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Path decoherence of charged and neutral particles near surfaces

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We show how the spatial decoherence properties of coherently split beams of charged and neutral particles are influenced by their (quantum) electromagnetic interaction with a metallic surface. Our investigation makes use of macroscopic quantum electrodynamics based on linear-response theory, which provides an alternative to microscopic models of the decoherence mechanism. We find strikingly simple expressions for the decoherence rate that admit an interpretation in terms of *which path* information due to image-charge interactions.

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Introduction. According to the superposition principle of quantum mechanics—itself a straightforward consequence of the linearity of the Schrödinger equation—linear superpositions of solutions of the Schrödinger equation are also solutions themselves. However, in nature these superpositions of quantum states are hard, if not impossible, to observe. Interactions with an unobserved environment destroy any coherent superpositions and leave a quantum system in a mixed state with no (or strongly reduced) quantum coherence. Understanding decoherence is important in the study of the transition from quantum to classical dynamics in quantum systems interacting with an environment [1]. A variety of experiments have been designed to push the boundaries of observation of quantum decoherence effects, most notably in molecule interferometry [2], cavity quantum electrodynamics (QED) [3], and ground-state cooling of nanomechanical oscillators [4].

A particular environment, ubiquitous to many experiments in atomic physics, consists of the modes of the electromagnetic field that are modified by the presence of macroscopic bodies such as mirrors or other dielectric or metallic structures. The exchange of photons between coherently split electron beams is a source of decoherence already in free space [5–7] and sets upper limits to the distance or time the electrons can travel coherently. However, experiments with coherently split electron beams close to a metallic surface have shown much shorter coherence lengths [8]. These were found to be caused by the field fluctuations emanating from the metal surface [9] attributed to the dissipative environment obtaining *which path* information [10].

In this Rapid Communication, we study the path decoherence of charged and neutral particles near metallic surfaces from the viewpoint of macroscopic quantum electrodynamics without resorting to microscopic descriptions of the environment. In both instances we concentrate on the long-time (Markovian) dynamics in the slow-velocity limit. We obtain intuitive results that can be interpreted as the interaction of the particles with their mirror images, thus affirming the interpretation of *which path* dephasing.

Path decoherence of electron beams. The starting point is the decoherence functional theory described in Ref. [11], Chap. 12. As shown therein, the time evolution of the matrix elements of the reduced density matrix $\hat{\varrho}(t)$ of a charged particle moving along two paths 1 and 2 from an initial time t_i to a final time t_f can be approximated by

$$\hat{\varrho}(t_f) \simeq \hat{\varrho}_{11}(t_i) + \hat{\varrho}_{22}(t_i) + e^{i\Phi}\hat{\varrho}_{12}(t_i) + e^{-i\Phi^*}\hat{\varrho}_{21}(t_i). \quad (1)$$

Here, Φ is the influence phase functional that arises from the interaction with the (medium-assisted) electromagnetic field. The decoherence properties, i.e., the degrading of the off-diagonal density matrix elements $\hat{\varrho}_{ij}(t_i)$, are contained in the imaginary part of this phase $|e^{i\Phi}| = e^{\Gamma[c]}$, where the functional $\Gamma[c]$ depends only on the classical current density difference $c^\mu(x) = [j_1^\mu(x) - j_2^\mu(x)]/\sqrt{2}$ along the two paths 1 and 2. It can be written as

$$\Gamma[c] = -\frac{1}{2\hbar^2}\int_{t_i}^{t_f} d^4x \int_{t_i}^{t_f} d^4x'\, D_+(x,x')_{\mu\nu} c^\mu(x) c^\nu(x') \quad (2)$$

in terms of the field anticommutator function $D_+(x,x')_{\mu\nu} = \langle\{\hat{A}_\mu(x), \hat{A}_\nu(x')\}\rangle$.

We focus first on the spatial components of the anticommutator function. They are easily obtained by writing out the vector potential in terms of its frequency components as $\hat{A}_i(\mathbf{r},t) = \int_0^\infty d\omega\, \hat{A}_i(\mathbf{r},\omega)e^{-i\omega t} + \text{H.c.}$ with (thermal) expectation values ($\perp$: transverse part) [12]

$$\langle \hat{A}_i^\dagger(\mathbf{r},\omega)\hat{A}_j(\mathbf{r}',\omega')\rangle = \frac{\hbar\mu_0}{\pi}\bar{n}_{\text{th}}(\omega)\text{Im}\, G_{ij}^{\perp\perp}(\mathbf{r},\mathbf{r}',\omega)\delta(\omega-\omega'), \quad (3)$$

$$\langle \hat{A}_i(\mathbf{r},\omega)\hat{A}_j^\dagger(\mathbf{r}',\omega')\rangle = \frac{\hbar\mu_0}{\pi}[\bar{n}_{\text{th}}(\omega)+1]\text{Im}\, G_{ij}^{\perp\perp}(\mathbf{r},\mathbf{r}',\omega) \times \delta(\omega-\omega'), \quad (4)$$

from which follows that $[\boldsymbol{G}(\mathbf{r},\mathbf{r}',\omega) = \boldsymbol{G}^T(\mathbf{r}',\mathbf{r},\omega)]$

$$D_+^\perp(x,x')_{ij} = \frac{\hbar\mu_0}{\pi}\int_0^\infty d\omega[2\bar{n}_{\text{th}}(\omega)+1]\text{Im}\, G_{ij}^{\perp\perp}(\mathbf{r},\mathbf{r}',\omega) \times [e^{i\omega(t-t')} + e^{-i\omega(t-t')}]. \quad (5)$$

The mean thermal photon number $\bar{n}_{\text{th}}(\omega)$ obeys the Bose-Einstein distribution, $\bar{n}_{\text{th}}(\omega) = [e^{\hbar\omega/(k_BT)} - 1]^{-1}$.

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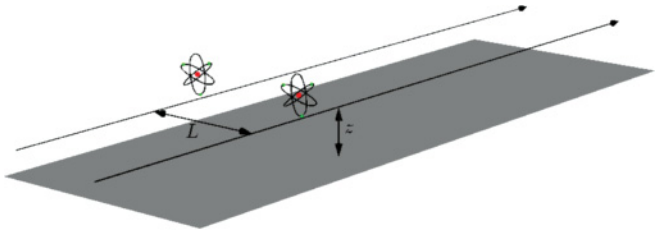

FIG. 1. (Color online) Schematic setup of two electron beams traveling along the x direction with velocity v close to a dielectric or metallic surface. The two beams are separated by a lateral distance L and travel at a distance z from the surface.

For a single charged particle (e.g., an electron) traveling along a path $\mathbf{r}_e(t)$, the classical current density is simply $\mathbf{j}_e(x) = -e\mathbf{v}\delta[\mathbf{r} - \mathbf{r}_e(t)]$ so that the spatial part $\Gamma_{(ij)}[c]$ of the decoherence functional (2) becomes [$\mathbf{r}_i \equiv \mathbf{r}_i(t)$, $\mathbf{r}'_i \equiv \mathbf{r}_i(t')$]

$$\Gamma_{(ij)}[c] = -\frac{e^2\mu_0}{4\hbar\pi}\int_0^\infty d\omega \int_{t_i}^{t_f} d^4x \int_{t_i}^{t_f} d^4x' [2\bar{n}_{\rm th}(\omega)+1] \times e^{i\omega(t-t')}\mathbf{v}\cdot\{\mathrm{Im}\,\boldsymbol{G}^{\perp\perp}(\mathbf{r}_1,\mathbf{r}'_1,\omega) + \mathrm{Im}\,\boldsymbol{G}^{\perp\perp}(\mathbf{r}_2,\mathbf{r}'_2,\omega) - \mathrm{Im}\,\boldsymbol{G}^{\perp\perp}(\mathbf{r}_1,\mathbf{r}'_2,\omega) - \mathrm{Im}\,\boldsymbol{G}^{\perp\perp}(\mathbf{r}_2,\mathbf{r}'_1,\omega)\}\cdot\mathbf{v} + \text{c.c.} \tag{6}$$

For simplicity, we have assumed that both electron beams travel in parallel with the same velocity $\mathbf{v}$. The temporal and mixed spatio-temporal contributions vanish identically for such motion when also being parallel to the surface, so that $\Gamma[c] \equiv \Gamma_{(ij)}[c]$. In principle, these contributions exist for divergent and convergent beams and lead to bremsstrahlung [11].

At this point it is impossible to simplify the expressions further without referring to a specific geometry. As an example, we envisage a situation in which a wave packet is initially split, and its two components then propagate parallel to a macroscopic body before being recombined. We are interested in the decoherence during the parallel transport as depicted in Fig. 1. Decomposing the Green's function into a free-space part $\boldsymbol{G}^{(0)}$ and a scattering part $\boldsymbol{G}^{(S)}$, we note that the (double-sided transverse) scattering part dominates at small particle-surface distances. For a particle traveling along the x direction parallel to a plane dielectric or metallic surface, we use the Weyl expansion to write the scattering part of the relevant Green's tensor component as [12] ($\mathbf{r}_i = \boldsymbol{\rho}_i + z_i\mathbf{e}_z$, $\mathbf{k} = \mathbf{k}_\parallel + k_z\mathbf{e}_z$)

$$G^{(S)}_{xx}(\mathbf{r}_1,\mathbf{r}_2,\omega) = \int \frac{d^2k_\parallel}{(2\pi)^2} e^{i\mathbf{k}_\parallel\cdot(\boldsymbol{\rho}_1-\boldsymbol{\rho}_2)} R_{xx}(\mathbf{k}_\parallel,z,\omega), \tag{7}$$

$$R_{xx}(\mathbf{k}_\parallel,z,\omega) = \frac{i}{2k_z}e^{2ik_z z}\left[\frac{r_p c^2}{\omega^2}\left(-k_z^2\frac{k_x^2}{k_\parallel^2}\right) + r_s\frac{k_y^2}{k_\parallel^2}\right], \tag{8}$$

where $r_{s,p}(k_\parallel,\omega)$ are the Fresnel reflection coefficients for (s,p)-polarized waves. Here we have also assumed that both electron beams travel across the surface at the same height z. Note that the spatial arguments in the Green's tensor refer to the time-dependent trajectories of the two electron beams, $\mathbf{r}_{1,2}(t) = \mathbf{r}_{1,2} + \mathbf{v}t$. Hence, in the Weyl expansion there are terms of the form $e^{i\mathbf{k}_\parallel\cdot[\boldsymbol{\rho}_i(t)-\boldsymbol{\rho}_j(t)]} = e^{i\mathbf{k}_\parallel\cdot(\boldsymbol{\rho}_i-\boldsymbol{\rho}_j)}e^{i\mathbf{k}_\parallel\cdot\mathbf{v}(t-t')}$, implying that spatio-temporal dependencies decouple.

The time integrals in the decoherence functional can then be performed to give

$$\int_{t_i}^{t_f} dt \int_{t_i}^{t_f} dt'\, e^{i(\mathbf{k}_\parallel\cdot\mathbf{v}\pm\omega)(t-t')} = \frac{2}{(\mathbf{k}_\parallel\cdot\mathbf{v}\pm\omega)^2}[1-\cos(\mathbf{k}_\parallel\cdot\mathbf{v}\pm\omega)(t_f-t_i)]. \tag{9}$$

The remaining exponential functions containing only the spatial variables combine to

$$e^{i\mathbf{k}_\parallel\cdot(\boldsymbol{\rho}_1-\boldsymbol{\rho}_1)} + e^{i\mathbf{k}_\parallel\cdot(\boldsymbol{\rho}_2-\boldsymbol{\rho}_2)} - e^{i\mathbf{k}_\parallel\cdot(\boldsymbol{\rho}_1-\boldsymbol{\rho}_2)} - e^{i\mathbf{k}_\parallel\cdot(\boldsymbol{\rho}_2-\boldsymbol{\rho}_1)} = 2[1-\cos\mathbf{k}_\parallel\cdot(\boldsymbol{\rho}_1-\boldsymbol{\rho}_2)]. \tag{10}$$

With these results, the decoherence functional reads

$$\Gamma[c] = -\frac{e^2v^2\mu_0}{\hbar\pi}\int_0^\infty d\omega\,\mathrm{Im}\int\frac{d^2k_\parallel}{(2\pi)^2}[1-\cos\mathbf{k}_\parallel\cdot(\boldsymbol{\rho}_1-\boldsymbol{\rho}_2)] \times [2\bar{n}_{\rm th}(\omega)+1]\frac{i}{2k_z}e^{2ik_z z}\left[\frac{r_p c^2}{\omega^2}\left(-k_z^2\frac{k_x^2}{k_\parallel^2}\right) + r_s\frac{k_y^2}{k_\parallel^2}\right] \times\left[\frac{1-\cos(\mathbf{k}_\parallel\cdot\mathbf{v}+\omega)t}{(\mathbf{k}_\parallel\cdot\mathbf{v}+\omega)^2} + \frac{1-\cos(\mathbf{k}_\parallel\cdot\mathbf{v}-\omega)t}{(\mathbf{k}_\parallel\cdot\mathbf{v}-\omega)^2}\right], \tag{11}$$

where we have set $t \equiv t_f - t_i$.

As it stands, Eq. (11) is still not analytically tractable. Hence, some more approximations are in order. Insisting that the decoherence functional should be linear in the interaction time, we have to employ the long-time limit. This is equivalent to performing a Markov approximation which can be done if we assume that the surface response function is not sharply peaked. In this limit, one approximates

$$\frac{[1-\cos(\mathbf{k}_\parallel\cdot\mathbf{v}\pm\omega)t]}{(\mathbf{k}_\parallel\cdot\mathbf{v}\pm\omega)^2} \mapsto \pi t\,\delta(\mathbf{k}_\parallel\cdot\mathbf{v}\pm\omega). \tag{12}$$

The frequency integration can then be performed. Note that the integral only runs over positive frequencies whereas the projection of the wave vector onto the direction of the velocity can be both positive or negative. Hence, we have to write

$$\int_0^\infty d\omega[\delta(\mathbf{k}_\parallel\cdot\mathbf{v}+\omega)+\delta(\mathbf{k}_\parallel\cdot\mathbf{v}-\omega)]f(\omega) = f(|\mathbf{k}_\parallel\cdot\mathbf{v}|). \tag{13}$$

In this way we obtain

$$\Gamma[c] = -\frac{\mu_0 q^2 v^2 t}{\hbar}\int\frac{d^2k_\parallel}{(2\pi)^2}[1-\cos\mathbf{k}_\parallel\cdot(\boldsymbol{\rho}_1-\boldsymbol{\rho}_2)] \times[2\bar{n}_{\rm th}(|\mathbf{k}_\parallel\cdot\mathbf{v}|)+1]\mathrm{Im}\,R_{xx}(\mathbf{k}_\parallel,z,|\mathbf{k}_\parallel\cdot\mathbf{v}|). \tag{14}$$

Combining all of the above yields an expression for the decoherence functional in the long-time limit as

$$\Gamma[c] = -t\frac{q^2}{\varepsilon_0\hbar}\iint \frac{k_\parallel dk_\parallel d\varphi}{(2\pi)^2}[1-\cos(k_\parallel L\sin\varphi)][2\bar{n}_{\rm th}(k_\parallel v|\cos\varphi|)+1]\frac{e^{-2k_\parallel z\gamma(\varphi)}}{2k_\parallel\gamma(\varphi)} \times \mathrm{Im}\left\{r_p(k_\parallel,k_\parallel v|\cos\varphi|) - \frac{v^2}{c^2}[r_p(k_\parallel,k_\parallel v|\cos\varphi|)\cos^2\varphi - r_s(k_\parallel,k_\parallel v|\cos\varphi|)\sin^2\varphi]\right\}, \quad (15)$$

where we have set $\gamma^2(\varphi) = 1-\frac{v^2}{c^2}\cos^2\varphi$, with φ being the angle between the vectors $\mathbf{v}$ and $\mathbf{k}_\parallel$. We have also assumed for simplicity that $\mathbf{v}\perp(\boldsymbol{\rho}_1-\boldsymbol{\rho}_2)$. For slow velocities, we set $\gamma(\varphi)\simeq 1$ such that $\varepsilon(\omega)\omega^2/c^2 = \varepsilon(k_\parallel v|\cos\varphi|)k_\parallel^2 v^2\cos^2\varphi/c^2 \ll k_\parallel^2$ and hence $r_p(k_\parallel,\omega)\simeq r_p(\infty,\omega)\equiv r_p(\omega)$ and neglect the last line in Eq. (15) to obtain

$$\Gamma[c] = -t\frac{q^2}{\varepsilon_0\hbar}\iint\frac{k_\parallel dk_\parallel d\varphi}{(2\pi)^2}[1-\cos(k_\parallel L\sin\varphi)] \times[2\bar{n}_{\rm th}(k_\parallel v|\cos\varphi|)+1]\frac{e^{-2k_\parallel z}}{2k_\parallel}\mathrm{Im}\,r_p(k_\parallel v|\cos\varphi|). \quad (16)$$

Note that in this limit the reflection coefficient for s-polarized waves does not contribute. This is entirely due to the low-frequency dominance of the reflection coefficient for p-polarized waves, and not a geometric artefact.

At this point one might be tempted to set the velocity to zero in order to find the slow-particle limit explicitly. This proves to be tricky on two grounds. First, we have already performed the Markov approximation in which we assumed that all time scales are long compared to the inverse velocity (times a transverse wave number). Second, the imaginary part of the reflection coefficients vanish at zero argument by virtue of the Schwarz reflection principle, $r_p(\omega) = r_p^*(-\omega)$. However, the mean thermal photon number grows over all bounds, $\bar{n}_{\rm th}(k_\parallel v|\cos\varphi|)\mapsto k_BT/(\hbar k_\parallel v|\cos\varphi|)$. This is compensated by the fact that the reflection coefficient can be expanded for small arguments as $r_p(\omega)\simeq r_p(0)+i\omega r_p'(0)$ with $r_p(0), r_p'(0)\in\mathbb{R}$. In particular, for metals that can be described by a Drude permittivity $\varepsilon(\omega) = 1+i\sigma/(\varepsilon_0\omega)$ with conductivity σ, one finds that $r_p'(0) = 2\varepsilon_0/\sigma$.

Let us first look at what happens at finite temperature. Using standard integrals we find that

$$\Gamma_{\rm th}[c] = -t\frac{q^2k_BTr_p'(0)}{2\pi\varepsilon_0\hbar^2}\left[\frac{1}{2z}-\frac{1}{\sqrt{(2z)^2+L^2}}\right]. \quad (17)$$

Equation (17) is, due to the way in which the velocity and temperature combine, in fact a high-temperature expansion. The distance-dependent terms in brackets are easy to understand; the first is clearly the distance between one electron and its image in the metal (distance $2z$), and the second is the distance between one electron and the image of the other [distance $\sqrt{(2z)^2+L^2}$]. If the separation between the two electron paths vanishes ($L\to 0$), the spatial decoherence effect vanishes as L^2, as expected. Note that in this limit the spatial decoherence rate is independent of the velocity of the particles.

In the zero-temperature limit, there is nothing to compensate the velocity dependence of the reflection coefficient r_p, hence the decoherence functional turns into

$$\Gamma_0[c] = -t\frac{q^2vr_p'(0)}{2\pi^2\varepsilon_0\hbar}\left[\frac{1}{(2z)^2}-\frac{1}{(2z)^2+L^2}\right]. \quad (18)$$

Its structure is similar to the thermal result, Eq. (17), with the difference that the frequency $k_BT/\hbar$ has been replaced by the velocity divided by another power of the separation of the electrons from their mirror images. In both cases, the spatial decoherence vanishes as soon as $r_p'(0)$ vanishes. As noted above, this is the case for perfect conductors with $\sigma\to\infty$.

Comparing Eqs. (17) and (18) provides a means to distinguish between high-temperature and low-temperature limits (both are understood for slow velocities). The high-temperature limit is thus obtained whenever the condition $k_BT/\hbar \gg v/z$ is met. At room temperature, this implies that $z/v \gg \hbar/(k_BT) = 2.55\times 10^{-14}$ s is the time the electrons would need to reach the surface at normal incidence. Hence, already at room temperature the high-temperature limit (17) is valid for experimentally achievable distances and velocities. The zero-temperature limit (18) can be observed at very low temperatures (on the order of μK) with typical velocities on the order of 100 ms^{-1}.

Decoherence of neutral atoms. We now briefly discuss the path decoherence of neutral particles. The main difference between globally neutral systems of charges and charged particles such as electrons is that for the latter one cannot define a length scale associated with any transition, making it hard to perform meaningful approximations. In contrast, globally neutral systems of charges such as atoms inherently possess length scales associated with multipole transitions. For example, in a two-level atom with a dipole-allowed transition of energy $\hbar\omega_A$, the time dependence of the dipole moment operator $\hat{\mathbf{d}}$ is as $e^{-i\omega_At}$, which allows to define fast and slowly varying amplitudes as well as a length scale c/ω_A.

The decoherence functional $\Gamma[d]$ for neutral atoms can be derived in two equivalent ways. First, one can start from the current density $\mathbf{j}(\mathbf{r},t)$, expanded in terms of dipole moments. In this way, the current density takes the form $\mathbf{j}(\mathbf{r},t) = \frac{d}{dt}[\mathbf{d}(t)\delta(\mathbf{r}-\mathbf{r}_A(t))]$, where magnetic effects are omitted. Partial integration with respect to t and t' transforms the anticommutator function for the vector potential into that of the electric field, since $-\dot{\mathbf{A}} = \mathbf{E}$. The second, equivalent, way is to go back to the derivation of the decoherence functional itself and use the Hamiltonian $\hat{H} = -\hat{\mathbf{d}}\cdot\hat{\mathbf{E}}$ instead of $\hat{H} = -\hat{\mathbf{j}}\cdot\hat{\mathbf{A}}$.

The calculation proceeds as before, except that the thermal expectation value of the electric field [13]

$$\langle \hat{E}_i(\mathbf{r},\omega)\hat{E}_j^\dagger(\mathbf{r}',\omega')\rangle = \frac{\hbar\mu_0}{\pi}\omega^2[\bar{n}_{\rm th}(\omega)+1]{\rm Im}\, G_{ij}(\mathbf{r},\mathbf{r}',\omega)\delta(\omega-\omega') \quad (19)$$

is needed, with a similar expression holding for the normally ordered expectation values. For the classical dipole moments we insert $\mathbf{d}(x)=\mathbf{d}e^{-i\omega_A t}\delta(\mathbf{r}-\mathbf{r}_A)$, assuming that the dipole moment is associated with a transition of frequency ω_A and strength $\mathbf{d}$. Using again the Weyl expansion for the Green's function we find for the double time integral in the Markov approximation

$$\iint_{t_i}^{t_f} dt\, dt'\, e^{i(\mathbf{k}_\parallel\cdot\mathbf{v}-\omega_A\pm\omega)(t-t')} \mapsto \pi t\,\delta(\mathbf{k}_\parallel\cdot\mathbf{v}-\omega_A\pm\omega), \quad (20)$$

with $t=t_f-t_i$.

Assuming further that the velocity of the atom is low enough then, in the near-field limit, $|\mathbf{k}_\parallel\cdot\mathbf{v}|\ll\omega_A$. The frequency integral then picks out the atomic transition frequency ω_A which, together with the remaining prefactors, combines neatly to the expression for the decoherence functional

$$\Gamma[d]\simeq-\frac{t}{2}[\Gamma_{\rm th}(\omega_A,0)-\Gamma_{\rm th}(\omega_A,L)], \quad (21)$$

where $\Gamma_{\rm th}(\omega_A,L)$ is the thermal spatial coherence function at frequency ω_A and separation L [14]. At vanishing separation this reduces to the spontaneous decay rate $\Gamma_{\rm th}(\omega_A,0)$. From previous works [14,15] we know that we can write in the near-field limit

$$\Gamma_{\rm th}(\omega_A,L)\simeq\left[1-\frac{L^2}{32}\frac{\partial^2}{\partial z^2}\right]\Gamma_{\rm th}(\omega_A,0), \quad (22)$$

which, combined with the distance scaling in the near field of $\Gamma_{\rm th}(\omega_A,0)\propto 1/z_A^3$, leads to

$$\Gamma_{\rm th}[d]\simeq-t\frac{3L^2}{16z_A^2}\Gamma_{\rm th}(\omega_A,0). \quad (23)$$

The result is as expected; the path decoherence rate is again quadratic in the path separation L, and it is proportional to the spontaneous decay rate $\Gamma_{\rm th}(\omega_A,0)$ in the near-field limit. Note that this result expresses the decoherence of external degrees of freedom in terms of relaxation rates of the internal dynamics.

Conclusion. We have shown that the effect of path decoherence of both charged and neutral particles near macroscopic bodies can be traced back to the interaction of the particles with their mirror images inside the material. As expected, as soon as the material could in principle determine the positions of the images by performing a measurement, i.e., due to nonzero resistivity, path decoherence sets in. This is consistent with the interpretation that the material gains *which path* information and thus destroys the coherence between the beams. Regardless of whether or not the particles are charged, the spatial decoherence rate for small path separations L is proportional to L^2. This result confirms previous theoretical predictions based on decoherent histories [6] and experimental observations [3,8].

The lack of an intrinsic length scale for charged particles such as electrons means that taking the slow-velocity limit is equivalent to the zero-frequency limit. The transition frequency of a neutral particle, however, sets a time scale and therefore a length scale with respect to which one can define a nonretarded limit.

The approach based on macroscopic quantum electrodynamics provides an alternative viewpoint to microscopic models of path decoherence. It covers all choices of surface material, from metals to semiconductors and dielectrics, and describes the decoherence of charged and neutral particles. It also allows for the inclusion of different geometries, if the appropriate Green's tensor is specified. The insight that one gains by using this approach is an effective way to describe experimental evidence of either existing experiments using electrons [8] of planned experiments with ions.

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