon interferometry and discovered interesting and attractive features of two-photon interference [3]. A recent study of two-photon interference of a thermal field has also explored attractive features of a two-photon interferometer for high sensitivity-accuracy measurements in "noisy" environments [4]. Advances in laser cooling and trapping have led to interest in higher order coherence correlations of atomic wave packets [5,6]. Two-atom correlations were first demonstrated in an ultracold atomic beam in 1996 [7] and later for ultracold atoms and Bose-Einstein condensates (BEC) [8–11] and atoms in an optical lattice [12].

The development of atom interferometers which yield significantly higher sensitivity to external potentials than photon interferometers [13–17] leads to the question of whether it is possible to adapt the mechanism of two-photon interference of thermal light to two-atom interferometry. Can randomly paired atoms in thermal state interfere with the pair itself, and if so, is this interference measurable by means of a two-atom interferometer? If possible, a two-atom interferometer should achieve much higher sensitivity and accuracy in acceleration and/or rotation sensing than is achievable by a two-photon interferometer.

Figure 1 schematically illustrates the working mechanism of the proposed two-atom light-pulse atom interferometer, which is the same as a classic one-atom interferometer [18–20], except for the use of two detectors for joint atom-detection of two atoms. This type of interferometer uses optical fields to create two different yet indistinguishable probability amplitudes for an atom-detection event or a joint atom-detection event of two atoms. First, the one-atom interferometer will be described, followed by the two-atom interferometer and its unique properties. The interferometer uses three excitation lasers, L_1, L_2, and L_3, corresponding to $\pi/2$, π, and $\pi/2$ Raman pulses, analogous in purpose to the matter-based beam splitters and mirrors of optical interferometers [21,22].

Assume the interferometer is sourced by a continuous atom beam transversely cooled with respect to its propagation direction [23]; the beam contains a large ensemble of randomly distributed atoms in their ground state with an initial momentum distribution about momentum $\mathbf{p}_0$. In the basis for which the atomic state is a tensor product of internal energy and momentum Hilbert spaces [21], the state of the mth atom

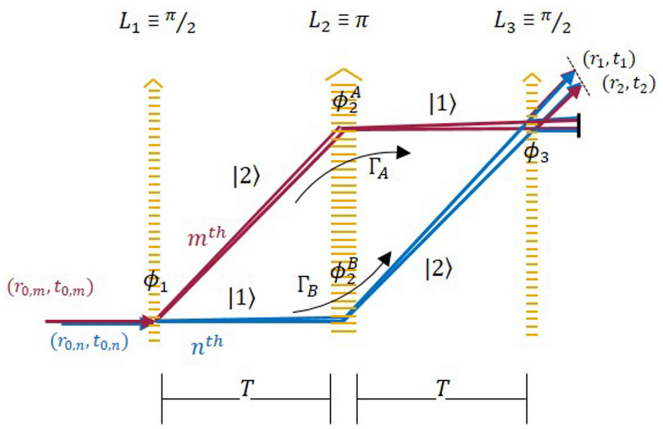

FIG. 1. Two-atom $\pi/2$-π-$\pi/2$ light pulse interferometer in which the cross-correlation terms are independent of path-dependent phase noise and the one-atom phase contrast is eliminated.

is given as $|1\rangle \otimes |\mathbf{p}_{1,m}\rangle$ in the ground state and $|2\rangle \otimes |\mathbf{p}_{2,m}\rangle$ in its excited state. Due to the interaction with the spatial dependence of the electric field, absorption or emission of a photon results in a momentum transfer of $\hbar\mathbf{k}$ with $\hbar\mathbf{k} \perp \mathbf{p}_0$, as shown in Fig. 1. With the normalized complex amplitudes, $c_{1,m}$ and $c_{2,m}$, the state of the mth atom is

$$|\Psi_m(t)\rangle = c_{1,m}(t)\,|1\rangle \otimes |\mathbf{p}_{0,m}\rangle + c_{2,m}(t)\,|2\rangle \otimes |\mathbf{p}_{0,m} + \hbar\mathbf{k}\rangle\,, \tag{1}$$

where $\mathbf{p}_{0,m}$ is the initial momentum of the atom. The probability amplitudes prior to the first laser field are $c_{1,m} = 1$ and $c_{2,m} = 0$.

Due to the interaction with the first laser field, the wave function is superposed into both the ground and excited states. The ground-state wave function keeps its propagation direction (path Γ_B); however, the excited state probability amplitude is kicked into a slightly different direction (path Γ_A). During propagation along Γ_A or Γ_B, the wave packet acquires phase contributions due to the interaction with the excitation lasers, designated ϕ_1, ϕ_2, and ϕ_3, as well as during free evolution between the laser fields. Note that the phase contribution from L_2 is added into both the upper path ϕ_2^A and lower path ϕ_2^B trajectories.

In the interaction picture, without considering the contribution from the spatial dependence of the electric field and in the limit that spontaneous emission is negligible, after the first $\pi/2$ light pulse with time duration τ, the mth atom is in a coherent superposition of $|1\rangle$ and $|2\rangle$ as follows:

$$|\Psi_m(\tau)\rangle = c_{1,m}(\tau)\,|1\rangle + c_{2,m}(\tau)e^{-i(\phi_{1,m}+\pi/2)}\,|2\rangle\,. \tag{2}$$

In the idealized situation, $c_1(\tau) = c_2(\tau) = 1/\sqrt{2}$. The excited state acquired an addition phase term due to the atom-photon interaction while the phase of the ground state is left unaffected by the first laser field.

Following a period of free evolution, T, the wave function interacts with the second laser field. Unlike the first field, which superposes the wave function into both the ground and excited states, the Raman field ideally drives a π pulse resulting in complete inversion of the probability amplitude from c_2 to c_1 for Γ_A and vice versa for Γ_B. Along Γ_A,

$$\left|\Psi_m^{\Gamma_A}(3\tau + T)\right\rangle = c_{1,m}(3\tau + T)\,e^{-i(\phi_{1,m}-\phi_{2,m}^A)}\,|1\rangle\,, \tag{3}$$

where the wave function is redirected along the initial momentum vector. On path Γ_B, the wave function becomes

$$\left|\Psi_m^{\Gamma_B}(3\tau + T)\right\rangle = c_{2,m}(3\tau + T)\,e^{-i(\phi_{2,m}^B+\pi/2)}\,|2\rangle \tag{4}$$

and acquires a momentum component through the absorption of a photon.

After a second period of free evolution, T, the wave function undergoes another $\pi/2$ pulse, further superposing each arm into both the ground- and excited-state output ports of the interferometer. The wave function of the excited state output port (which is the only port that will be considered here), at time $t = 4\tau + 2T$, has contributions from both Γ_A and Γ_B given separately by

$$\left|\Psi_m^{\Gamma_A}(t)\right\rangle = c_{2,m}^A(t)e^{-i(\phi_{1,m}-\phi_{2,m}^A+\phi_{3,m})}\,|2\rangle\,, \tag{5}$$

$$\left|\Psi_m^{\Gamma_B}(t)\right\rangle = c_{2,m}^B(t)e^{-i\phi_{2,m}^B}\,|2\rangle\,. \tag{6}$$

This result will be used throughout the remainder of this paper and therefore several substitutions are made for simplification. The upper and lower path laser phase contributions for the mth atom will be written as $\phi_{L,m}^A \equiv \phi_{1,m} + \phi_{3,m} - \phi_{2,m}^A$ and $\phi_{L,m}^B \equiv \phi_{2,m}^B$ respectively (the common $\pi/2$ phase factor has been omitted). The laser contributions result in a constant phase shift between the two arms of the interferometer dependent on the relative phases of the three Raman fields. Additionally, the phase evolves during the periods, T, between the Raman fields according to the classical action [24,25]. The accumulated phase ϕ^{Γ_i} for $i = A, B$ is path dependent and defined by

$$\phi^{\Gamma_i} = \frac{1}{\hbar}\int_{\Gamma_i} L(\mathbf{r}, \dot{\mathbf{r}}; t)dt, \tag{7}$$

where $L(r, \dot{r}; t)$ is the Lagrangian of the kinematic motion of the atom in the limit that the path length is large compared to $\hbar$. This results in the path-dependent phase accumulation for the mth atom

$$\phi^{\Gamma_i} = \int_{\Gamma_i} (\mathbf{k}_m \cdot d\mathbf{r} - \omega_m dt), \tag{8}$$

where $\mathbf{p} = \hbar\mathbf{k}_m$ is the kinetic momentum and $E_m = \hbar\omega_m$ is the kinetic energy of the atom. In measurements of inertial potentials, the output phase of the interferometer relies on the ability to measure the differences in phase evolution between ϕ^{Γ_A} and ϕ^{Γ_B} at the detection point. Therefore, the full propagation over Γ_A and Γ_B can be written in terms of the phase contributions from the first excitation field to the closure at the third field. The state $|\Psi_m(\mathbf{r}, t)\rangle$ of the mth atom at space-time coordinate $(\mathbf{r}, t)$ of the atom-detection event is thus the superposition of all possible interferometer paths, $i = A, B$:

$$\begin{aligned}|\Psi_m(\mathbf{r}, t)\rangle &= e^{-i\phi_{0m}}\sum_i \left|\Psi_m^{\Gamma_i}(\mathbf{r}, t)\right\rangle \\ &= c_{2,m}^A(t)e^{-i(\phi_{L,m}^A+\phi_m^{\Gamma_A})}e^{-i\phi_{0m}}\,|2\rangle \\ &\quad + c_{2,m}^B(t)e^{-i(\phi_{L,m}^B+\phi_m^{\Gamma_B})}e^{-i\phi_{0m}}\,|2\rangle\,,\end{aligned} \tag{9}$$

where $e^{-i\phi_{0m}}$ is the initial random phase of the mth atom.

Now, consider the state of a mixed ensemble of atoms which can be involved in an atom-detection event or a joint atom-detection event,

$$\begin{aligned}|\Psi(\mathbf{r},t)\rangle &= \prod_m (|0\rangle + \epsilon\,|\Psi_m(\mathbf{r},t)\rangle)\\ &\simeq |0\rangle + \epsilon \sum_m |\Psi_m(\mathbf{r},t)\rangle\\ &\quad + \epsilon^2 \sum_{m,n} |\Psi_m(\mathbf{r},t)\rangle\,|\Psi_n(\mathbf{r},t)\rangle + \cdots, \end{aligned} \tag{10}$$

where $\epsilon \ll 1$, $|0\rangle$ results in no detection and m runs over the total number of atoms. Note that the state is written in terms of the interaction picture.

The effective wave function of an atom is defined by considering the probability of an atom observation, or detection, event occurring at $(\mathbf{r}, t)$:

$$\begin{aligned}P(\mathbf{r},t) &= \langle\langle\Psi(\mathbf{r},t)|\Psi(\mathbf{r},t)\rangle\rangle_{En}\\ &\simeq \left\langle \sum_n \langle\Psi_n(\mathbf{r},t)| \sum_m |\Psi_m(\mathbf{r},t)\rangle \right\rangle_{\mathrm{En}}\\ &= \left\langle \sum_{m,n} \langle\Psi_n(\mathbf{r},t)| \sum_f |f\rangle\langle f|\Psi_m(\mathbf{r},t)\rangle \right\rangle_{\mathrm{En}}\\ &= \sum_m |\langle 2|\Psi_m(\mathbf{r},t)\rangle|^2\\ &\equiv \sum_m |\psi_m(\mathbf{r},t)|^2, \end{aligned} \tag{11}$$

where the outer bracket denotes the ensemble average and

$$\psi_m(\mathbf{r},t) \equiv \langle 2|\Psi_m(\mathbf{r},t)\rangle \tag{12}$$

is defined as the effective wave function of the mth atom. In the above derivation, the identity operator $\mathbb{1} = \sum_f |f\rangle\langle f| = |0\rangle\langle 0| + |1\rangle\langle 1| + |2\rangle\langle 2|$ is used.

Utilizing the effective wave functions defined above, the probability of observing an atom at space-time coordinate $(\mathbf{r}, t)$ is as follows:

$$\begin{aligned}P(\mathbf{r},t) &= \langle|\Psi(\mathbf{r},t)|^2\rangle_{\mathrm{En}}\\ &= \left\langle \sum_n \psi_n^*(\mathbf{r},t) \sum_m \psi_m(\mathbf{r},t) \right\rangle_{\mathrm{En}}\\ &= \left\langle \sum_{m=n} \psi_m^*(\mathbf{r},t)\psi_m(\mathbf{r},t) \right\rangle_{\mathrm{En}}\\ &\quad + \left\langle \sum_{m\neq n} \psi_n^*(\mathbf{r},t)\psi_m(\mathbf{r},t) \right\rangle_{\mathrm{En}}\\ &\propto \langle n(\mathbf{r},t)\rangle + \langle\Delta n(\mathbf{r},t)\rangle\\ &= \langle n(\mathbf{r},t)\rangle, \end{aligned} \tag{13}$$

where $\langle n(\mathbf{r},t)\rangle$ is the mean atom number and $\langle\Delta n(\mathbf{r},t)\rangle$ is the atom number fluctuation. The atom number fluctuation is averaged to zero by taking into account all possible random phases of the wave packets. Expanding the effective wave function into its upper and lower path components and assuming $c_{2,m}^A$ and $c_{2,m}^B$ are equal, the mean atom number is thus

$$\begin{aligned}\langle n(\mathbf{r},t)\rangle &\propto \left\langle \sum_m \left|\psi_m^A(\mathbf{r},t) + \psi_m^B(\mathbf{r},t)\right|^2 \right\rangle_{\mathrm{En}}\\ &\propto [1 + \cos(\Delta\phi_\Gamma + \Delta\phi_L)], \end{aligned} \tag{14}$$

where $\Delta\phi_\Gamma = (\phi^{\Gamma_B} - \phi^{\Gamma_A})$ and $\Delta\phi_L = (\phi_L^B - \phi_L^A)$.

The modulation depends on both the differences in phase path evolution $\Delta\phi_\Gamma$ and phase accumulated from interaction with the Raman fields for each path $\Delta\phi_L$. $\Delta\phi_\Gamma$ is the only phase term that carries information about rotation and acceleration of the interferometer and thus the expected observable. The other phases such as $\Delta\phi_L$ are at best a constant offset. However, in the case for which there is path-dependent random phase variations, for instance, on the Raman fields, this noise is indistinguishable from the measurement of $\Delta\phi_\Gamma$ and leads to a reduction in the interference contrast [26], equivalent to path-dependent turbulence. The phase shifts in the one atom interferometer are limited by random path-dependent phase fluctuations, namely laser noise. Additionally, as the interference for the thermal state is the result of a single atom interfering with itself, the path difference between the arms of the interferometer $\Gamma_B - \Gamma_A$ must be less than the coherence length of the atom l_c.

Interestingly, our recent work on two-photon interferometry [4] has shown that these path-dependent random-phase variations do not contribute to the two-photon interference modulation (namely, the two-photon interferometer is free of turbulence) and that the path difference $\Gamma_B - \Gamma_A$ may be larger than the coherence length of an individual atom. This Letter will show a two-atom interferometer and by extension other two-matter-wave interferometers exhibit similar turbulence-free nature.

The second-order correlation measurement involves the joint two-atom observation at space-time points $(\mathbf{r}_1, t_1)$ and $(\mathbf{r}_2, t_2)$. The second-order correlation measurement results in a product of the first-order interferences, if the first-order measurements are both first-order coherent; however, this is not the interest of this Letter. The interest here is for the case in which no interference is observable from the classic first-order measurements. The following shows that even if both first-order measurements have lost their interferences due to noisy environments, or the atom beam itself is first-order incoherent, the second-order atom number fluctuation measurement still produces observable two-atom interference.

The probability of a joint atom-observation or atom-detection event occurring at $(\mathbf{r}_1, t_1)$ and $(\mathbf{r}_2, t_2)$ is calculated as follows:

$$\begin{aligned}&P(\mathbf{r}_1,t_1;\mathbf{r}_2,t_2)\\ &= \left\langle \left|\sum_m \psi_m(\mathbf{r}_1,t_1)\right|^2 \left|\sum_p \psi_p(\mathbf{r}_2,t_2)\right|^2 \right\rangle_{\mathrm{En}}\\ &= \left\langle \sum_{m,n,p,q} \psi_m^*(\mathbf{r}_1,t_1)\psi_n(\mathbf{r}_1,t_1)\psi_p^*(\mathbf{r}_2,t_2)\psi_q(\mathbf{r}_2,t_2) \right\rangle_{\mathrm{En}}, \end{aligned} \tag{15}$$

in which m, n, p, and q each label individual atoms of the ensemble. When taking into account all possible random phases

of the wave packets, the only cross-correlation terms which survive are for $m = q$ and $n = p$ and the self-correlation terms survive for $m = n$ and $p = q$. The second-order correlation function reduces to

$$\begin{aligned} &P(\mathbf{r}_1, t_1; \mathbf{r}_2, t_2) \\ &= \left\langle \sum_m \psi_m^*(\mathbf{r}_1, t_1)\psi_m(\mathbf{r}_1, t_1) \sum_n \psi_n^*(\mathbf{r}_2, t_2)\psi_n(\mathbf{r}_2, t_2) \right\rangle_{\text{En}} \\ &\quad + \left\langle \sum_{m\neq n} \psi_m^*(\mathbf{r}_1, t_1)\psi_n(\mathbf{r}_1, t_1)\psi_n^*(\mathbf{r}_2, t_2)\psi_m(\mathbf{r}_2, t_2) \right\rangle_{\text{En}} \\ &\propto \langle n(\mathbf{r}_1, t_1)\rangle\langle n(\mathbf{r}_2, t_2)\rangle + \langle \Delta n(\mathbf{r}_1, t_1)\Delta n(\mathbf{r}_2, t_2)\rangle, \end{aligned} \tag{16}$$

indicating a set of two-atom superposition between the following two-atom amplitudes: (1) the mth atom is observed at $(\mathbf{r}_1, t_1)$ while the nth atom is observed at $(\mathbf{r}_2, t_2)$ and (2) the mth atom is observed at $(\mathbf{r}_2, t_2)$ while the nth atom is observed at $(\mathbf{r}_1, t_1)$. In Eq. (16), the cross terms of the two-atom superposition results in the atom-number fluctuation correlation $\langle \Delta n(\mathbf{r}_1, t_1)\Delta n(\mathbf{r}_2, t_2)\rangle$, while the mean atom number at $(\mathbf{r}_1, t_1)$ and $(\mathbf{r}_2, t_2)$ contributes a constant to the joint-atom measurement. In the first-order incoherent case, the atom number fluctuation correlation still results in a two-atom interference modulating based on the phase difference between the two paths.

Considering the two indistinguishable interferometer paths, the atom-number fluctuation correlation at the excited state output port of the interferometer is calculated as follows:

$$\begin{aligned} &\langle \Delta n(\mathbf{r}_1, t_1)\Delta n(\mathbf{r}_2, t_2)\rangle \\ &\propto \left\langle \sum_{m\neq n} \psi_m^*(\mathbf{r}_1, t_1)\psi_n(\mathbf{r}_1, t_1)\psi_n^*(\mathbf{r}_2, t_2)\psi_m(\mathbf{r}_2, t_2) \right\rangle_{\text{En}} \\ &= \Big\langle \sum_{m\neq n} \left[\psi_m^{A*}(\mathbf{r}_1, t_1) + \psi_m^{B*}(\mathbf{r}_1, t_1)\right]\left[\psi_n^{A}(\mathbf{r}_1, t_1) + \psi_n^{B}(\mathbf{r}_1, t_1)\right] \\ &\quad \times \left[\psi_n^{A*}(\mathbf{r}_2, t_2) + \psi_n^{B*}(\mathbf{r}_2, t_2)\right]\left[\psi_m^{A}(\mathbf{r}_2, t_2) + \psi_m^{B}(\mathbf{r}_2, t_2)\right] \Big\rangle, \end{aligned} \tag{17}$$

which yields 16 ensemble averages to calculate. In the case of first-order measurements for which the contrast is completely eliminated, potentially by the condition $(\Gamma_B - \Gamma_A) > l_c$, only four terms survive:

$$\begin{aligned} &\langle \Delta n(\mathbf{r}_1, t_1)\Delta n(\mathbf{r}_2, t_2)\rangle \\ &\quad = \langle \Delta n_{AB}(\mathbf{r}_1, t_1)\Delta n_{AB}(\mathbf{r}_2, t_2)\rangle \\ &\quad + \langle \Delta n_{BA}(\mathbf{r}_1, t_1)\Delta n_{BA}(\mathbf{r}_2, t_2)\rangle \\ &\quad + \langle \Delta n_{AA}(\mathbf{r}_1, t_1)\Delta n_{AA}(\mathbf{r}_2, t_2)\rangle \\ &\quad + \langle \Delta n_{BB}(\mathbf{r}_1, t_1)\Delta n_{BB}(\mathbf{r}_2, t_2)\rangle. \end{aligned} \tag{18}$$

The first subscript designates the path of the mth atom while the second designates the nth atom path. Therefore, the first term in Eq. (18) is the cross-interference term of the two-atom superposition

$$\sum_{m\neq n} \left|\psi_m^A(\mathbf{r}_1, t_1)\psi_n^B(\mathbf{r}_2, t_2) + \psi_m^A(\mathbf{r}_2, t_2)\psi_n^B(\mathbf{r}_1, t_1)\right|^2, \tag{19}$$

which corresponds to the superposition of the following two different yet indistinguishable two-atom amplitudes: (1) mth atom taking Γ_A to $(\mathbf{r}_1, t_1)$ and the nth atom taking Γ_B to $(\mathbf{r}_2, t_2)$, and (2) the mth atom taking Γ_A to $(\mathbf{r}_2, t_2)$ and the nth atom taking Γ_B to $(\mathbf{r}_1, t_1)$, as shown in Fig. 1. The second term is the same as the first, except for the exchange of the paths for atom m and n. The third term and four term contribute a constant to the measurement of atom number fluctuation correlation.

In terms of the path-dependent interferometer phases, the quantum atom-number fluctuation for the joint detection measurement is

$$\begin{aligned} &\langle \Delta n(\mathbf{r}_1, t_1)\Delta n(\mathbf{r}_2, t_2)\rangle \\ &\quad \propto 1 + \cos[\Delta\phi_\Gamma(\mathbf{r}_2, t_2) - \Delta\phi_\Gamma(\mathbf{r}_1, t_1)]. \end{aligned} \tag{20}$$

Notably, (20) only includes terms which are dependent on the phase difference between the two paths as calculated by (7). The measurement is independent of the relative phases of the Raman lasers as well as any path-dependent noise over Γ_A and Γ_B. This cancellation is due to the measurement of the two-atom wave packet traversing the same paths, simultaneously allowing for full phase contrast in spite of noise, which would normally result in error in the interferometer readout making the sensor useless for tracking inertial potentials. This case is distinctly different from the first-order measurement (14) for which the Raman phases and any other path-dependent phases contribute fully to the interferometer measurement resulting in contrast loss. In the intermediate case for which the atom beam is not fully first-order incoherent, more terms than given in (18) will contribute to the atom number fluctuation correlation with varying degrees of contrast.

The benefit to using an atom interferometer versus a photon interferometer lies in the comparison of the mass of an atom versus the relative mass of a photon. The corresponding phase change under a rotation rate of $\mathbf{\Omega}$ is

$$\Delta\phi_{\text{rot}} = \frac{2M}{\hbar}\mathbf{A}\cdot\mathbf{\Omega}, \tag{21}$$

where M is the mass of the particle and $\mathbf{A}$ is the enclosed area. A comparison of a rubidium atom to the relative mass of a photon leaves the atom interferometer with a factor of $\approx 10^{10}$ times larger phase shift to rotation rate. There is further benefit in the two-atom interference measurement over the one-atom interferometer considering the decreasing coherence of atoms as they approach higher temperatures and velocity spread which is common for atomic beams. In a classic one-atom interferometer, the path difference between its two arms must be less than the coherence length of the atom beam. However, for the two-atom interferometer, the path difference can exceed the coherence length of the atom beam. Two-photon interference for path differences well beyond the coherence length of thermal field has already been described in detail [27] and demonstrated [28], and by application this would allow for atom interferometry measurements in which the path length difference is significantly larger than the atomic coherence length. In a two-atom interferometer, $\mathbf{A}$ may achieve orders greater than that of a one-atom interferometer because of its peculiar nature of first-order coherence independence. The ability to increase the area of an atom interferometer makes

the two-atom interferometer unique for the idea of precision measurements via matter waves.

In summary, (1) two-atom interference is observable with the time delay and/or when spatial separation of the interferometer is greater than the spatial or temporal coherence of the atom beam; (2) the joint two-atom number fluctuation measurement is independent of the same contrast-reducing phase fluctuations found in the first-order measurement; and (3) the sinusoidal modulation of the two-atom interferometer is a function of the temporal or spatial separation of the two atom detection events $(\mathbf{r}_1, t_1; \mathbf{r}_2, t_2)$ while the one-atom interferometer is a function of the temporal or spatial coordinate of the one-atom detection event. The presented concept and mechanism of two-atom interference can be adapted to other matter-wave interferometers, such as two-neutron interferometry and two-electron interferometry. The concept and mechanism of two-atom and two-matter-wave interference is fundamentally interesting and we are currently implementing it experimentally. While significant experimental work is still required to achieve sensor measurements for thermal two-atom and two-matter-wave interferometers, the potential for high-sensitivity measurements with path-dependent phase noise that renders one-atom and one-matter-wave interferometry unfeasible is a worthwhile and interesting endeavor. Given the high sensitivity of atoms to inertial potentials, the two-atom and two-matter-wave interferometer sensors would outperform both the classic atom and matter-wave interferometers for sensing applications.

[1] R. Hanbury Brown and R. Q. Twiss, Nature (London) **177**, 27 (1956).

[2] R. J. Glauber, Phys. Rev. **131**, 2766 (1963).

[3] Y. H. Shih, *An Introduction to Quantum Optics: Photons and Biphoton Physics*, 1st ed. (CRC Press, Boca Raton, FL, 2011).

[4] T. A. Smith and Y. Shih, Phys. Rev. Lett. **120**, 063606 (2018).

[5] J. Viana Gomes, A. Perrin, M. Schellekens, D. Boiron, C. I. Westbrook, and M. Belsley, Phys. Rev. A **74**, 053607 (2006).

[6] M. Naraschewski and R. J. Glauber, Phys. Rev. A **59**, 4595 (1999).

[7] M. Yasuda and F. Shimizu, Phys. Rev. Lett. **77**, 3090 (1996).

[8] M. Schellekens, R. Hoppeler, A. Perrin, J. Viana Gomes, D. Boiron, A. Aspect, and C. I. Westbrook, Science **310**, 648 (2005).

[9] D. Hellweg, L. Cacciapuoti, M. Kottke, T. Schulte, K. Sengstock, W. Ertmer, and J. J. Arlt, Phys. Rev. Lett. **91**, 010406 (2003).

[10] A. Öttl, S. Ritter, M. Köhl, and T. Esslinger, Phys. Rev. Lett. **95**, 090404 (2005).

[11] S. Flling, F. Gerbier, A. Widera, O. Mandel, T. Gericke, and I. Bloch, Nature (London) **434**, 481 (2005).

[12] J. Grondalski, P. Alsing, and I. Deutsch, Opt. Express **5**, 249 (1999).

[13] C. Monroe, W. Swann, H. Robinson, and C. Wieman, Phys. Rev. Lett. **65**, 1571 (1990).

[14] D. J. Wineland and H. Dehmelt, Bull. Am. Phys. Soc. **20**, 637 (1975).

[15] T. W. Hänsch and A. L. Schawlow, Opt. Commun. **13**, 68 (1975).

[16] A. Ashkin, Phys. Rev. Lett. **40**, 729 (1978).

[17] S. Chu, L. Hollberg, J. E. Bjorkholm, A. Cable, and A. Ashkin, Phys. Rev. Lett. **55**, 48 (1985).

[18] C. J. Bordé, Phys. Lett. A **140**, 10 (1989).

[19] M. Kasevich and S. Chu, Phys. Rev. Lett. **67**, 181 (1991).

[20] A. Peters, K. Y. Chung, and S. Chu, Metrologia **38**, 25 (2001).

[21] B. Young, M. Kasevich, and S. Chu, in *Atom Interferometry*, 1st ed., edited by P. R. Berman (Academic Press, San Diego, 1997).

[22] A. J. Dunning, *Coherent Atomic Manipulation and Cooling*, Springer Theses (Springer International Publishing, Switzerland, 2015), 1st ed.

[23] J. Schoser, A. Batär, R. Löw, V. Schweikhard, A. Grabowski, Y. B. Ovchinnikov, and T. Pfau, Phys. Rev. A **66**, 023410 (2002).

[24] R. P. Feynman and A. R. Hibbs, *Quantum Mechanics and Path Integrals* (McGraw-Hill, New York, 1965).

[25] A. Peters, K. Y. Chung, B. Young, J. Hensley, and S. Chu, Philos. Trans. Math., Phys. Eng. Sci. **355**, 2223 (1997).

[26] A. Lenef, T. D. Hammond, E. T. Smith, M. S. Chapman, R. A. Rubenstein, and D. E. Pritchard, Phys. Rev. Lett. **78**, 760 (1997).

[27] Y. H. Shih, *An Introduction to Quantum Optics: Photons and Biphoton Physics*, 2nd ed. (CRC Press, Boca Raton, FL, 2021), Chap. 11.

[28] Y. S. Ihn, Y. Kim, V. Tamma, and Y.-H. Kim, Phys. Rev. Lett. **119**, 263603 (2017).