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Decoherence and recoherence from vacuum fluctuations near a conducting plate

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The interaction between particles and the electromagnetic field induces decoherence generating a small suppression of fringes in an interference experiment. We show that if a double-slit-like experiment is performed in the vicinity of a conducting plane, the fringe visibility depends on the position (and orientation) of the experiment relative to the conductor's plane. This phenomenon is due to the change in the structure of vacuum induced by the conductor and is closely related to the Casimir effect. We estimate the fringe visibility both for charged and for neutral particles with a permanent dipole moment. The presence of the conductor may tend to increase decoherence in some cases and to reduce it in others. A simple explanation for this peculiar behavior is presented.

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I. INTRODUCTION

The interaction of a quantum system with its environment is responsible for the process of decoherence, which is one of the main ingredients to understand the quantum-classical transition [1]. In some cases, the interaction with the environment cannot be switched off. This is the case for charged particles that unavoidably interact with the electromagnetic field. As this interaction is fundamental, its effect is present in any interference experiment. In this paper we will analyze the influence of a conducting boundary in the decay of the visibility of interference fringes in a double slit experiment performed with charged particles (or neutral particles with a dipole moment). The reduction of fringe visibility is induced by the interaction between the particles and the electromagnetic field. Some aspects of this problem have been analyzed before. In fact, it is known that for charged particles, the interaction between the system (the particle) and the environment (the electromagnetic field) induces a rather small decoherence effect even if the initial state of the field is the vacuum [2–10]. A particularly simple expression for the decay in the fringe visibility was obtained in Refs. [2,3]: Assuming an electron in harmonic motion (with frequency Ω) along the relevant trajectories of the double slit experiment, the fringe visibility decays by a factor $(1-P)^2$, where P is the probability that a dipole $p=eR$ oscillating at frequency Ω emits a photon (R is the characteristic size of the trajectory). This result is in accordance with the idea that decoherence becomes effective when a record of the state of the system is irreversibly imprinted in the environment. In this case, after photon emission, if the electron follows the trajectory $\vec{X}_1(t)$ of the double slit experiment (see Fig. 1) it becomes correlated with a state of the environment $|E_1(t)\rangle$. This state is different from the one with which the electron correlates if it follows the trajectory $\vec{X}_2(t)$. The absolute value of the overlap between these two different states is precisely given by $(1-P)^2$.

In this paper we will analyze how the fringe visibility is modified when performing a double slit interference experiment in the vicinity of a conducting plane. One could ask what is the reason why the presence of a perfect conductor could possibly modify the fringe visibility in a double slit experiment. The answer to this question is not complicated: It is well known that the presence of a conducting plane enforces nontrivial boundary conditions on the electromagnetic field. These conditions strongly affect the nature of the space of physical states of the quantum field. As we argued above, the fringe visibility is determined by the absolute value of the overlap between two physical states of the field (see also below). Therefore, as the field states are affected by the boundary conditions, it is not unexpected that the overlap between them is affected as well. For this reason, one expects decoherence to be influenced by the presence of a conducting plane.

The effect we study here is as fundamental as the Casimir force between two conductors [11], which is also caused by the nontrivial boundary conditions. Being the effect of fundamental origin, one expects that the impact of the conductors on decoherence should be understandable in simple and somewhat intuitive terms. However, until now this was not the case. In fact, the way in which a conductor modifies the fringe visibility in an electron double slit experiment was analyzed before by Ford in Refs. [7–9]. The results contained in those papers are far from intuitive. The analysis we present here will serve not only to correct these previous results [7–9], which turned out to be erroneous. Thus, we will also show that the effect of the conductor is quite remarkable and simple to understand. As we will see, the presence of the conducting plane may produce more decoherence in some cases and less decoherence in others. For example, we will show that for a double slit experiment with charged particles, if a conducting plate is placed perpendicular to the trajectories of the interfering particles, the fringe visibility decreases with respect to the vacuum case (absence of conducting plate). However, if the plate lies parallel to the trajectories of the charges, the contrast increases (the system recoheres). We will show that this peculiar behavior can be understood in simple terms by using a variation of the method of images. We will show a similar result for the case of neutral particles with permanent electric or magnetic dipole moment. Also, we will show that the magnitude of the

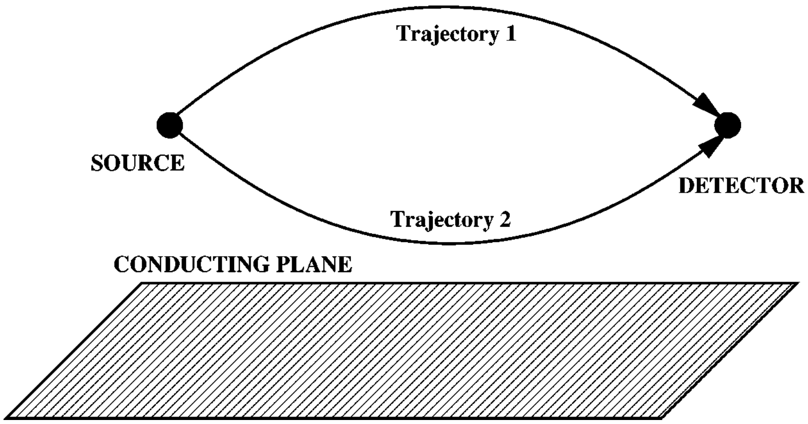

FIG. 1. Scheme for a double-slit-like experiment near a conducting plane. The component of the velocity in the direction from the source to the detector is assumed to be constant.

effect can be easily estimated.

There are several interesting physical effects connected with the one we are analyzing here. As we mentioned above, it is well known that a conducting boundary modifies the properties of the zero-point fluctuations, and therefore could affect the interference experiments of particles that interact with the electromagnetic field. Apart from the above mentioned Casimir force between conductors, other consequences of this same phenomenon is the Casimir-Polder force [12] affecting a probe particle in the vicinity of a conductor. These phenomena, that have been experimentally verified [13], are close relatives of the process we are studying here. In fact, the Casimir-Polder force can be thought as the dispersive counterpart of the decoherence effect we will discuss. The influence of boundaries on the electromagnetic vacuum is also responsible for changes in atomic lifetimes and interference phenomena for light emitted by atoms near conducting surfaces [14].

The paper is organized as follows: In Sec. II we outline the main calculation that needs to be done to compute the fringe visibility in a double slit experiment with charges and neutral particles with permanent dipole moment. In Sec. III we show how to evaluate the fringe visibility in vacuum. In Sec. IV we show how to compute the effect induced by the nontrivial boundary condition generated by the perfect conductor. We also discuss there how to evaluate the fringe visibility for experimentally relevant situations estimating the significance of the effect and showing how it can be intuitively understood. In Sec. V we present a short summary of our results.

II. FRINGE VISIBILITY IN A DOUBLE SLIT EXPERIMENT

Let us first outline a simple method to compute the effect of electromagnetic interactions on the fringe contrast. We consider two electron wave packets that follow well defined trajectories $\vec{X}_1(t)$ and $\vec{X}_2(t)$ that coincide at initial $(t=0)$ and final $(t=T)$ times as shown in Fig. 1. In the absence of environment, the interference depends on the relative phase between both the wave packets at $t=T$. Because of the interaction with the quantum electromagnetic field, the interference pattern is affected. This effect can be calculated as follows: We assume an initial state of the combined particle-field system of the form $|\Psi(0)\rangle=(|\phi_1\rangle+|\phi_2\rangle)\otimes|E_0\rangle$. Here $|E_0\rangle$ is the initial (vacuum) state of the field and $|\phi_{1,2}\rangle$ are two states of the electron that are localized around the initial point and that in the absence of other interaction continue to be localized along the trajectories $X_{1,2}(t)$, respectively. At later times, due to the particle-field interaction, the state of the combined system becomes

$$|\Psi(t)\rangle=(|\phi_1(t)\rangle\otimes|E_1(t)\rangle+|\phi_2(t)\rangle\otimes|E_2(t)\rangle). \quad (1)$$

Thus, the two localized states $|\phi_1(t)\rangle$ and $|\phi_2(t)\rangle$ become correlated with two different states of the field. Therefore, the probability of finding a particle at a given position turns out to be

$$\begin{aligned}\mathrm{prob}(\vec{X},t)=&|\phi_1(\vec{X},t)|^2+|\phi_2(\vec{X},t)|^2+2\mathrm{Re}(\phi_1(\vec{X},t)\phi_2^*(\vec{X},t)\\&\times\langle E_2(t)|E_1(t)\rangle).\end{aligned} \quad (2)$$

The overlap factor $F=\langle E_2(t)|E_1(t)\rangle$ is responsible for two effects. Its phase produces a shift of the interference fringes (the Aharonov-Bohm effect can be obtained in this way when the initial state of the magnetic field has a nonzero expectation value). The absolute value $|F|$ is responsible for the decay in the fringe contrast, which is the phenomenon we will analyze here. Calculating it is conceptually simple since F is nothing but the overlap between two states of the field that arise from the vacuum under the influence of two different sources (this factor is identical to the Feynman-Vernon influence functional [15]). Each of the two states of the field can be written as

$$|E_a(t)\rangle=T\left[\exp\left(-i\int d^4x J_a^\mu(x)A_\mu(x)\right)\right]|E_0\rangle, \quad (3)$$

where $J_a^\mu(x)$ is the conserved four-current generated by the particle following the classical trajectory $\vec{X}_a(t)$, i.e., $J_a^\mu(\vec{X},t)=(e,e\dot{\vec{X}}_a(t))\delta^3(\vec{X}-\vec{X}_a(t))$, $(a=1,2)$. Using this, it is simple to derive an expression for the overlap: The simplest way to do this is based on the observation that as the QED action is quadratic in the fields, the overlap must be a Gaussian functional of the two currents J_1 and J_2. Thus, we can write the most general Gaussian functional ansatz for F as

$$F=\exp\left(-i\int\int d^4x_1 d^4x_2 J_a^{\mu}(x_1)G_{\mu\nu}^{ab}(x_1,x_2)J_b^{\nu}(x_2)\right)\times\exp\left(-i\int d^4x J_a^{\mu}(x)C_{\mu}^{b}(x)\right), \quad (4)$$

where a summation over the indices $a,b=1,2$ is implicit. On the other hand, we can explicitly write down the expression for the overlap as

$$F=\langle E_0|\tilde{T}\left[\exp\left(i\int d^4x J_2^{\mu}(x)A_{\mu}(x)\right)\right]\times T\left[\exp\left(-i\int d^4x J_1^{\mu}(x)A_{\mu}(x)\right)\right]|E_0\rangle. \quad (5)$$

The kernels $G_{\mu\nu}^{ab}$ and C_{μ}^{a} appearing in Eq. (4) can be determined by identifying the functional derivatives of Eqs. (4) and (5). In this way, one can relate C_a and G_{ab} with the one and two point functions of the field operators. As we are only interested in the absolute value of the overlap, we will only present the result for this quantity here. Denoting $|F|=\exp(-W_c)$, we get

$$W_c=\frac{1}{2}\int d^4x\int d^4y(J_1-J_2)^{\mu}(x)D_{\mu\nu}(x,y)(J_1-J_2)^{\nu}(y), \quad (6)$$

where $D_{\mu\nu}$ is the expectation value of the anticommutator of two field operators

$$D_{\mu\nu}(x,y)=\tfrac{1}{2}\langle\{A_{\mu}(x),A_{\nu}(y)\}\rangle. \quad (7)$$

From this derivation, it is clear that the probability for vacuum persistence in the presence of a source J_1^{μ}, which is given by $|\langle 0|E_i\rangle|^2$, can be obtained from the above expression by simply setting $J_2^{\mu}=0$. Taking this into account, it is possible to attach a simple physical interpretation to the overlap. The square of the overlap, $F^2=|\langle E_1(t)|E_2(t)\rangle|^2$, is equal to the vacuum persistence probability in the presence of a source $j_{\mu}=(J_1-J_2)_{\mu}$. This source corresponds to a time dependent electric dipole $e(\vec{X}_1(t)-\vec{X}_2(t))$, which is directed from one trajectory towards the other [5]. Then, decoherence arises when this (fictitious) dipole emits a photon imprinting a record in the electromagnetic environment that could in principle be used to distinguish between the two paths.

A conceptually similar and physically interesting problem can be analyzed along the same lines: the decoherence of neutral particles with a nonvanishing permanent dipole moment. In such case we can model the particle-field interaction using a Lagrangian $L_{int}=P_{\mu\nu}(x)F^{\mu\nu}(x)$. Here, $F_{\mu\nu}$ is the field strength tensor and $P_{\mu\nu}$ is a totally antisymmetric tensor whose nonvanishing components can be interpreted as the electric and magnetic dipole densities (in the laboratory frame). In fact, the electric dipole moment $\vec{p}$ and magnetic dipole moment $\vec{m}$ of the particles moving along a trajectory $\vec{X}(t)$ are obtained in terms of the dipolar tensor as $P_{0i}=p_i(t)\delta^3(\vec{X}-\vec{X}(t))/2$ and $P_{ij}=\epsilon_{ijk}m_k(t)\delta^3(\vec{X}-\vec{X}(t))/2$. In this case we can perform a calculation which is, *mutatis mutandi,* similar to the one above and show that the overlap $F=\exp(-W_d)$ is

$$W_d=\frac{1}{2}\int\int d^4x d^4y(P_1-P_2)^{\mu\nu}(x)K_{\mu\nu\rho\sigma}(x,y)\times(P_1-P_2)^{\rho\sigma}(y). \quad (8)$$

The kernel appearing in this equation, defined as $K_{\mu\nu\rho\sigma}(x,y)=\langle\{F_{\mu\nu}(x),F_{\rho\sigma}(y)\}\rangle$, can be expressed in terms of derivatives of $D_{\mu\nu}(x,y)$

$$K^{\mu\nu\rho\sigma}(x,y)=[\partial_x^{\mu}\partial_y^{\rho}D^{\nu\sigma}(x,y)+\partial_x^{\nu}\partial_y^{\sigma}D^{\mu\rho}(x,y)-\partial_x^{\mu}\partial_y^{\sigma}D^{\nu\rho}(x,y)-\partial_x^{\nu}\partial_y^{\rho}D^{\mu\sigma}(x,y)]. \quad (9)$$

III. EVALUATING THE FRINGE VISIBILITY IN VACUUM

In what follows we will present results for the *decoherence factors* W_c and W_d (the subscripts stand for "charges" and "dipoles"). To compute W_c we need the two point function appearing in Eq. (6). In the Feynman gauge and in the absence of conducting plates it is

$$D_{\mu\nu}^{(0)}(x,y)=-\eta_{\mu\nu}\int\frac{d^3\vec{k}}{(2\pi)^3 2k}e^{i\vec{k}(\vec{x}-\vec{y})}\cos(k(x_0-y_0)), \quad (10)$$

where the superscript (0) identifies this as the vacuum contribution. We will assume that the trajectories are symmetric and write $\vec{X}_1(t)=-\vec{X}_2(t)=x(t)\hat{\mathrm{x}}$. This is enough to describe a typical double slit experiment from the point of view of an observer moving at constant velocity from the source to the detector. In such case we can evaluate the overlap and obtain a relatively simple expression for the decoherence factor W_c,

$$W_c^{(0)}=e^2\int\frac{d^3\vec{k}}{8\pi^3 k}\left(1-\frac{k_j^2}{k^2}\right)\left|\int_{-\infty}^{\infty}dt\,\dot{x}(t)\cos[k_x x(t)]e^{ikt}\right|^2. \quad (11)$$

To obtain this equation from Eq. (6) one should first explicitly perform the spatial integration, write the kernel $D_{\mu\nu}^{(0)}(x,y)$ in momentum space and finally use the conservation of the four-current to cancel the contribution of the temporal and longitudinal components of the current. This result was obtained first in Ref. [2] using the Coulomb gauge and a slightly different but equivalent method. It can be simplified further by assuming the validity of the dipole approximation $\cos[k_x x(t)]\simeq 1$ (which is consistent in the nonrelativistic limit). Doing this, one can evaluate the decoherence factor for some special trajectories. In fact, for adiabatic trajectories, where $x(t)=R\exp[-t^2/T^2]$, we find that $W_c^{(0)}=2e^2v^2/3\pi$, where $v=R/T$ is a characteristic velocity. This result is finite and free of any cutoff dependence. However, for trajectories evolving over a finite time the situation is different. Thus, assuming that the motion starts at $t=0$, ends at $t=T$, and that is composed of periods of constant velocity v, or constant acceleration v/τ, we obtain a result that diverges logarithmically when $\tau\to 0$, $W_c^{(0)}=2e^2v^2\log[T/\tau]/\pi^2$ (if $\tau/T\ll 1$). Previous results [5] were obtained for trajectories with discontinuous velocity using a natural UV cutoff arising from the finite size of the electron. The results of Ref. [5] agree with ours if the high-frequency cutoff is identified with $1/\tau$ (the results of Refs. [7–9] are incorrect, see below). Thus, the cutoff dependence disappears in the adiabatic case and is a consequence of abrupt changes in velocity and the instantaneous preparation of the initial state.

The concept of a decoherence rate can also be introduced in this context as follows: Let us now assume that the two wave packets are superposed after oscillating N times, and that the time to complete one oscillation is much shorter than the period between oscillations. If $v(t)$ denotes the velocity during one oscillation of period T for a sequence of N identical oscillations separated by $\Delta T\gg T$, we can write the velocity as

$$\dot{x}_N(t)=\sum_{n=0}^{N-1} v(t-n\Delta T). \tag{12}$$

Therefore, the temporal integral appearing in the decoherence factor (11) can be evaluated as

$$I_N\equiv\int_{-\infty}^{\infty}dt\ \ \dot{x}_N(t)e^{ikt}=I\,\frac{e^{-ikN\Delta T}-1}{e^{ik\Delta T}-1}, \tag{13}$$

where I is the result of the temporal integral for a single oscillation. Inserting Eq. (13) into Eq. (11) we find that, after a large number of oscillations, the decoherence factor is proportional to N: $W_c^{(0)}=NW_c^{(0)}(1)$, where $W_c^{(0)}(1)$ is the decoherence factor in a single oscillation.

We now describe the results for the case of neutral particles with permanent dipole moments. The calculation is a bit more tedious than for the case of charges. To avoid cumbersome details we will evaluate the decoherence factor only under somewhat simplified assumptions. We will consider that the dipole moments $\vec{p}$ and $\vec{m}$ remain constant along the trajectories. We will also assume symmetric trajectories and write $\vec{X}_1(t)=-\vec{X}_2(t)=x(t)\hat{\jmath}$. In this case, we can perform the spatial integration in Eq. (8) and use the antisymmetric nature of the polarization tensor $P_{\mu\nu}$ (which has the same effect than the conservation of the four-current in the previous case and enables us to cancel the contribution of temporal and longitudinal modes of the field). Finally, we can use the momentum representation for the two point function obtained by replacing Eq. (10) into Eq. (9). In this way we obtain a relatively simple final expression for the decoherence factor. It reads

$$W_d^{(0)}=\int\ \frac{d^3\vec{k}}{8\pi^3}k\left\{\vec{p}^{\,2}\left(1-\frac{k_p^2}{k^2}\right)+\vec{m}^{\,2}\left(1-\frac{k_m^2}{k^2}\right)\right\}\times\left|\int_0^t dt'\sin[k_j x(t')]e^{ikt'}\right|^2. \tag{14}$$

Using again the dipole approximation for the adiabatic trajectory we can show that the above decoherence factor is such that $W_d^{(0)}/W_c^{(0)}\simeq p^2/e^2T^2$ for a purely electric dipole (a similar expression is obtained for the purely magnetic case). This ratio is typically much smaller than one (i.e., it is of the order of the square of the typical dipole length in units of the total length of the trajectory).

IV. FRINGE VISIBILITY IN THE PRESENCE OF A CONDUCTOR

We will now show how the above results are modified by the presence of a perfect conductor located in the plane $z=0$. To consider the effect of the conductor we only need to use the appropriate two point functions that obey the correct boundary conditions. For the case of charges, the kernel $D_{\mu\nu}$ is the sum of two terms [16]

$$D_{\mu\nu}=D_{\mu\nu}^{(0)}+D_{\mu\nu}^{(B)}. \tag{15}$$

The vacuum term is the same as in Eq. (10). The contribution of the boundary conditions [identified by the superscript (B)] can be obtained by the method of images and is

$$D_{\mu\nu}^{(B)}(x,y)=(\eta_{\mu\nu}+2n_\mu n_\nu)\int\frac{d^3\vec{k}}{(2\pi)^3 2k}\exp(i\vec{k}(\vec{x}-\vec{y}\,'))\times\cos(k(x_0-y_0)). \tag{16}$$

Here n^μ is the normal to the plane and $\vec{y}\,'$ is the position of the image point of $\vec{y}$ (a prime denotes a vector reflected with respect to the plane, i.e., $\vec{y}\,'=(y_x,y_y,-y_z)$. Inserting Eq. (16) into Eq. (6) we can derive a formula for the contribution of the boundary to the decay of the interference fringes. The complete equation is involved but it becomes considerably simpler if we restrict to the case where the trajectories are either perpendicular or parallel to the conductor's plane. In this case we can write $\vec{X}_{1,2}=z_0\hat{z}\pm x(t)\hat{\jmath}$, where $\hat{\jmath}$ defines a fixed vector aligned either along the $\hat{z}$ axis or along the plane perpendicular to it. The contribution of the conductor to the decoherence factor is

$$W_c^{(B)}=-\hat{\jmath}\hat{\jmath}'e^2\int\ \frac{d^3\vec{k}}{8\pi^3 k}\left(1-\frac{k_j^2}{k^2}\right)\times e^{2ik_z z_0}\left|\int_0^t dt'\dot{x}(t')\cos[k_j x(t')]e^{ikt'}\right|^2. \tag{17}$$

It is interesting to note that the sign of $W_c^{(B)}$ is determined by the orientation of $\hat{\jmath}'$ relative to $\hat{\jmath}$. $W_c^{(0)}$ is negative when the trajectories are parallel to the conductor's plane (since in that case $\hat{\jmath}'=\hat{\jmath}$). On the other hand, $W_c^{(0)}$ is positive when the trajectories are perpendicular to the plane (where $\hat{\jmath}'=-\hat{\jmath}$). At small distances to the plane ($z_0\simeq 0$) we can see from Eq. (17) that $|W_c^{(B)}|\simeq W_c^{(0)}$. Therefore, if the trajectories are perpendicular to the plane, in the limit of small distances the decoherence factor is $W_c=W_c^{(0)}+W_c^{(B)}\simeq 2W_c^{(0)}$. The effect of the conductor is to double the decoherence factor. However, if the trajectories are parallel to the conductor the effect is exactly the opposite. As $W_c^{(B)}$ is negative, the conductor produces *recoherence* increasing the contrast of the fringes. In fact, for small distances the decoherence factor tends to vanish since $W_c = W_c^{(0)} + W_c^{(B)} \simeq 0$.

It is remarkable that these results can be understood using the method of images. For this, we should take into account that, as mentioned above, the decoherence factor for a double slit experiment with charge can be related to the probability of photon emission from a source characterized by a four-current $j_\mu = (J_1 - J_2)_\mu$, which is the difference between the two interfering currents. This source corresponds to that of a varying dipole $p = ex(t)$. So, to understand our result one has to analyze how does the conductor affects the photon emission from this fictitious dipole. When the conducting plane is parallel to the dipole, the image dipole is $\vec{p}_{im} = -\vec{p}$. Therefore, the total dipole moment vanishes, and so does the probability to emit a photon. The image dipole cancels the effect of the real dipole and this produces the recovery of the fringe contrast. On the other hand, when the conductor is perpendicular to the trajectories, the image dipole is equal to the real dipole $\vec{p}_{im} = +\vec{p}$. Therefore, the total dipole is twice the original one. This in principle would lead us to conclude that the total decoherence factor $W_c = W_c^{(0)} + W_c^{(B)}$ should be four times larger than $W_c^{(0)}$. However, one should take into account that in the presence of a perfect mirror photons can only be emitted in the $z \geqslant 0$ region. This introduces an additional factor of 1/2 that gives rise to the final result $W_c \simeq 2W_c^{(0)}$. One should remark that the dipole used in the above reasoning is not a real dipole but the effective (fictitious) dipole created by the opposite charges following the two interfering trajectories.

In the case when the interfering particles are neutral but carry a permanent dipole moment, the effect of the conductor can also be taken into account using the method described above. For simplicity we will only consider trajectories that are parallel to the plane [i.e., $\vec{X}_{1,2} = z_0 \hat{z} \pm x(t)\hat{\jmath}$] and assume that the dipole moments are either perpendicular or parallel to the conductor (the general case is more complex but the essential features can be seen here). Using this we obtain

$$W_d^{(B)} = -\int \frac{d^3\vec{k}}{32\pi^3} k \left\{ \vec{p}\vec{p}\,' \left(1 - \frac{k_p^2}{k^2}\right) - \vec{m}\vec{m}\,' \left(1 - \frac{k_m^2}{k^2}\right) \right\} \times e^{2ik_z z_0} \left| \int_0^t dt' \sin[k_j x(t')] e^{ikt'} \right|^2. \tag{18}$$

Thus, if the reflected dipole $\vec{p}\,'$ has the opposite direction than $\vec{p}$ (which is the case when $\vec{p}$ is parallel to the plate) the conductor tends to increase decoherence (since the contribution of the electric dipole to $W_d^{(B)}$ is positive). Likewise, when $\vec{p}$ is perpendicular to the plane, $\vec{p} = \vec{p}\,'$ and the contribution of the electric dipole to $W_d^{(B)}$ is negative. Therefore, in this case the conductor produces *recoherence* instead of decoherence. The opposite effect is found for the magnetic dipole. Indeed, when $\vec{m}\,' = -\vec{m}$ (magnetic dipole perpendicular to the plane) the conductor produces recoherence while more decoherence is produced if the magnetic dipole is parallel to the plane. These features can also be understood by thinking in terms of the image dipoles that are generated by the conductor. Thus, both when $\vec{p}$ is perpendicular to the plane or when $\vec{m}$ is parallel, the direction of the image dipoles coincide with the source dipoles. In such case the decoherence increases. In the opposite situation ($\vec{p}$ parallel or $\vec{m}$ perpendicular to the plane) the effect of the conductor is to introduce recoherence. Again, in the limit of small distances the absolute value of $W_d^{(0)}$ and $W_d^{(B)}$ coincide and therefore the decoherence factor doubles with respect to the vacuum case.

Both for the case of charges and for the case of dipoles the boundary contribution to the decoherence factor decays algebraically with the distance to the conductor (in the limit of large distances). In the opposite case, for small separations, explicit expressions can be obtained. For example, for charges moving close and parallel to the conductor, the lowest order contribution of W_c can be shown to depend quadratically on z_0. As expected, it exactly coincides with the decoherence factor produced by an electric dipole $p = 2ez_0$ in vacuum (with an additional factor of 1/2 that takes into account that photons can only be emitted with $z \geqslant 0$).

V. CONCLUSIONS

The impact of conducting boundaries on the fringe visibility for the case of electrons was previously examined in Refs. [7–9]. The results reported in such papers are not correct because of the use of an erroneous approximation for low velocities. In fact, the author performed a low velocity expansion by neglecting the spatial components of the four-currents appearing in Eq. (6), which are indeed linear in velocity. However, this approximation clearly violates the conservation of the four-current and makes the final result unphysical. In fact, the final expressions obtained in Refs. [7–9] are gauge dependent and could violate unitarity allowing for $|F| > 1$ (thus, the decoherence factor computed in such papers is not positive defined). As mentioned above, in the correct result the contribution of the temporal component of the four-current is canceled out by the one corresponding to the longitudinal component. Thus, only the transverse modes are physical and determine the value of the decoherence factor W, which being manifestly positive ensures that $|F| \leqslant 1$ [note that positivity of the decoherence factor W is manifested from Eqs. (11) and (14)].

Another interesting point discussed in Refs. [7,8] is the possibility of observing decoherence induced by vacuum fluctuations in a vetolike experiment. In such experiment one ideally surrounds the setup with photon detectors and keeps only the data coming from those runs where no photons were actually detected. In Ref. [8] the author argues that the decoherence factor in such experiments is different from the one we analyzed here. In our view, in this kind of veto experiment the decoherence factor actually vanishes and the fringe visibility is not affected. Indeed, in an ideal vetolike experiment one is really performing a projective measurement of the state of the electromagnetic field (which is detected to be the vacuum). Therefore, the state of the combined system (particle plus field) is obtained from Eq. (1) by projecting the electromagnetic field into the vacuum state (and normalizing accordingly). The probability of finding a particle at a given position turns out to be

$$\mathrm{prob}(\vec{X},t)=|Z_1\phi_1(\vec{X},t)|^2+|Z_2\phi_2(\vec{X},t)|^2 + 2\mathrm{Re}(Z_1\phi_1(\vec{X},t)Z_2\phi_2^*(\vec{X},t)), \qquad (19)$$

where

$$Z_i=\frac{1}{\sqrt{|\langle 0|E_1\rangle|^2+|\langle 0|E_2\rangle|^2}}\langle 0|E_i\rangle. \qquad (20)$$

This shows that, although the particular values of Z_i depend on the position and orientation of the conducting plane, the visibility of the fringes is not affected by the interaction with the field (note also that for the simplest case of symmetric trajectories we always have $Z_i=1/\sqrt{2}$). Thus, as there is no decoherence factor in the last term of Eq. (19), contrary to what was claimed in Refs. [7–9], no decoherence can be seen in this kind of veto experiment.

As we mentioned in the Introduction, our work not only serves the purpose of correcting previous results. In fact, we showed that the effect of conducting boundaries on the fringe visibility of double slit interference experiments can be intuitively understood in simple terms. The way in which decoherence is affected is similar to the manner in which atomic emission properties are modified by the presence of conducting boundaries. Thus, the effect of the boundaries does not have a well defined sign and may produce either more decoherence or complete recoherence (i.e., smaller or higher fringe visibility than in vacuum) depending on the orientation of the relevant trajectories with respect to the conductor's plane. Most notably, one can predict the sign of the effect (i.e., whether the conductor decreases or enhances the fringe visibility) by using a reasoning based on the method of images. The effect discussed here is conceptually important due to its fundamental origin (i.e., it is always present), but its magnitude is too small to be under the reach of current experiments involving interference of neutral atoms in the vicinity of conducting planes [17]. However it may be possible to enhance the effect by considering other geometries, like a periodically corrugated conducting plane. Work in this direction is in progress.

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