

Coupling Free Electrons to a Trapped-Ion Quantum Computer

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Freely propagating electrons may serve as quantum probes that can become coherently correlated with other quantum systems, offering access to advanced metrological resources. We propose a setup that coherently couples free electrons in an electron microscope to a trapped-ion quantum processor, enabling nondestructive, quantum-coherent detection and the accumulation of information across multiple electrons. Our analysis shows that single electrons can induce resolvable qubit excitations, establishing a platform for practical applications such as quantum-enhanced, dose-efficient electron microscopy.

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Electron microscopy (EM) [1] has reached extraordinary spatial [2,3] and temporal [4–8] resolution, enabling direct visualization of matter at the atomic scale. These advances have been driven by major developments in electron optics, including aberration correction and high-brightness sources [2,9], revolutions in electron detection [2,10–12], and increasingly sophisticated computational reconstruction techniques [13]. As a result, EM has become a powerful, high-throughput imaging method in both biology [14,15] and materials science [16]. In the near future, further advances may be provided, for example, by laser-based electron wave packet shaping, which has already been experimentally demonstrated [17–19].

Despite this progress, the information that can be extracted from a specimen remains fundamentally constrained by the interaction between the electron beam and

the sample. Inelastic scattering processes induce structural and chemical modifications, imposing a finite tolerable electron dose [20–23] and, consequently, an upper bound on the achievable signal-to-noise ratio [24,25]. This limitation is particularly severe for radiation-sensitive specimens such as living cells, vitrified biological material, viruses, and individual proteins. Improving EM under such dose constraints therefore requires strategies that increase the information gained per electron, rather than simply increasing beam intensity.

At the same time, electron microscopy has largely operated within a semiclassical measurement paradigm, in which quantum correlations involving the probing electrons are neglected or treated as detrimental. Yet, the entanglement of electrons with other quantum systems, including electromagnetic fields, detectors, and in principle the specimen itself, may be harnessed to open up an additional and largely unexplored avenue for enhancing metrological performance, enabling new measurement modalities and protocols that go beyond classical bounds [26–34].

Recent theoretical proposals and experimental demonstrations illustrate how both challenges, i.e., low-dose operation and quantum-enhanced measurements, can be addressed: approaches such as employing entangled multi-electron probe states [35], coherent multipass electron microscopy [36], interaction-free measurements [37–40], entanglement-assisted phase contrast [41], and quantum search-inspired measurement protocols [42] aim to maximize information extraction per damaging interaction.

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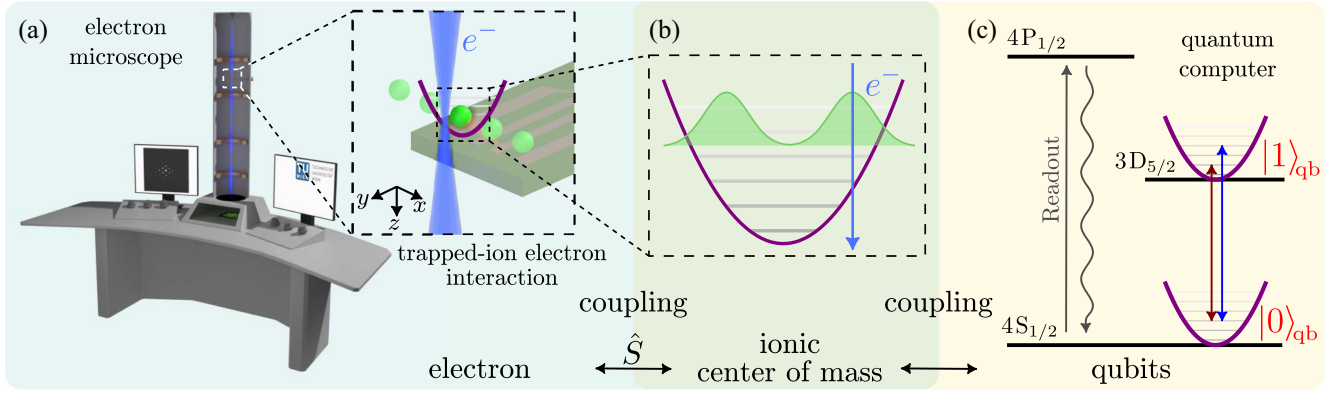


FIG. 1. (a) Transmission electron microscope (TEM). Within a plane conjugate to the specimen plane, a trapped-ion quantum computer is inserted. (b) The electron e^- couples to the ionic center-of-mass degree of freedom via the Coulomb interaction (scattering matrix $\hat{S}$). If the ion is prepared in a superposition of coherent states, the electron induces a relative phase shift between these states. (c) Simplified level diagram of $^{40}\text{Ca}^+$ with the relevant transitions used to encode the qubit. The center-of-mass degree of freedom of the ion is coupled to the $4S_{1/2}$ and $3D_{5/2}$ states of the valence electron (qubit degree of freedom) via the application of bichromatic laser pulses (red- and blue-detuned transitions in the figure). The rapidly decaying fluorescent transition from $4P_{1/2}$ to $4S_{1/2}$ is used to perform projective readout of the qubit state.

Complementarily, coherent interactions of free electrons with other quantum systems such as plasmons [43], spins [44–48], and photons [49–54] including electron-photon entanglement [55,56] have been experimentally demonstrated, highlighting the feasibility of coherent electron-quantum-system interfaces. Together, these developments point toward a new regime of quantum-enhanced electron microscopy [29], in which dose efficiency and quantum metrology are intrinsically linked.

Motivated by the rapid progress of quantum information processors, we propose a concrete implementation of this paradigm by coherently coupling free electrons in an electron microscope to a trapped-ion quantum computer. We analyze the interaction between electrons and trapped ions and propose a protocol to entangle free electrons with trapped-ion qubits. Beyond enabling the preparation of advantageous quantum probe states, this interface grants access to the full space of quantum measurements through the universal state-manipulation capabilities of the quantum processor, establishing a versatile platform for low-dose, quantum-enhanced electron microscopy. The proposed platform can also be useful for coherent state preparation and readout in other setups that employ free electrons, for example, quantum free electron lasers [57,58] and nano-scale electron accelerators [59]. Furthermore, the coupling scheme may be generalized to different highly controlled charged particles ranging from focused ion beams to the protons used in high energy physics.

In the following, we assume a focused electron wave function, as commonly produced and coherently shaped within a transmission electron microscope (TEM) [see Fig. 1(a)], interacting with ions trapped in a harmonic oscillator potential. A promising ion-trap platform consists of a planar surface-electrode Paul trap on top of an integrated photonic chip, which both confines $^{40}\text{Ca}^+$ ions around hundred micrometers above the surface and

manipulates them with lasers through grating couplers [see Fig. 1(a); more details on integrating an ion trap experiment into a TEM are given in Supplemental Material (SM) Sec. A [60]]. This choice has the advantage that the ions can be separated by several hundred micrometers from each other [78], allowing for the electron to couple to single ions in a larger quantum register. The ion trap includes a sufficiently large opening around the trapping region, allowing the electron beam to pass by. Typical trap frequencies Ω lie between 0.5 MHz and 5 MHz, which correspond to harmonic-oscillator ground-state sizes $R_0/\sqrt{2}$ of 40 nm and 13 nm, respectively, where $R_0 = \sqrt{\hbar/m_{\text{ion}}\Omega}$. After resolved-sideband cooling, the ion can be initialized in its motional and electronic ground state [79].

The $^{40}\text{Ca}^+$ ions provide an established platform for quantum information processing, including coherent manipulation on the narrow S-D quadrupole transition and projective readout on the S-P cycling transition [Fig. 1(c)]. These quantum operations have been demonstrated in compact experimental systems with more than 15 ions [80]. Operations between multiple qubits are performed by bringing the ions close together and coupling their qubit states to their common motion in a single trap potential. The same tools can be used to create nonclassical sensing states involving the qubit of a single ion and its motion [81].

The coupling between the electrons and the qubits is then mediated by the ionic center-of-mass motion, which couples directly to the free electron via the Coulomb interaction and to the qubits via the application of laser pulses; see Fig. 1. The focus of the present Letter is the analysis of the first stage of this coupling scheme described in the following by the scattering operator $\hat{S}$.

We first model the interaction of a single electron with a single trapped ion. We start from the Hamiltonian $\hat{H} = \hat{H}_{\text{el}} + \hat{H}_{\text{cmi}} + \hat{H}_{\text{qb}} + \hat{H}_{\text{int}}$, where $\hat{H}_{\text{el}} = c\sqrt{c^2 + \hat{\mathbf{p}}^2}$ is the relativistic energy of the free electron (in atomic units, $\hbar = m_e = |e| = 4\pi\epsilon_0 = 1$), $\hat{\mathbf{p}}$ is the electron momentum operator, $\hat{H}_{\text{cmi}}$ is the kinetic and potential energy of the center-of-mass degree of freedom of a harmonically trapped ion with trap frequencies $\Omega_x = \Omega_y = \Omega$ and Ω_z , and $\hat{H}_{\text{qb}} = \omega\hat{\sigma}_{\text{qb}}^z$ is the free Hamiltonian for the two-level system used to encode the qubit. The operator $\hat{H}_{\text{int}} = -1/|\hat{\mathbf{r}} - \hat{\mathbf{R}}|$ is the Coulomb interaction of the free electron ($\hat{\mathbf{r}}$) with the ionic center of mass ($\hat{\mathbf{R}}$). Note that the probability for internal transitions of the ion induced by the free electron is negligibly small in the relevant parameter regime (see SM Sec. E [60]). Therefore, we do not include the direct interaction of the electron with the internal degrees of freedom of the ion into the model.

In a TEM, the electron momentum is typically dominated by a large component in the direction of $\mathbf{e}_z$, where we fix the coordinate system such that the z axis is aligned with the beam axis [Fig. 1(a)]. To leading order in the relative transverse and longitudinal momentum spreads of the electron wave packet, this allows us to replace the electron momentum operator $\hat{\mathbf{p}}$ in the interaction Hamiltonian by its average $\mathbf{p}_{\text{el}} = \gamma v_{\text{el}} \mathbf{e}_z$ corresponding to the paraxial approximation. Furthermore, the interaction time with the ion is much shorter than the timescale of the ionic center-of-mass motion, such that we can use the stroboscopic approximation, i.e., neglecting the time evolution of the ion during the scattering event, to derive the scattering operator (see SM Sec. B [60]):

$$\hat{S} = \exp\left(-\frac{2i}{v_{\text{el}}} \log|\hat{\mathbf{r}}_{\perp} - \hat{\mathbf{R}}_{\perp}^{(1)}(t_{\text{el}})|\right). \quad (1)$$

Here, $\hat{\mathbf{r}}_{\perp}$ is the transverse position operator for the electron; $\hat{\mathbf{R}}_{\perp}^{(1)}(t_{\text{el}}) = \hat{\mathbf{R}}_{\perp} \cos(\Omega t_{\text{el}}) + \hat{\mathbf{P}}_{\perp}/m_{\text{ion}} \sin(\Omega t_{\text{el}})$ is the transverse position operator of the ion in the interaction picture; $\hat{\mathbf{R}}_{\perp}$ and $\hat{\mathbf{P}}_{\perp}$ are the transverse position and momentum operators of the ion, respectively; and t_{el} is the electron arrival time.

With the approximations made (paraxial and stroboscopic), the scattering operator only couples the transverse degrees of freedom of the electron and the ion. For the subsequent analysis, we thus express the joint wave function of the electron, the center-of-mass motion of the ion, and the qubits in terms of states of the form $|\psi_{\mathbf{r}_{\perp}}\rangle_{\text{el}}|\alpha\rangle_{\text{cmi}}|\sigma\rangle_{\text{qb}}$. Here, $|\psi_{\mathbf{r}_{\perp}}\rangle_{\text{el}}$ denotes a focused Gaussian electron state in the plane of the ion trap, with midpoint $\mathbf{r}_{\perp} = (x, y)$ and width $\delta_{\mathbf{r}_{\perp}}$. We assume the width $\delta_{\mathbf{r}_{\perp}}$ to be significantly smaller than the zero-point fluctuation of the trapped ion $\delta_{\mathbf{R}_{\perp}} = R_0/\sqrt{2} = 1/\sqrt{2m_{\text{ion}}\Omega}$. The ion center-of-mass state $|\alpha\rangle_{\text{cmi}}$ corresponds to a coherent state in the

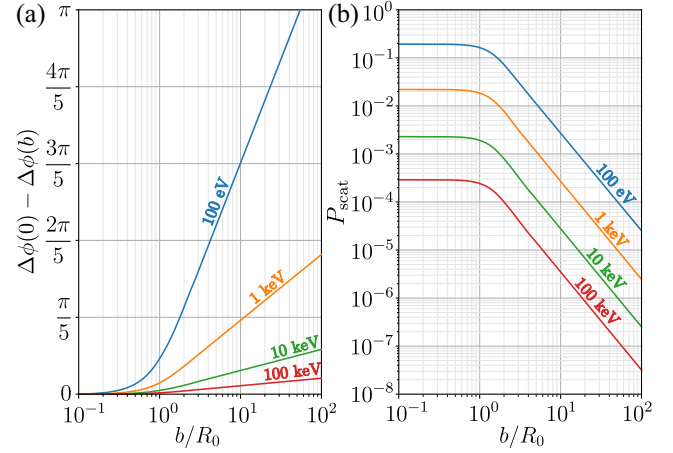


FIG. 2. (a) Phase shift $\Delta\phi(b)$ [Eq. (2)], assuming $\delta_{\mathbf{r}_{\perp}} \ll R_0$, of the combined electron-ion wave function due to the interaction with impact parameter b compared to the corresponding phase shift $\Delta\phi(0)$ when the electron is focused to the center of the harmonic trap. Note that the phase shift will approach a finite limit eventually due to screening of the ionic charge. Here only the phase shift for small relative impact parameters relevant to the proposed scheme is shown. (b) Probability $P_{\text{scat}} = 1 - |\mathcal{S}(\mathbf{r}_{\perp}, \alpha)|^2$, assuming $\delta_{\mathbf{r}_{\perp}} \ll R_0$, for the ion to be scattered out of state $|\alpha\rangle_{\text{cmi}}$ upon impact of an electron as a function of the impact parameter b .

harmonic trap with displacements $\alpha = (\alpha_x, \alpha_y)$. Finally, $|\sigma\rangle_{\text{qb}} \in \{|0\rangle_{\text{qb}}, |1\rangle_{\text{qb}}\}$ is the qubit encoded in the internal electronic degrees of freedom of the ion.

We now study the effect of this operator on the basis states of the form $|\psi_{\mathbf{r}_{\perp}}\rangle_{\text{el}}|\alpha\rangle_{\text{cmi}}|\sigma\rangle_{\text{qb}}$. As long as the spread of the electronic wave function is small enough ($\delta_{\mathbf{r}_{\perp}} \lesssim 0.1\delta_{\mathbf{R}_{\perp}} \lesssim 1\text{nm}$), the focused electronic states $|\psi_{\mathbf{r}_{\perp}}\rangle_{\text{el}}$ are approximate eigenstates of the scattering operator. In this setting, we may ignore any back action, including momentum kicks on the electron. In SM Sec. D [60] we discuss the effect of a finite extent of the electronic wave function and give upper bounds to the decoherence arising from the back action to the electron.

The main effect of the scattering operator Eq. (1) acting on coherent states is to imprint a displacement-dependent phase shift onto them:

$$\hat{S}|\psi_{\mathbf{r}_{\perp}}\rangle_{\text{el}}|\alpha\rangle_{\text{cmi}} \approx e^{i\Delta\phi(b)}|\psi_{\mathbf{r}_{\perp}}\rangle_{\text{el}}|\alpha\rangle_{\text{cmi}}, \quad (2)$$

with the impact parameter $b = |\mathbf{r}_{\perp} - \sqrt{2}R_0\Re(\alpha e^{i\Omega t_{\text{el}}})|$. The phase shift $\Delta\phi = \arg[\mathcal{S}(\mathbf{r}_{\perp}, \alpha)]$ is given by the complex argument of the state overlap of the initial and final state: $\mathcal{S}(\mathbf{r}_{\perp}, \alpha) = \langle\psi_{\mathbf{r}_{\perp}}|_{\text{el}}\langle\alpha|_{\text{cmi}}\hat{S}|\alpha\rangle_{\text{cmi}}|\psi_{\mathbf{r}_{\perp}}\rangle_{\text{el}}$. In Fig. 2(a) we plot $\Delta\phi$ as a function of the impact parameter and assuming $\delta_{\mathbf{r}_{\perp}} \ll R_0$, and we give an analytical expression for $\Delta\phi(b)$ in SM Sec. C [60]. There is also a finite probability for the scattering process to change the state of the ion by more than just a phase, given by $P_{\text{scat}} = 1 - |\mathcal{S}(\mathbf{r}_{\perp}, \alpha)|^2$ and plotted, again assuming $\delta_{\mathbf{r}_{\perp}} \ll R_0$, in Fig. 2(b). The probability of

this happening is approximately constant for impact parameters $b < R_0$, with $P_{\text{scat}} = 1 - \pi/[v_{\text{el}} \sinh(\pi/v_{\text{el}})]$ for when the electron is focused to the center of the trap ($b = 0$) (see SM Sec. C [60]), and it decays quadratically for larger impact parameters. For $b = 0$ and an electron energy of 100 eV the probability becomes approximately 10^{-1} , decreasing rapidly for higher energies. In the following, we focus on the induced phase shift and assume that the probability of state changes is negligible.

The phase imprinted onto the joint electron-ion wave function, Eq. (2), is a global phase that is undetectable. However, one may prepare the ion in a qubit-motion entangled cat state, that is, a superposition of two coherent states with displacements $-\alpha$ and α , entangled with the qubit degree of freedom: $|\psi_{\text{cat}}\rangle = (|-\alpha\rangle_{\text{cmi}}|-\rangle_{\text{qb}} + |\alpha\rangle_{\text{cmi}}|+\rangle_{\text{qb}})/\sqrt{2}$. Here, $|\pm\rangle_{\text{qb}} = (|0\rangle_{\text{qb}} \pm |1\rangle_{\text{qb}})/\sqrt{2}$ are the eigenstates of the Pauli-X operator $\hat{\sigma}_{\text{qb}}^x$. It has been experimentally shown that such a state can be prepared by using bichromatic laser pulses [see Fig. 1(c)], resulting in a qubit-dependent displacement [81]. This operation can be represented as

$$\hat{D}(\hat{\sigma}_{\text{qb}}^x \alpha) = \exp[\hat{\sigma}_{\text{qb}}^x (\alpha \cdot \hat{\mathbf{a}}^\dagger - \alpha^* \cdot \hat{\mathbf{a}})], \quad (3)$$

where $\mathbf{a}$ and $\mathbf{a}^\dagger$ are the (transverse) ladder operators in the harmonic trap. We assume the electron arrival to be synchronized with the harmonic trap such that, at time t_{el} , the ion is at its turning point: $\alpha e^{i\Omega t_{\text{el}}} \in \mathbb{R}^2$. For the sake of readability, we also fix $t_{\text{el}} = 0$, such that α is real. Then, the final ionic wave function after tracing out the electronic degrees of freedom becomes (up to a global phase)

$$|\psi'_{\text{cat}}\rangle = \frac{|-\alpha\rangle_{\text{cmi}}|-\rangle_{\text{qb}} + e^{ig(\mathbf{r}_\perp, \alpha)}|\alpha\rangle_{\text{cmi}}|+\rangle_{\text{qb}}}{\sqrt{2}}. \quad (4)$$

Here, $g(\mathbf{r}_\perp, \alpha) = \Delta\phi(|\mathbf{r}_\perp - \sqrt{2}R_0\alpha|) - \Delta\phi(|\mathbf{r}_\perp + \sqrt{2}R_0\alpha|)$ is the relative phase shift due to the difference in impact parameter. Thus, the final ionic state is distinguishable from the initial state, with $\arccos(|\langle\psi_{\text{cat}}|\psi'_{\text{cat}}\rangle|) = g(\mathbf{r}_\perp, \alpha)/2$. The quantum computer (in form of the ion) has thus acquired knowledge about the presence of an electron. The maximal coupling strength is reached when the electron beam is focused to the center of one of the two coherent states [see Fig. 1(b)]. For $\mathbf{r}_\perp = \sqrt{2}R_0\alpha$ we get $g(\sqrt{2}R_0\alpha, \alpha) = \Delta\phi(0) - \Delta\phi(2\sqrt{2}R_0|\alpha|)$. Because of the flat plateau at the origin and the logarithmic behavior at larger impact parameters of the relative phase [Fig. 2(a)], the coupling strength, $g(\sqrt{2}R_0\alpha, \alpha)$, is not very sensitive to precise alignment of the electron beam to the center of the coherent states or to precise control of $|\alpha|$.

After the interaction with the electron, the two coherent states can be recombined by inverting the operation Eq. (3). Altogether, assuming the ion's center of mass to initially be in the ground state, this results in an effective unitary

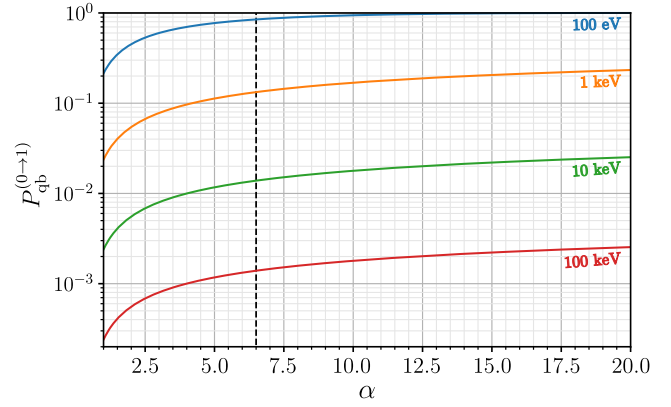


FIG. 3. Probabilities to find the qubit in state $|1\rangle$ after the interaction with an electron of different energies as a function of the cat state size α . It is assumed that the qubit is initially prepared in state $|0\rangle$. The dashed line indicates $|\alpha| = 6.5$, which has already been demonstrated experimentally [82].

$\hat{U}_{\text{e-qb}} = \langle 0|_{\text{cmi}} \hat{D}(-\hat{\sigma}_{\text{qb}}^x \alpha) \hat{S} \hat{D}(\hat{\sigma}_{\text{qb}}^x \alpha) |0\rangle_{\text{cmi}}$ that couples the qubit to the electron by performing a rotation on the qubit that depends on the position of the electron $\mathbf{r}_\perp$:

$$\hat{U}_{\text{e-qb}} \approx e^{i\kappa(\mathbf{r}_\perp, \alpha)} e^{\frac{i}{2}\hat{\sigma}_{\text{qb}}^x g(\mathbf{r}_\perp, \alpha)}. \quad (5)$$

The qubit-independent phase κ is thereby given by $\kappa(\mathbf{r}_\perp, \alpha) = [\Delta\phi(|\mathbf{r}_\perp - \sqrt{2}R_0\alpha|) + \Delta\phi(|\mathbf{r}_\perp + \sqrt{2}R_0\alpha|)]/2$. By implementing this protocol, it is thus possible to produce an effective coupling between the free probe electron and the bound electronic qubit. The strength of the coupling can be understood in terms of the probability to measure the qubit in the $|1\rangle_{\text{qb}}$ state after the interaction, assuming that it was initially prepared in the $|0\rangle_{\text{qb}}$ state. We plot this probability in Fig. 3. For slow electrons (100 eV–1 keV) and experimentally plausible values for $|\alpha|$, the probability for a bit flip reaches the order of 0.1–1. In an ideal, noiseless experiment, this probability would increase with larger $|\alpha|$. In practice, however, the measured signal is limited by atomic and motional decoherence since the sensitivity to noise increases with the cat state amplitude [81,82]. In SM Sec. A [60] we discuss the impact of experimental imperfections on our setup. While not usually employed in TEM, electron energies of 100 eV–1 keV are well matched to the technological advances of the last decades in electron microscopy. Atomic resolution has been demonstrated down to approximately 15 keV [83,84], and even lower electron energies, down to effectively 0 eV, are utilized in high-resolution EMs, achieving spatial resolutions down to 2 nm via beam acceleration and deceleration stages [85,86]. Spot sizes down to 2 nm can be reached in scanning electron microscopy [87], showing that it is possible to focus the electron to the center of the ionic wave function. Thus, by combining the setup with additional acceleration and deceleration stages it would be possible to reach the strongly interacting regime without sacrificing the

high resolution. Furthermore, in multipass geometries [36], the effect due to multiple interactions can add up coherently, rendering even 100 keV an attractive parameter regime.

For the interaction with multiple electrons, Eq. (5) readily generalizes to the product

$$\hat{U}_{N_{\text{el}}-\text{qb}} \approx e^i \sum_{k=1}^{N_{\text{el}}} \kappa(\hat{\mathbf{r}}_{\perp}^{(k)}, \boldsymbol{\alpha}) e^{\frac{i}{2} \hat{\sigma}_{\text{qb}}^x} \sum_{k=1}^{N_{\text{el}}} g(\hat{\mathbf{r}}_{\perp}^{(k)}, \boldsymbol{\alpha}). \quad (6)$$

This establishes the trapped ion as a quantum detector that is able to interact with multiple electrons in a coherent way. Such a detector may be used to generate many-electron entanglement as well as exotic ionic and electronic states in a two-path setup [88].

So far we have only discussed the interaction with a single ion. On a quantum computer based on surface-ion traps, one typically has many ions that can be arranged in linear or even quadratic arrays [89]. In a linear arrangement of the ions aligned with the beam axis, the coherent interaction with the ion chain results in a collective enhancement of the coupling strength, scaling with the number of ions. Arranging the ions perpendicular to the beam axis as visualized in Fig. 1(a), on the other hand, can give rise to many new measurement schemes. We envision a setup in which the electron is prepared in a superposition of distinct paths, where each of the paths is focused onto a single ion. This could be achieved by using electron diffraction gratings, as proposed in [39]. In each path, the electron would then interact with the corresponding ion according to Eq. (5), yielding the one-electron-many-qubit coupling operator:

$$\hat{U}_{e-N_{\text{qb}}} \approx e^{i\kappa} \sum_{k=1}^{N_{\text{qb}}} e^{\frac{i}{2} g \hat{\sigma}_{\text{qb},k}^x} |\psi_k\rangle \langle \psi_k|_{\text{el}}. \quad (7)$$

Here, $|\psi_k\rangle$ denotes the k th electron path and $\hat{\sigma}_{\text{qb},k}^x$ acts on the k th qubit. The coupling strength $g = \Delta\phi(0) - \Delta\phi(2\sqrt{2}R_0|\boldsymbol{\alpha}|)$ as well as the global phase κ is constant here as we assume the individual beams to be focused to the center of the corresponding coherent state $|\boldsymbol{\alpha}_k\rangle_{\text{cmi}}$ and we fix the magnitude of $\boldsymbol{\alpha}_k$. The initial state on the quantum computer can be a product state or a highly entangled state, depending on the metrological use case. The operator Eq. (7) can be understood as follows: if the electron is in the k th path, the k th qubit is subjected to a partial bit flip. By these means, the quantum computer's state becomes entangled with the electron path degree of freedom. If the electron subsequently traverses a weak-phase object, i.e., a biological specimen, the joint wave function acquires a slight phase shift. The electron-ion entanglement can now be used to transfer this phase shift to the qubits. This can be achieved by measuring the electron in momentum space and performing appropriate measurement-result-dependent operations on the quantum computer; see SM Sec. F [60].

Acquiring phase information about the specimen in this way might have many benefits. Since the coupling is coherent, phase shifts due to multiple probe electrons can add up [41], making it possible to surpass the standard quantum limit. In SM Sec. G [60] we explain briefly how the scheme proposed in [41] can be adapted to the ion-trap platform using only a single ion and discuss how to generalize it to allow for electron losses and nonideal coupling strengths. Furthermore, by coupling the electron to the trapped-ion array one can perform arbitrary quantum operations to extract as much information as possible out of a single electron. A similar advantage has been claimed in [90], where specially engineered diffractive elements are proposed as a means to perform optimal measurements, and in [42], where it has been proposed to implement Grover-like detection algorithms.

As a future avenue, integration into a multipass electron-microscopy setup might provide even more advantages. When the electron interacts only once with the ion and is then measured by the classical detector of the electron microscope, coherence can only be assured if the measurement result is used to compensate for the phase imprinted on the ion by the interaction. In a multipass setup, this requirement would drop out as the electron can interact many times with both sample and quantum computer without being detected. Furthermore, in the multipass electron microscopical setups currently envisioned [36,91], the probe electron has to be coupled out of the resonator in order to extract information. By using a trapped-ion quantum computer as a detector, a continuous coherent information transfer from the specimen to the quantum computer, mediated by a single or few electrons, could be possible. This would allow for weak or strong measurements without coupling the electron out of the resonator.

Our approach thus represents a significant step toward the integration of a well-established quantum-computing platform [92] into the powerful and versatile framework of electron microscopy.

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Data availability—The data that support the findings of this article are not publicly available. The data are available from the authors upon reasonable request.

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