

Quantum friction through monolayers

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(Received 22 December 2025; accepted 26 February 2026; published 13 April 2026)

Quantum friction is a friction force caused by electromagnetic fluctuations as an atom moves relative to macroscopic bodies. This effect is only predicted theoretically and requires high velocity and close proximity to give a significant signal. To date, experiments have not been able to measure this effect. Recent experiments with ultrafast atoms have sent neutral atoms through the crystal structure of monolayer graphene [Science **389**, 724 (2025)]. We present a theoretical framework describing quantum friction as atoms pass through monolayers and calculate the force on helium passing through graphene. Once the atom passes the monolayer, we find that the velocity-dependent force no longer opposes the direction of motion due to strong interactions caused by having been closer to the monolayer in the past. For the experimental parameters, we find a kinetic energy loss caused by quantum friction on the order of ten to hundred μeV . We find that quantum friction cannot be measured with this setup with current techniques.

DOI: [10.1103/787l-w86q](https://doi.org/10.1103/787l-w86q)

I. INTRODUCTION

The Casimir-Polder force is a force acting on atoms in the presence of macroscopic bodies [1–3]. In free space, the atom experiences changes to its energy levels caused by fluctuations in the form of emitted and reabsorbed virtual photons, known as the Lamb shift. In the presence of a macroscopic body, these fluctuations interact with it, creating a position-dependent Lamb shift for the atom, leading to an energy gradient and, consequently, a force between the atom and the body.

While the Casimir-Polder forces are well studied experimentally [4–9], theoretical calculations predict a velocity-dependent force when the atom is in relative motion to the object [10–15]. The early calculations focused on parallel motion relative to a surface. In this scenario, one expects a quantum friction force acting in the opposite direction of motion. The predicted force is very small and cannot be measured successfully.

Further calculations predict a much stronger quantum friction force when the atom moves perpendicular to the macroscopic body [16,17]. Due to the proportionality, the effect is strongly enhanced when the atom is moving very fast and in close proximity to the body. A further possibility, enhancing the effect, can be realized when the atom and the body are highly polarizable. As quantum friction depends on

all previous states and positions, a description of the entire motion is required to derive the actual effect [13,14].

Recent experiments on the diffraction of neutral atoms through the lattice structure of graphene have been performed [18]. These atoms have to move fast to overcome the short-range electric repulsion between the atom and the graphene monolayer [18,19] to tunnel through the monolayer. As atoms move very rapidly in close proximity to the macroscopic body, this scenario is a suitable candidate for the investigation of quantum friction.

In this work, we use a quantum master equation to theoretically model the quantum friction force experienced by an atom passing through a monolayer material. We present a model of internal atomic dynamics that is modified with respect to energy levels and state lifetimes.

After deriving the theory, we apply it to the specific dynamics of a helium atom passing through a graphene monolayer, where we find that after passing through a monolayer, the quantum friction force alternates between opposing and increasing the atom's velocity. The magnitude of the friction force in the helium-graphene case is found to be negligible compared to the other forces involved.

The quantum friction effect is qualitatively different on the two sides of the monolayer. On the escaping side, it is dominated by the residual interactions from when the atom was close to the monolayer and does not always oppose the direction of motion.

This paper aims to do the following:

- (1) Create a formalism for normal passage through a monolayer membrane, including the dependence on its history.
- (2) Introduce the sign-changing behavior of the quantum friction force, which happens after passing the monolayer.
- (3) Predict the energy loss to quantum friction force in current experiments, which we find to be on the μeV scale and thus undetectable.

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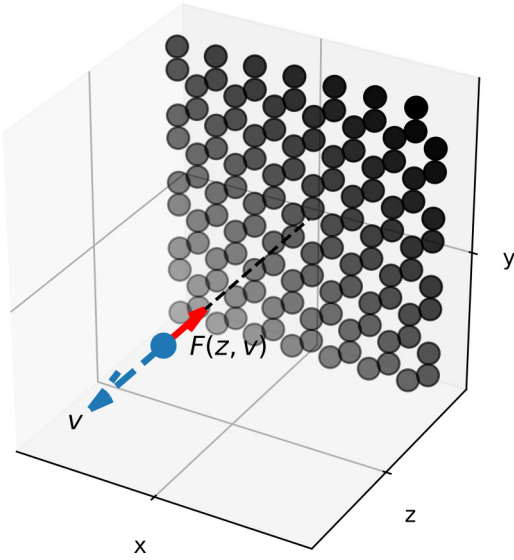


FIG. 1. Overview of the system being modeled. An atom (blue dot) travels with a velocity v (blue dashed arrow), through a monolayer experiencing the combined quantum friction and Casimir-Polder force $F(z, v)$ (orange arrow).

II. MODEL

Figure 1 shows the considered setup. An atom is accelerated in the z direction to a velocity v toward a monolayer lying in the xy plane. At high speeds, the atoms can tunnel through the monolayer and continue along their trajectories on the other side. The atom experiences the long-range Casimir-Polder force [2,3]. This force is modified when the atom is moving, leading to quantum friction [20]. This gives us a combined quantum friction and Casimir-Polder force $F(z, v)$, which is a function of position z and velocity v .

This one-dimensional approximation describes the mean dispersion force interaction between the atom and the graphene monolayer, in the regime where the kinetic energy of the atom is high, and the transverse velocity caused by diffraction can be considered an additional perturbation. At lower kinetic energies, nearer to the threshold where the atoms stop passing through the graphene, one has to account for the transverse velocity imparted to the atom. All kinetic energies considered here are well above this threshold.

A. Atomic dynamics

The Casimir-Polder and quantum friction follow from $F = -\nabla_z \hat{H}_I$, with $\hat{H}_I$ being the dipole interaction Hamiltonian. Since the dipole operator acts in the internal Hilbert space and carries no explicit spatial dependence, the gradient acts only on the electric field evaluated at the atomic position, yielding

$$F(t) = \langle \hat{\mathbf{d}}(t) \cdot \nabla \hat{\mathbf{E}}(\mathbf{r}, t) \rangle, \quad (1)$$

with the angle brackets indicating the expectation value of the atom and electromagnetic field states. Here, the dipole operator $\hat{\mathbf{d}}$ and the electric field operator $\hat{\mathbf{E}}(\mathbf{r})$ were evaluated at the particle's position $\mathbf{r}$. For constant velocity motion, $z = vt$ and we use $F(z, v)$ and $F(t)$ interchangeably.

At low velocity, this leads to the Casimir-Polder force. In contrast, the quantum friction force is the difference between this force and the one caused by the atom moving at a certain velocity relative to a macroscopic body that changes the dynamics of the electromagnetic field.

The Casimir-Polder and quantum friction forces arise from the atom's interaction with its induced field; thus, we first need to model how the atomic states evolve in an electromagnetic field. The atom's dynamics is given by

$$\dot{\hat{A}}_{mn}(t) = i\omega_{mn}\hat{A}_{mn}(t) + \frac{1}{i\hbar}[\hat{A}_{mn}(t), \hat{H}_I(t)], \quad (2)$$

with $\hat{A}_{mn} = |m\rangle\langle n|$ being the flip operators describing transitions from atomic energy state m to n . $\hat{H}_I = -\hat{\mathbf{d}} \cdot \hat{\mathbf{E}}(\mathbf{r})$ is the interaction Hamiltonian, describing the coupling between the atom and the electric field. This differential equation can be solved iteratively by resubstituting the operators into itself, leading to the next order [16]

$$\dot{\hat{A}}_{mn}^{(2)}(t) = \dot{\hat{A}}_{mn}^{(0)}(t) + \frac{1}{i\hbar}[\hat{A}_{mn}^{(0)}(t), \hat{H}_I^{(1)}(t)], \quad (3)$$

with the higher-order interaction Hamiltonian [16]

$$\hat{H}_I^{(1)}(t) = -\sum_{mn} \hat{A}_{mn}^{(0)}(t) \hat{\mathbf{d}}_{mn} \cdot \hat{\mathbf{E}}^{(1)}(\mathbf{r}, t). \quad (4)$$

The operator $\hat{\mathbf{d}}_{mn}$ is the atomic dipole operator in the flip operator basis. The operator $\hat{\mathbf{E}}^{(1)}$ is the electric field induced by the atomic transitions $\hat{A}_{mn}^{(0)}(\tau)$:

$$\begin{aligned} \hat{\mathbf{E}}^{(1)}(\mathbf{r}, t) &= \frac{i\mu_0}{\pi} \sum_{mn} \int_{t_0}^t d\tau \int_0^\infty d\omega \omega^2 e^{i(\omega - \omega_{mn})(t - \tau)} \\ &\quad \times \text{Im } \mathbf{G}(\mathbf{r}, \mathbf{r}', \omega) \cdot \hat{\mathbf{d}}_{mn} \hat{A}_{mn}^{(0)}(\tau), \end{aligned} \quad (5)$$

with $\mathbf{G}$ being the electromagnetic Green tensor. $\mathbf{r}$ and $\mathbf{r}' = \mathbf{r} - \mathbf{v}\tau$ are the positions of the atoms at time t and $t - \tau$, respectively. μ_0 is the vacuum permeability. This represents a photon being emitted at time τ . Next, we substitute Eqs. (4) and (5) into Eq. (3). Doing so, one must sum over all intermediate transitions $\hat{A}_{nk}$ and $\hat{A}_{km}$, keeping in mind that $\hat{A}_{mk}\hat{A}_{kn} = \hat{A}_{mn}$. The resulting time evolution, for the operator averages, takes the form [16]

$$\langle \dot{\hat{A}}_{mn}^{(2)}(t) \rangle = i\omega_{mn} \langle \hat{A}_{mn}^{(2)}(t) \rangle - \sum_k [C_{nk}(t) + C_{mk}^*(t)] \langle \hat{A}_{mn}^{(2)}(t) \rangle, \quad (6)$$

with the coefficients $C_{nk}(z, v, t)$ being given by [16,17]

$$\begin{aligned} C_{nk}(z, v, t) &= \frac{\mu_0}{\pi\hbar} \int_{t_0}^t d\tau \int_0^\infty d\omega \omega^2 e^{-i(\omega - \omega_{nk})(t - \tau)} \\ &\quad \times \hat{\mathbf{d}}_{nk} \cdot \text{Im } \mathbf{G}(\mathbf{r}, \mathbf{r}', \omega) \cdot \hat{\mathbf{d}}_{kn}. \end{aligned} \quad (7)$$

The nonretarded Green's tensor at a planar interface reads [2]

$$\mathbf{G}(\mathbf{r}, \mathbf{r}', \omega) = \int_0^\infty \frac{d^2 k_\parallel}{k_\parallel} \frac{r(k_\parallel, \omega)c^2}{8\pi^2\omega^2} (\mathbf{k}_\parallel \otimes \mathbf{k}_\parallel^*) e^{-k_\parallel(|z| + |z - vt|)}, \quad (8)$$

where $k_\parallel$ is the wave vector in the plane parallel to the monolayer. Substituting this Green's function into Eq. (7) and

performing the integrals, one obtains

$$C_{nk}(z, v, t) = -\frac{1}{4\pi^2\epsilon_0\hbar} \int_0^\infty dk_\parallel \int_0^\infty d\omega \hat{\mathbf{d}}_{nk} \cdot \frac{r(k_\parallel, \omega)(\mathbf{k}_\parallel \otimes \mathbf{k}_\parallel^*)}{k_\parallel v - i(\omega - \omega_{nk})} e^{-2k_\parallel(|z|)} \cdot \hat{\mathbf{d}}_{kn} \\ + \frac{1}{4\pi^2\epsilon_0\hbar} \int_0^\infty dk_\parallel \int_0^\infty d\omega \hat{\mathbf{d}}_{nk} \cdot \frac{r(k_\parallel, \omega)(\mathbf{k}_\parallel \otimes \mathbf{k}_\parallel^*)}{k_\parallel v - i(\omega - \omega_{nk})} e^{-k_\parallel(|z|+|z-v(t-t_0)|)-i(\omega-\omega_{nk})(t-t_0)} \cdot \hat{\mathbf{d}}_{kn}. \quad (9)$$

Equation (9) represents the finite initial time solution. Taking the limit $t_0 \rightarrow -\infty$ removes the transient term and yields Eq. (11).

The imaginary and real parts of the coefficients C_{nk} describe the energy change $\delta\omega$ and rate of spontaneous decay of the atom

$$i\delta\omega_n + \frac{1}{2}\Gamma_n = \sum_k C_{nk}. \quad (10)$$

For an approaching atom, we can set the initial time to $-\infty$, and thus the second term of Eq. (9) vanishes, leaving

$$C_{nk}(z, v) = -\frac{1}{4\pi^2\epsilon_0\hbar} \int_0^\infty dk_\parallel \int_0^\infty d\omega \hat{\mathbf{d}}_{nk} \cdot \frac{r(k_\parallel, \omega)(\mathbf{k}_\parallel \otimes \mathbf{k}_\parallel^*)}{k_\parallel v - i(\omega - \omega_{nk})} e^{-2k_\parallel(|z|)} \cdot \hat{\mathbf{d}}_{kn}. \quad (11)$$

Once the atom passes the monolayer, the absolute value in the exponential of Eq. (9) changes the sign of v , which means we must perform two time integrals, one from $-\infty$ to $t - z/v$ and one from $t - z/v$ to t , thus leading to a second term

$$C_{nk}(z, v) = -\frac{1}{4\pi^2\epsilon_0\hbar} \int_0^\infty dk_\parallel \int_0^\infty d\omega \hat{\mathbf{d}}_{nk} \cdot \frac{r(k_\parallel, \omega)(\mathbf{k}_\parallel \otimes \mathbf{k}_\parallel^*)}{k_\parallel v - i(\omega - \omega_{nk})} e^{-2k_\parallel(|z|)} \cdot \hat{\mathbf{d}}_{kn} \\ + \frac{1}{2\pi^2\epsilon_0\hbar} \int_0^\infty dk_\parallel \int_0^\infty d\omega \hat{\mathbf{d}}_{nk} \cdot \frac{r(k_\parallel, \omega)(\mathbf{k}_\parallel \otimes \mathbf{k}_\parallel^*)k_\parallel v}{k_\parallel^2 v^2 + (\omega - \omega_{nk})^2} e^{-k_\parallel(|z|)-i(\omega-\omega_{nk})z/v} \cdot \hat{\mathbf{d}}_{kn}. \quad (12)$$

B. Forces

The force experienced by the atom is given by

$$F(t) = \text{Re} \frac{i\mu_0}{\pi} \int_0^\infty d\tau \int_0^\infty dk_\parallel \int_0^\infty d\omega \nabla \langle \omega^2 \hat{\mathbf{d}}(t) \cdot \text{Im} \mathbf{G}(\mathbf{r}(t), \mathbf{r}(t-\tau), \omega) \cdot \hat{\mathbf{d}}(t-\tau) e^{i\omega\tau} \rangle. \quad (13)$$

Using

$$\langle \hat{\mathbf{d}}(t) \cdot \hat{\mathbf{d}}(t-\tau) \rangle = \sum_{nk} \hat{\mathbf{d}}_{nk} \hat{\mathbf{d}}_{kn} p_n(t) e^{i\omega'_{nk}\tau - \frac{1}{2}(\Gamma'_n + \Gamma'_k)\tau}, \quad (14)$$

where the prime signifies that the effects of the modified atomic dynamics have to be added, the force separates into the contribution from each transition,

$$F(t) = \sum_{nk} F_{nk}(t). \quad (15)$$

Here, we introduced the population of the state $p_n(t) = \langle n | \hat{\rho}_A(t) | n \rangle$.

For an atom that remains in state n , and is approaching the monolayer, the time integral in Eq. (13) can be performed from 0 to ∞ , which yields

$$F_{nk}(t) = \text{Re} \frac{ip_n(t)|d_{nk}|^2}{2\pi^2\epsilon_0} \int_0^\infty dk_\parallel \int_0^\infty d\omega \frac{\text{Sign}(z)\text{Im}[r(k_\parallel, \omega)]k_\parallel^3}{k_\parallel v - i(\omega - \omega'_{nk}) + \frac{1}{2}(\Gamma'_n + \Gamma'_k)} e^{-2k_\parallel|z|}, \quad (16)$$

as discussed in the literature [16]. For a departing atom, one has to carry out a time integral from 0 to z/v , and another time integral from z/v to ∞ , which leads to

$$F_{nk}(t) = \text{Re} \left\{ \frac{ip_n(t)|d_{nk}|^2}{2\pi^2\epsilon_0} \int_0^\infty dk_\parallel \int_0^\infty d\omega \frac{\text{Sign}(z)\text{Im}[r(k_\parallel, \omega)]k_\parallel^3}{k_\parallel v - i(\omega - \omega'_{nk}) + \frac{1}{2}(\Gamma'_n + \Gamma'_k)} e^{-2k_\parallel|z|} \right\} \\ - \text{Re} \left\{ \frac{ip_n(t)}{2\pi^2\epsilon_0} \int_0^\infty dk_\parallel \int_0^\infty d\omega \frac{\text{Sign}(z)\text{Im}[r(k_\parallel, \omega)]k_\parallel^3}{-k_\parallel v - i(\omega - \omega'_{nk}) + \frac{1}{2}(\Gamma'_n + \Gamma'_k)} e^{-k_\parallel|z|+i(\omega-\omega'_{nk})z/v - \frac{1}{2}(\Gamma'_n + \Gamma'_k)z/v} \right\}, \quad (17)$$

where the additional second term leads to oscillations of the quantum friction.

III. HELIUM THROUGH GRAPHENE

We will look into the specific case of helium going through graphene. We consider ground-state helium, and thus set $p_0(t) = 1$. Furthermore, to simplify the calculations, the atom is assumed to be a two-level system with a ground and a single excited 2^1P state. The energy difference between the two levels is $E_A = 3.40 \times 10^{-18} \text{ J} = 21.2 \text{ eV}$ [21], and the spontaneous decay rate is $1.80 \times 10^{-7} \text{ s}^{-1}$ [21]. We define the atomic angular frequency $\omega_A = E_A/\hbar$.

For the reflection coefficient of graphene, we use the $(2+1)$ -dimensional Dirac model, as it is described in Ref. [22]. This model is sufficient for low-energy excitations below about 2 eV [22,23]. As the dispersion interaction with the atom is primarily dominated by lower-energy contributions, this should be sufficient [23]. More complete calculations of the reflection coefficient, such as full inter- and intraband calculations or finite-temperature calculations, are therefore not required. In this model, the reflection coefficient of graphene is given by

$$r(k_{\parallel}, \omega) = \frac{4\pi\alpha\sqrt{\omega^2/c^2 - k_{\parallel}^2}}{4\pi\alpha\sqrt{\omega^2/c^2 - k_{\parallel}^2} + 8\sqrt{\omega^2/c^2 - k_{\parallel}^2}v_F^2}, \quad (18)$$

with the fine structure constant $\alpha = \frac{1}{137}$ and the Fermi velocity $v_F = 1/300$, which we use as the dimensionless ratio between the Fermi velocity and the speed of light [22]. For distances below 1 nm, the relevant $k_{\parallel}$ satisfy $\omega/c \ll k_{\parallel}$, justifying the nonretarded approximation, since the dominant contribution to the $k_{\parallel}$ integral comes from $k_{\parallel} \propto 1/|z|$. Thus, the reflection coefficient becomes

$$r(k_{\parallel}, \omega) = \frac{1}{1 + \frac{2}{\pi\alpha}\sqrt{v_F^2 - \frac{\omega^2}{c^2k_{\parallel}^2}}}. \quad (19)$$

We can then find the force by substituting this into Eqs. (16) and (17). Once we get within a few angstroms of the monolayer, the Green's tensor approach is no longer valid. This regime is dominated by short-range repulsion forces. The forces in this regime have been treated using density-functional theory [18,19] or atomistic potentials. We exclude these forces from our fluctuational-electrodynamics calculations of the Casimir-Polder and quantum friction force. Within the van der Waals radius of the carbon (1.7 Å) in the monolayer, we define a thickness of the monolayer, where we make it an interaction-free tunneling of the atom, as our Green's tensor approach becomes invalid. This prescription does not affect the asymptotic energy loss because the quantum friction contribution is negligible compared to short-range forces in this region.

For the considered weak interactions, there is no significant change in the force from the modified internal dynamics. Using this fact and substituting Eq. (19) into Eq. (16), and using the fact that the ground-state helium has a negligible decay rate, simplifies the approaching part to

$$F(z, v) = \text{Re} \frac{i|d_{01}|^2}{2\pi^2\epsilon_0} \int_0^\infty dk_{\parallel} \int_0^\infty d\omega \times \frac{\text{Sign}(z)\text{Im}[r(k_{\parallel}, \omega)]k_{\parallel}^3}{k_{\parallel}v - i(\omega + \omega_A)} e^{-2k_{\parallel}|z|}. \quad (20)$$

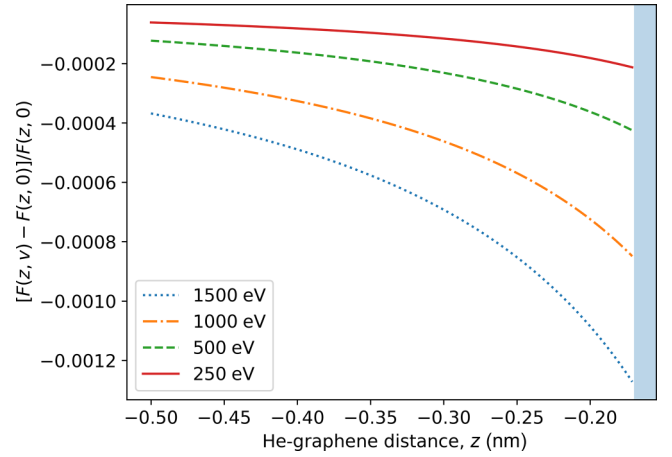


FIG. 2. The quantum friction relative to the Casimir-Polder force as a function of distance for an approaching atom. The plots show atoms with kinetic energies of 250, 500, 1000, and 1500 eV. The blue region represents the inside of the monolayer; it is $2r_{\text{vdW}}$ wide. We assume there is no friction force within this area.

Substituting into Eq. (17) yields a departing part of the quantum friction, which reads

$$F(z, v) = \text{Re} \frac{i|d_{01}|^2}{2\pi^2\epsilon_0} \int_0^\infty dk_{\parallel} \int_0^\infty d\omega \times \frac{\text{Sign}(z)\text{Im}[r(k_{\parallel}, \omega)]k_{\parallel}^3}{k_{\parallel}v - i(\omega + \omega_A)} e^{-2k_{\parallel}|z|} - \text{Re} \frac{i}{2\pi^2\epsilon_0} \int_0^\infty dk_{\parallel} \int_0^\infty d\omega \times \frac{\text{Sign}(z)\text{Im}[r(k_{\parallel}, \omega)]k_{\parallel}^3}{-k_{\parallel}v - i(\omega + \omega_A)} e^{-k_{\parallel}|z| + i(\omega + \omega'_A)z/v}. \quad (21)$$

IV. RESULTS

We now evaluate the force expressions (16) and (17) numerically for helium passing graphene under experimental conditions of Ref. [18]. Helium has been sent through graphene at kinetic energies between 390 and 1600 eV [18]. Theoretical calculations predict a minimal energy loss to the monolayer membrane at 250 eV. In Figs. 2 and 3, we see the strength of the quantum friction relative to the Casimir-Polder potential plotted against the distance between the helium and the monolayer. The y axis is given by the formula

$$\tilde{F}_{\text{rel}} = \frac{F(z, v) - F(z, 0)}{F(z, 0)}. \quad (22)$$

The Casimir-Polder force is shown in Fig. 4. The atom is approaching from the left. Before passing through the monolayer, the quantum friction force opposes the direction of motion and increases in magnitude with increasing velocity.

Once we are within the van der Waals radius, r_{vdW} , of the carbon atoms in the monolayer (1.7 Å), short-range repulsion between the atoms becomes dominant. We set the quantum friction force to 0 in this region.

The departing atom experiences an equivalent friction force on the other side of the monolayer. However, the

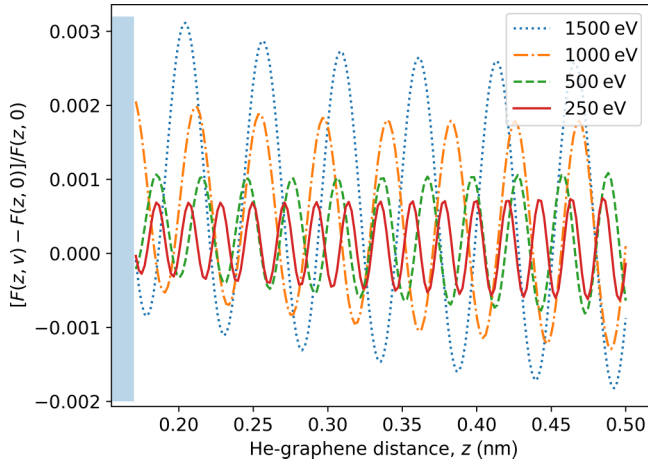


FIG. 3. The quantum friction relative to the Casimir-Polder force as a function of distance for an escaping atom. The plots show atoms with kinetic energies of 250, 500, 1000, and 1500 eV. The blue region represents the inside of the monolayer; it is $2r_{\text{vdW}}$ wide. We assume there is no friction force within this area.

contributions from the past trajectory before passing the layer give rise to a force that rapidly oscillates between accelerating and decelerating the atom. The magnitude of these oscillations increases with velocity.

Figures 5 and 6 show the relative quantum friction force plotted against the unitless quantity $\frac{v}{z\omega_A}$ for an approaching and a departing atom, respectively. For the approaching case, this seems to correspond to earlier work that predicts a v^2 dependence of the quantum friction force at low velocities [16]. Unlike others, we do not find a function of $\frac{v}{\omega_A z}$. This is probably due to the $k_{\parallel}$ dependence in the reflection coefficient.

At very high velocities, the atom will have been far away from the monolayer in the recent past; thus, it experiences a very weak force [17].

In the departing case, the direction of the friction force oscillates due to the imaginary term in the exponential. At higher

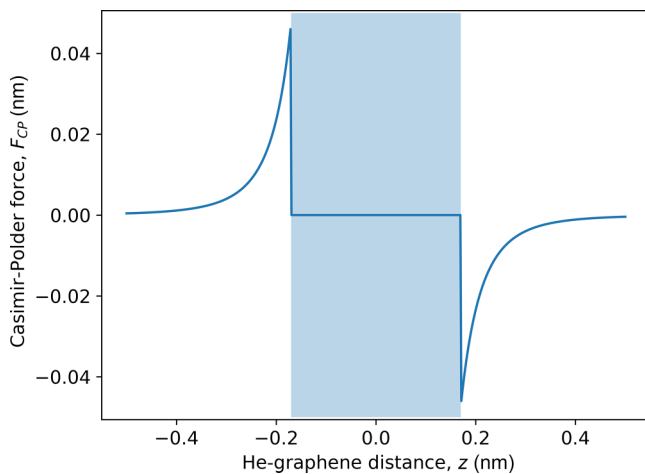


FIG. 4. Casimir-Polder force plotted against the atom's distance from the monolayer. The blue region represents the inside of the monolayer; it is $2r_{\text{vdW}}$ wide. We assume the Casimir-Polder force to disappear within this region.

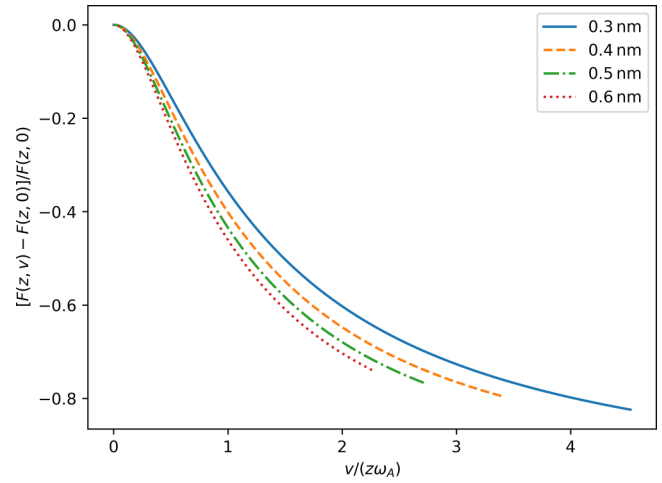


FIG. 5. Dependence of relative quantum friction force [Eq. (22)] on the unitless quantity $\frac{v}{z\omega_A}$, for $z = 0.3, 0.4, 0.5$, and 0.6 nm, for an approaching atom.

velocities, the oscillations happen over longer distances. As the atom was closer in the past, these forces can be stronger than the Casimir-Polder force. This is not possible otherwise. Then, it will at most be equal in magnitude, canceling out the Casimir-Polder force [17].

Because the quantum friction force depends on velocity, it is not conservative and can change the atom's kinetic energy as it passes the monolayer. Since the kinetic energy required to go through the monolayer is much greater than the potential energy loss, it is possible to calculate the energy loss by assuming that the atom's velocity remains constant. Integrating the force over the atom's trajectory gives us the change in its kinetic energy.

Figure 7 shows the energy difference compared to the initial kinetic energy of the helium atom. We see that it oscillates smoothly between positive and negative energy changes. The

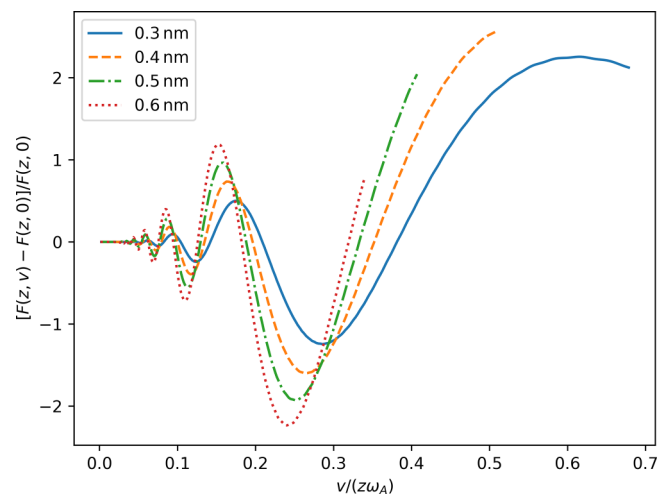


FIG. 6. Dependence of relative quantum friction force [Eq. (22)] on the unitless quantity $\frac{v}{z\omega_A}$, for $z = 0.3, 0.4, 0.5$, and 0.6 nm, for a departing atom.

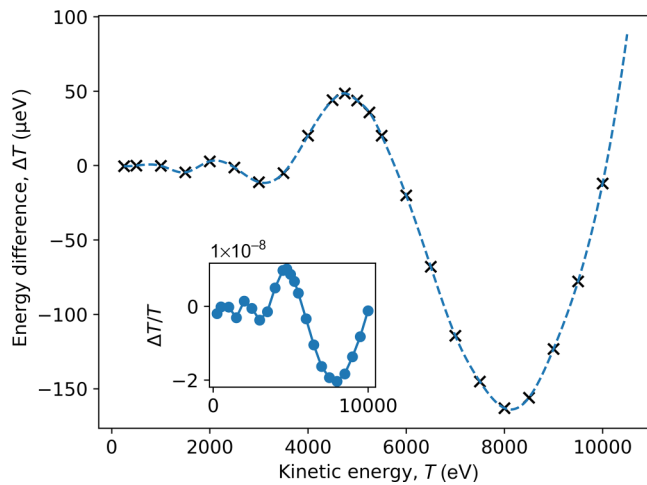


FIG. 7. Change in kinetic energy difference, ΔT , of a helium atom, plotted against initial kinetic energy, T , after it passes a graphene monolayer. Black crosses denote calculations, and the blue line is for guidance. The inset shows the relative energy change $\Delta T/T$ over the same range of T .

magnitude of the energy change grows with greater kinetic energies.

V. CONCLUSIONS

We have presented a model and calculations of the velocity-dependent changes to the Casimir-Polder forces between atoms and monolayers as the atoms pass the monolayer at high velocities. The case for helium and graphene is calculated at different velocities, and the energy change due to the velocity-dependent force is found.

We found that the strong past interactions with the monolayer create a force that changes between decelerating and accelerating the atom in the direction of motion. It changes between these in regular intervals. This means the force is no longer simply a friction force. This oscillating force is several times stronger than the quantum friction force on the approaching side.

For energies tested at current experiments, we see that the quantum friction force is negligible compared to the Casimir-Polder force, which is almost negligible itself. This means that measurements of quantum friction in experiments with atoms passing monolayers are unlikely to be possible.

Beyond this work, further studies should focus on identifying optimal combinations of velocity, monolayer type, and atomic species to enhance quantum friction forces. In addition, further studies should investigate whether excited atoms can be transmitted through a monolayer, as these exhibit much stronger dispersion interactions. However, it is not clear whether it is possible for an excited atom to pass a monolayer while staying excited.

Further work on quantum friction in atoms passing through holes in monolayers should also be done. For very small holes, we could retain the effects of past interactions without the very strong repulsive forces inside the monolayer. This could potentially lead to a more sensitive setup.

DATA AVAILABILITY

No new datasets were generated or analysed during this study. All equations, parameters, and methodological details required to reproduce the results are provided in the manuscript.

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