Physical picture of photodetachment in external fields: A way to its assessment

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The photodetachment of the negative ion H^- by an uv photon in the presence of a low-frequency (LF) electric field and a static magnetic field is investigated. The case is treated in which the ejected electron may exchange a large number of LF photons and the LF field period is much larger than the time required for the electron ejection. The LF field polarization is taken along the magnetic field direction, with the uv photon polarization either along or perpendicular to **B**. The intensities of the two external fields are taken to be sufficiently weak to neglect the Stark and diamagnetic shifts of the initial state. The reported results support the physical picture according to which (a) the main role of the LF field is that of creating a repulsive barrier in the direction of the ejected electron free motion and (b) the oscillations observed in the experiment [M. C. Baruch, T. F. Gallagher, and D. J. Larson, Phys. Rev. Lett. **65**, 1336 (1990)] are due to repeated reflections by such a barrier. The presence of a magnetic field enhances the interference effects and photodetachment experiments carried out with the additional presence of a magnetic field could provide more information on the precise role the LF field plays in the process.

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I. INTRODUCTION

A recent photodetachment experiment [1], in which a single uv photon with frequency ω_1 detached the electron of the Cl^- negative ion, has shown that the presence of a low-frequency (LF) electric field in the microwave range with frequency ω_2 gives rise, in the cross section, to oscillations reminiscent of those produced by a static (dc) electric field [2].

In analogy with the explanation proposed for the photodetachment in the presence of a dc electric field [3], the oscillations yielded by a LF electric field also have been interpreted [4–6] as due to the interference of the direct photoelectron wave with the wave reflected by the repulsive barrier of finite height originated by the LF field. A detailed analysis of the analogies and differences between the photodetachment in the presence of a dc field and that with a LF field has been reported in Ref. [7], in which the distributions of the ejected electrons are compared vs the kinetic energy E_T in the plane perpendicular to the direction of the assisting fields. The results of that analysis showed that the quasistatic (QS) approximation for the LF electric field is satisfactory for a limited interval of E_T values only, while the interference effects persist for a larger interval. If the oscillations exhibited by the cross section in the presence of a LF field are correctly interpreted as due to a barrier, then the additional presence of a magnetic field should enhance the interference effects, as it makes unidimensional the free motion of the ejected electron [8]. More to the point, the presence of a magnetic field should allow one to control the interference effects and, accordingly, the oscillations of the cross section. As a matter of fact, if the distance between two adjacent Landau levels (ω_c) is larger than the effective height of the barrier produced by the LF field, by increasing the uv photon energy the free motion energy may be increased to allow the ejected electron to overcome the barrier while still not being enough to excite the nearest Landau level. As the interference effects are due to barrier reflection, in such a case they should decrease and the oscillations on the cross sections disappear. In the opposite case, i.e., when the distance between two Landau levels is smaller than the height barrier, whatever the energy of the uv photon, we expect that most of the available energy after detachment will go into excitation of Landau levels. As a consequence, most of the ejected electrons will have insufficient free motion energy to overcome the barrier. The net outcome, in this case, should be repeated barrier reflection yielding not only the presence of interference effects, but also their enhancement. As the effective barrier height is of the order of $2\Delta_2$ [6], where $\Delta_2=E_2^2/4\omega_2^2$ is the ponderomotive shift of the LF electric field of frequency ω_2 and amplitude E_2, values of the parameter $\gamma=\omega_c/2\Delta_2$, greater or smaller than unity, may be taken to characterize respectively the two cases. For $\gamma>1$, the interference effects due to the LF electric field are well separated from the effects coming from the successive opening of Landau levels; depending on the uv photon energy, the interference effects may be absent. For $\gamma<1$, the interference effects are always present. The above considerations prompt the investigation of the photodetachment in the presence of a LF electric field with the simultaneous presence of a constant magnetic field as an interesting case study, in which one may elucidate, in a clearer way, the role of the LF field in the experiment of Ref. [1]. Furthermore, they offer interesting physics with the possibility of controlling the elementary process.

In the present work, we treat the photodetachment of the negative ion H^- in the presence of a magnetic field and a LF field. In particular we investigate the effects of the external fields on the ejected electron; accordingly, the intensities of such fields are taken to be small enough to neglect the Stark and the diamagnetic shifts of the ground state. The paper is organized as follows. In Sec. II we derive the cross section of photodetachment when the LF and magnetic fields are along the same z direction, while the uv photon polarization is either parallel or perpendicular to z. In Sec. III we report the cross sections for both geometries and analyze their basis characteristics, comparing them both with the cross section without the LF field and with the cross section in the presence of a QS field. Section IV reports the concluding remarks. The Appendix contains the derivation of the cross section in constant fields, using a gauge allowing a direct comparison with the case under consideration. Atomic units are used throughout the paper.

II. THEORY

In this section we derive the cross section of the photodetachment by absorption of an uv photon of frequency ω_1 and energy comparable to that of the negative-ion ground state I_0 in the presence of a constant magnetic field $\mathbf{B}$ and a radiation field $\mathbf{E}_2$ with frequency $\omega_2 \ll \omega_1$. During the photodetachment, the negative ion may exchange an arbitrary number of photons with the field $\mathbf{E}_2$.

We represent the negative ion by a one-electron bound system, in which the electron moves in a zero-range static potential $V(r)$. Within such a model, the exact S matrix of the photodetachment process by two radiation fields in the presence of a magnetic field $\mathbf{B}$ is written as [9]

$$S_{fi} = -i\int_{-\infty}^{+\infty} \langle \Psi_f(r,t)|V(r)|\Psi_i^+(r,t)\rangle dt \ . \tag{1}$$

In Eq. (1) Ψ_i^+ is the exact wave function of the negative ion in the presence of all the fields, for $t\to-\infty$, tending to the negative-ion bare state; Ψ_f is the wave function of the free electron dressed by the radiation and magnetic fields. Assuming $\mathbf{B}$ to be along z with the vector potential $\mathbf{A}_B = -By\mathbf{x}$, Ψ_f may be written as [10]

$$\Psi_f(r,t) = c_n \exp(-iE_n t)\exp(i\mathbf{p}\cdot\mathbf{r})\exp\left[-\frac{\xi^2}{2}\right]H_n(\xi) \times\exp\left\{i\frac{[\mathbf{A}_1(t)+\mathbf{A}_2(t)]\cdot\mathbf{r}}{c}\right\} \times\exp\left[-\frac{i}{2}\int^t P^2(\tau)d\tau\right] \ . \tag{2}$$

where

$$E_n(n+\tfrac{1}{2})\omega_c \quad (n=0,1,2,\dots) \ , \tag{3a}$$

$$\mathbf{p} = [p_x,\Pi_y,p_z] \ , \tag{3b}$$

$$P^2(t) = \left|\mathbf{p}+\frac{1}{c}[\mathbf{A}_1(t)+\mathbf{A}_2(t)]\right|^2 + [p_x-\Pi_x]^2 \ , \tag{3c}$$

$$\xi = \omega_c^{1/2}y - (\omega_c)^{-1/2}[p_x-\Pi_x]^2 \ , \tag{3d}$$

$$c_n = \left[\frac{\omega_c}{\pi}\right]^{1/4}\frac{1}{2\pi}[n!2^n]^{-1/2} \ , \tag{3e}$$

p_x and p_z are constant of motion, n is the quantum number identifying the Landau level, H_n is the Hermite polynomial, and $\mathbf{A}_j(t)$ $(j=1,2)$ is the vector potential of the two radiation fields. Finally, Π_x and Π_z are two real functions defined as

$$\Pi_x + i\Pi_y = -\frac{\omega_c}{c}\int^t \{[A_{1y}(\tau)+A_{2y}(\tau)] - i[A_{1x}(\tau)+A_{2x}(\tau)]\} \times\exp[i\omega_c(t-\tau)]d\tau \ . \tag{4}$$

As mentioned in the Introduction, the influence of the external fields on the initial state is neglected. Accordingly, the latter, in the zero-range approximation, may be written as [11]

$$\Psi_i^+ \approx \exp(-iI_0 t)N\frac{\exp(-br)}{r} \ , \tag{5}$$

where N is a normalization constant [11] and

$$b = \sqrt{2|I_0|} \ . \tag{6}$$

For H^-, $N=0.315\,52$ a.u. and $I_0=-0.027\,7509$ a.u. Assuming that the LF field is along z and that the uv radiation field is weak enough to be treated perturbatively, neglecting in the S matrix all the terms of order higher than those proportional to E_1, the cross section of the photodetachment by absorption of a photon ω_1 polarized along z has the form

$$(\sigma_\parallel)_{\rm LF} = \sum_{n_2}\sum_{n=0}^{n_{\max}}\left|\frac{1}{2\pi}\int_0^{2\pi}d\alpha\exp(in_2\alpha+i\lambda_2\sin\alpha + i\rho_2\sin 2\alpha) \times[p_z+K(\alpha)]\right|^2\frac{4\pi^2 I_0^2 A^2\omega_c}{c\omega_1^3 p_z} \ . \tag{7}$$

In the case in which the uv photon is polarized perpendicularly to z the cross section is

$$(\sigma_\perp)_{\rm LF} = \sum_{n_2}\sum_{n=0}^{n_{\max}} J_{n_2}^2(-\lambda_2,-\rho_2)(\sigma_\perp)_{FF} \ , \tag{8}$$

where σ_{FF} is the cross section in the absence of the LF radiation field, equal to

$$(\sigma_\perp)_{FF} = \frac{4\pi^2 I_0^2 A^2\omega_c^2}{c\omega_1(\omega_1+\omega)^2}\left[\frac{\omega_1^2+\omega_c^2}{(\omega_1-\omega)^2}n+\frac{1}{2}\right]\frac{1}{p_z} \ . \tag{9}$$

In addition,

$$\lambda_2 = (E_2/\omega_2)p_z \ , \tag{9a}$$

$$\rho_2 = \Delta_2/2\omega_2 \ , \tag{9b}$$

$$p_z = \sqrt{2[n_2\omega_2+\omega_1+I_0-\Delta_2-(n+\tfrac{1}{2})\omega_c]} \ , \tag{9c}$$

$$n_{\max} = \frac{(n_2\omega_2+\omega_1+I_0-\Delta_2)}{\omega_2} - \frac{1}{2} \ , \tag{9d}$$

$$K(\alpha) = (E_2/\omega_2)\cos\alpha \ . \tag{9e}$$

n_2 is the number of photon exchanged by the ejected electron; $J_n(x,y)$ is the generalized Bessel function [12,13]; $K(\alpha)$ is the oscillating momentum imparted to the electron by the LF field; $A=16.079\,000$. The cross sections $(\sigma_\parallel)_{\rm LF}$ and $(\sigma_\perp)_{\rm LF}$ are incoherent sums of different processes, each identified by the number n_2 of exchanged LF photons and by the Landau quantum number n. By an inspection of formula (8) one can tell that the cross section of a given channel is factored into a product in which the generalized Bessel function may be interpreted as the probability that the ejected electron exchanges n_2 photons with the LF field, while the second factor is the cross section without the LF field evaluated at the value of the momentum fixed by Eq. (9b). As $(\sigma_\perp)_{FF}$ exhibits resonances when $p_z=0$, i.e., when the opening of a Landau level is energetically possible, the presence of a LF field gives rise to replicas of such resonances. When the uv photon has a polarization parallel to **B**, the cross section is no longer factored and the presence of $\mathbf{E}_2$ produces another kind of resonance, brought about solely by $\mathbf{E}_2$. By an inspection of Eq. (7) it is seen that such resonances are originated by the fact that the final-state density is inversely proportional to the *average value* of the electron momentum along z (p_z), while the transition amplitude differs from zero even when p_z is zero because it is a coherent sum of amplitudes corresponding to processes in which the electron is ejected with instantaneous momentum $p_z+K(\alpha)$. It is worth noting that the presence of resonances in the cross sections is a result differing in an essential way from what is found when a dc electric field is present [8,14]. In the latter case, in fact, the field smears out the resonances, if these are present when there is no such field. The resonances of Eqs. (7) and (8) disappear if, more realistically, the broadening of the Landau levels is taken into account. Such broadening may be accounted for in our treatment using a complex energy for the Landau levels.

$$\varepsilon_n=(n+\tfrac{1}{2})\omega_c+i\frac{\Gamma}{2} \ . \tag{10}$$

In such a case, in the transition probability for a unit of time, the δ function, accounting for the conservation of energy, is replaced by a Lorentzian and following known procedures [15,16], cross sections are obtained differing from those reported in Eqs. (7) and (8) by the substitution of p_z with

$$p_z=\left[\frac{p_z^2+\sqrt{\Gamma^2+p_z^4}}{\Gamma^2+p^4}\right]^{-1/2} . \tag{11}$$

When this procedure is adopted, the divergences of Eqs. (7) and (8) are replaced by enhancements.

III. CALCULATIONS AND COMMENTS

In this section we report calculations of cross sections as functions of the relative energy $\omega_r=\omega_1+I_0-\omega_c/2$ of the uv photon. As in the experiment of Ref. [1], the parameters of the LF field $\mathbf{E}_2$ are chosen in such a way as to allow a large number of LF photons to be exchanged with the ejected electron. As discussed in Ref. [5], the multiphoton regime is identified by the condition $\eta=\Delta_2/\omega_2\gg 1$. In the calculations reported below, such a condition is always met. Besides, the period of the LF field is chosen to be larger than the time required for the electron to leave the ion; this latter condition allows one to consider the LF fields as quasistatic during the photodetachment. As mentioned in the Introduction, in the process under consideration, two distinct physical regimes may be identified and characterized, respectively, by values of the parameter γ larger or smaller than unity. In one regime ($\gamma>1$), the interference effects due to the LF electric field are well separated from the effects coming from the successive opening of the Landau levels. Depending on the uv photon energy, the interference effects may be absent. In the other regime ($\gamma<1$), the interference effects are always present. The reported calculations cover both regimes and the following two geometries: (a) the uv photon is polarized along **B**: (b) the uv photon is polarized perpendicularly to **B**. In both geometries the LF field $\mathbf{E}_2$ is polarized along **B**.

A. uv photon parallel geometry ($\mathbf{E}_1\|\mathbf{B}$) and $\mathbf{E}_2\|\mathbf{B}$

In Fig. 1(a) we report the cross section of photodetachment in the presence of the LF field only, while in Fig. 1(b) the cross section with magnetic field **B** in the regime $\gamma>1$ is shown. In Fig. 1(b) the cross section is also compared to that with **B**, but without the LF field $\mathbf{E}_2$. The field parameters are $I_2=10$ W/cm^2, $\omega_2=4\times10^{-5}$ eV, and $B=5\times10^5$ G. The comparison between the curve of Fig. 1(a) and the thin curve of Fig. 1(b) clearly shows that the presence of the magnetic field enhances the interference effects; in fact, more oscillations of larger amplitudes are present in the thin curve of Fig. 1(b). The most peculiar features exhibited by the curve representing the cross section with both fields **B** and $\mathbf{E}_2$ are the following. (i) The cross section displays structures that are absent in the cross section of photodetachment when only a magnetic field is present. These structures are repeated each time ω_r becomes larger than $m\omega_c$, with $m=1,2,3,\ldots$. At each repetition, the shape of these structures remains largely the same. This is a consequence of the fact that the magnetic effects are separated from those due to the LF field $\mathbf{E}_2$. (ii) For values $\omega_r<0$, the cross section does not exhibit oscillations and shows a behavior typical of tunneling processes. (iii) When $\omega_r\approx0$ the cross section has a maximum, and this value of ω_r separates the region of tunneling from that of oscillating behavior, which extends up to values of ω_r between $2\Delta_2$ and $3\Delta_2$. (iv) Finally, for even larger values of ω_r, the curve shows an increase in the average behavior, overlapping that of the cross section with the magnetic field only. As discussed in Sec. II, the enhancement, interrupting the regular behavior of the cross section, is due to the opening of the successive Landau channels.

In Fig. 2 we report the cross section with the magnetic field **B**, keeping for the LF field the same values of the parameters as in Fig. 1(b) and lowering by an order of magnitude the intensity of the magnetic field, which now is equal to 5×10^4 G. With this we now have $\gamma<1$. Examination of Fig. 2 shows that the cross section with **B** and $\mathbf{E}_2$ has an irregular behavior, as we have a mixing of the contributions coming from different Landau levels. As observed above, when $\gamma < 1$ the magnetic-field effects on photodetachment cannot be separated from those due to the LF field. The results reported in Figs. 1 and 2 may be interpreted according to the following model. In the absence of the LF field, the photodetachment has a threshold corresponding to the uv photon frequency $\omega_{th} = |I_0| + \omega_c/2$, which is the minimum value to reach the first Landau level. If ω_1 is larger than ω_{th}, the excess energy is converted into kinetic energy E_k of the free motion along **B** up to where ω_1 becomes sufficient to excite the second Landau level, and so on. The presence of the electric field $\mathbf{E}_2$ along **B** produces a barrier of height of about $2\Delta_2$ [6]. For values of $\omega_r < 0$ the electric field $\mathbf{E}_2$ makes possible the electron ejection by tunneling, while for values of $\omega_r > 0$ the interference region is entered, lasting until the kinetic energy of the free motion is smaller than or almost equal to the barrier height. Because of the interference between the direct and the reflected waves, oscillations appear in the cross section. For values of the free motion kinetic energy greater than the barrier height, the barrier effects disappear and the cross-section behavior becomes similar to the one with the magnetic field only.

To make more clear the role played by the barrier produced by the LF field, in Figs. 3 and 4 we compare the LF cross sections reported in Figs. 1(b) and 2 with those resulting by averaging the cross sections in the presence of a magnetic and a static electric field over the values taken by the LF field over a cycle (see the Appendix for the derivation). Figure 3 shows that $(\sigma_{\parallel})_{\rm LF}$, corresponding to the case $\gamma > 1$, exhibits the same behavior as $\langle \sigma_{\parallel} \rangle_{\rm dc}$, Eq. (A5), for $\omega_r < \Delta_2$, i.e., up to where the kinetic energy E_K of the free motion is much smaller than the

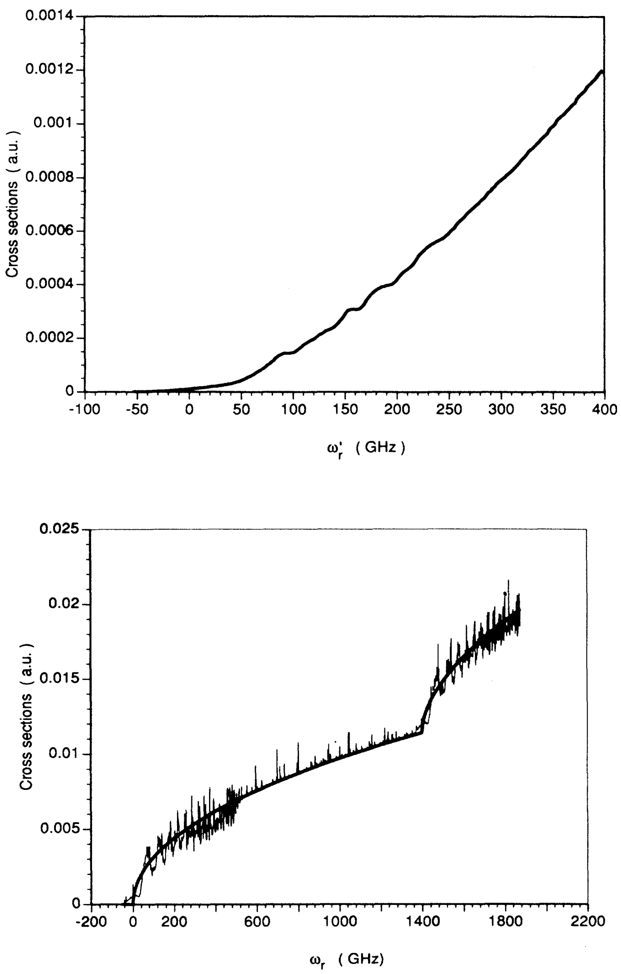

FIG. 1. (a) Photodetachment cross section in the presence of the LF field vs the relative energy $\omega_r = \omega_1 + I_0$ of the uv photon. The LF parameters are $I_2 = 10$ W/cm^2 and $\omega_1 = 4\times10^{-5}$ eV. Both the oscillating electric fields are directed along the z axis. (b) $(\sigma_{\parallel})_{\rm LF}$, Eq. (7), vs $\omega_r = \omega_1 + I_0 - \omega_c/2$ with (thin line) and without (thick line) the LF field. The parameters of the LF field are the same as in (a) and $B = 5\times10^5$ G. Both the oscillating electric fields are parallel to the static magnetic field.

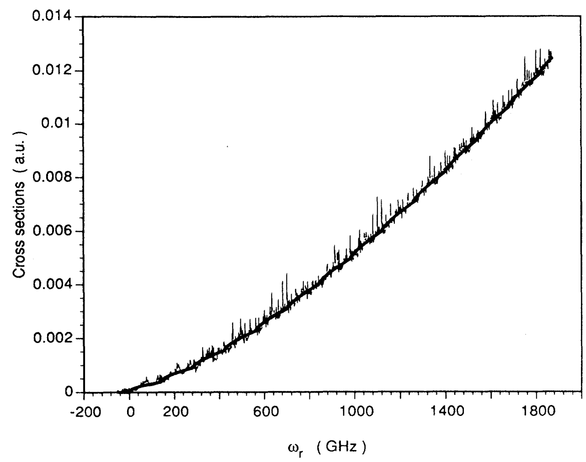

FIG. 2. Same as in Fig. 1(b) with $B = 5\times10^4$ G.

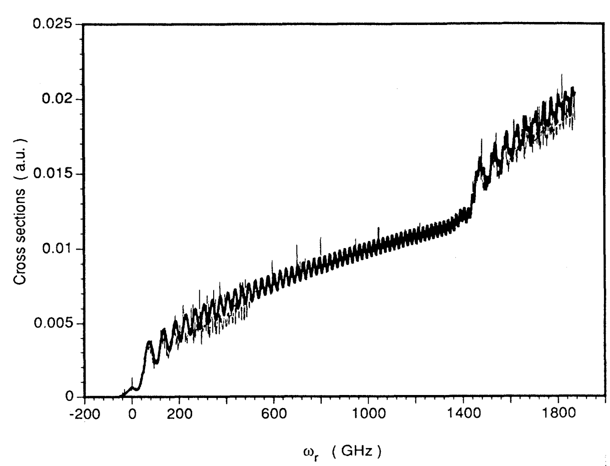

FIG. 3. $(\sigma_{\parallel})_{\rm LF}$ (thin line) and $\langle \sigma_{\parallel} \rangle_{\rm dc}$, Eq. (A5), (thick line) vs ω_r. The parameters of the LF and static magnetic fields are, respectively, the same as in Figs. 1(a) and 1(b). The field geometry is the same as in Fig. 1(b).

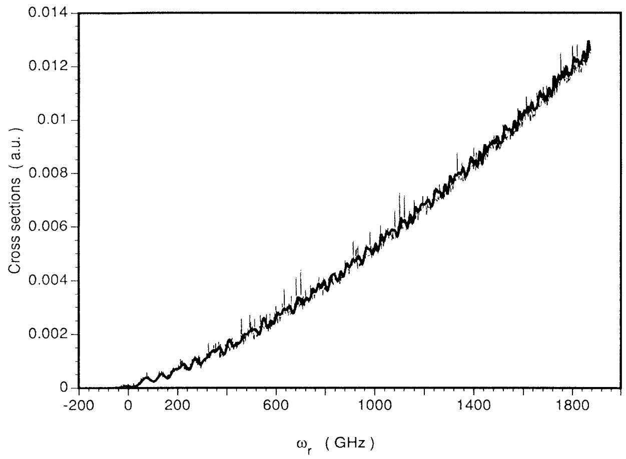

FIG. 4. Same as in Fig. 3 with $B = 5\times 10^4$ G.

height of the potential barrier. In particular, we want to point out that for $\gamma > 1$ at $\omega_r = 0$ the dominant channel of photodetachment is the one corresponding to the Landau level with $n = 0$. Hence, at the threshold, $\langle \sigma_\parallel \rangle_{dc}$ has a relative maximum because the squared derivative of the Airy function $\mathrm{Ai}(z)$ has a relative maximum for $z = 0$. Since near the threshold $(\sigma_\parallel)_{LF}$ and $\langle \sigma_\parallel \rangle_{dc}$ have the same behavior, it follows that the occurrence of the maximum at the threshold can be considered a general feature of the cross section of the process under consideration. We remark that this generality is confirmed by inspecting the curves reported in Ref. [7] (showing the photodetachment cross section in the presence of the sole LF field as a function of the electron energy E_T along the direction perpendicular to the LF field). Also in that case the curves exhibit a relative maximum when the electron longitudinal energy $\omega_1 + I_0 - E_T$ is zero. For larger values of E_k, the oscillations of the two curves go out of phase because of the finite height of the effective barrier and persist up to values of the free motion kinetic-energy approximately equaling the estimated height of the barrier. Accordingly, the quasistatic approximation is valid for values of the free motion energy smaller than Δ_2. When $\gamma < 1$, the predominant channels are the ones in which the electron is ejected with a kinetic energy along the magnetic field direction smaller than the barrier height for any energy of the uv photon giving rise to the interferences responsible for the oscillations appearing in the cross sections shown in Fig. 4.

B. uv photon perpendicular geometry ($\mathbf{E}_1 \perp \mathbf{B}$) and $\mathbf{E}_2 \parallel \mathbf{B}$

For this case as well, the LF field polarization is the same as before. The photodetachment cross sections vs ω_r are expected to have a rather different behavior as compared to the one discussed in the preceding subsection, just as it happens when the LF field is switched off or a static field is present [12]. However, as the main effect of the presence of the LF field remains that of creating a barrier of finite height in the direction of the magnetic field, it is still possible to characterize, as in the previous case, two different regimes depending on γ. In Fig. 5 we show the photodetachment cross sections vs ω_r for $\gamma > 1$. Their features indicate that the LF field-induced ripples are qualitatively similar to the ones found in the parallel geometry. Moreover, the quasistatic cross section $\langle \sigma_\perp \rangle_{dc}$ exhibits the same behavior as $(\sigma_\perp)_{LF}$ for values of the kinetic energy of the electron free motion E_k much smaller than the height of the effective potential barrier $2\Delta_2$. We note that, for $\gamma > 1$, the maximum of $(\sigma_\perp)_{LF}$ are more pronounced than the ones exhibited by $(\sigma_\parallel)_{LF}$ and are located beyond each Landau threshold, just where $(\sigma_\parallel)_{LF}$ have relative maxima. Further, the ripples induced by the LF field in the parallel geometry range over an interval of ω_r that is wider than the one found in the perpendicular geometry. This circumstance indicates that, as the kinetic energy of the electron free motion increases, the interference effects in the perpendicular geometry are less persistent than the ones in the parallel geometry. Also, we remark that for $\omega_r < \omega_c$, in the region where the cross sections oscillate, the maxima of $(\sigma_\perp)_{LF}$ are located in correspondence with the minima of $(\sigma_\parallel)_{LF}$. By exploiting the closeness of the values of $(\sigma)_{LF}$ and $\langle \sigma \rangle_{dc}$ in both the geometries investigated, in the region of oscillation, this last aspect may be easily evinced by the approximate expression of $\langle \sigma \rangle_{dc}$. In fact, by using the asymptotic expansion of the Airy function in the region where it oscillates and evaluating the integrals (A5) and (A7) by the steepest descent method (see Ref. [7]), $\langle \sigma_\parallel \rangle_{dc}$ and $\langle \sigma_\perp \rangle_{dc}$ read, respectively,

$$\langle \sigma_\parallel \rangle_{dc} \approx (\sigma_\parallel)_{FF} \left\{ 1 - \left[\frac{2}{\pi\beta} \right]^{1/2} \left[\sin(\beta + \pi/4) - \frac{1}{\beta}\cos(\beta + \pi/4) \right] \right\}, \tag{12}$$

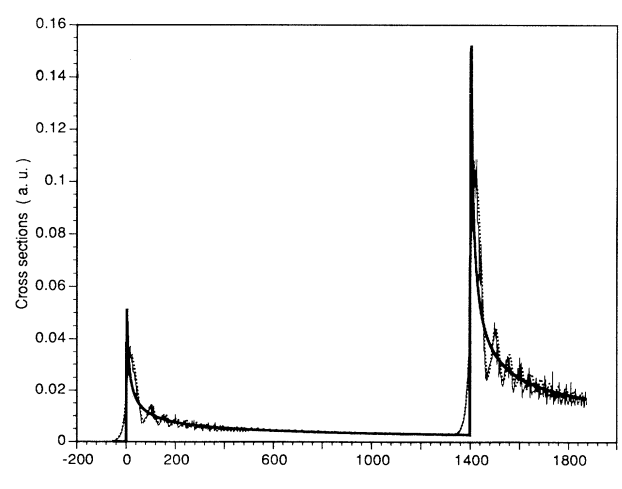

FIG. 5. $(\sigma_\perp)_{LF}$ (thin line) and $\langle \sigma_\perp \rangle_{dc}$, Eq. (A7), (dotted line) vs ω_r. The LF electric field is parallel to the static magnetic field and the uv field is perpendicular to both. The thick curve represents the photodetachment cross section without the LF field. The parameters of the LF and static magnetic fields are the same as in Fig. 3.

$$\langle \sigma_\perp \rangle_{dc} \approx (\sigma_\perp)_{FF} \left\{ 1 + \left[\frac{2}{\pi\beta} \right]^{1/2} \left[\sin(\beta+\pi/4) - \frac{1}{\beta}\cos(\beta+\pi/4) \right] \right\} \tag{13}$$

with

$$\beta = \frac{4\sqrt{2}}{3E_2}(\omega_r - n\omega_c)^{3/2} .$$

Equations (12) and (13) show that for $\beta > 1$ the effect of the LF electric field is to add oscillations to the cross section found in the absence of radiation whose amplitude is $(\sigma)_{FF}$. Moreover, whereas for one kind of geometry the interference is constructive, for the other it is destructive.

In Fig. 6 $(\sigma_\perp)_{LF}$ and $\langle \sigma_\perp \rangle_{dc}$ are shown as a function of ω_r for $\gamma < 1$. Once again, as for the parallel geometry, the effects of the magnetic and LF fields cannot be separated, giving rise to a quite irregular behavior of the cross sections.

IV. CONCLUSIONS

We have investigated the photodetachment of the negative ion H^- by an uv photon in the presence of a low-frequency electric field and a static magnetic field. The case has been treated in which the ejected electron may exchange a large number of LF photons and the LF field period is much larger than the time required for the electron ejection. The LF field polarization has been taken to be along the magnetic field direction, with the uv field polarization either along or perpendicular to **B**. The intensities of the two external fields have been taken to be sufficiently weak to neglect the Stark and diamagnetic shifts of the initial state. The reported results support the picture according to which the main role of the LF field is that of creating a repulsive barrier in the direction of the ejected electron free motion. Insofar as the uv photon energy is smaller than the threshold energy without the LF field, the detachment takes place through tunneling. When it is larger than the threshold energy, the excess energy is converted into kinetic energy of free motion E_k. If E_k is smaller than the effective height of the barrier (estimated to be about $2\Delta_2$), the quasistatic approximation for the LF field holds and the oscillations of the cross section overlap those obtained averaging the cross section in the presence of a static electric field over the values taken by the LF field in a cycle. For values of E_k approaching the barrier height, the presence of the LF field continues to give rise to interference between direct and reflected waves, producing in the cross sections oscillations out of phase as compared with those exhibited by cross sections in the presence of a quasistatic electric field. Finally, for values of E_k larger than the height of the effective barrier the interference effects disappear and the cross section overlaps that calculated with the presence of the magnetic field only. We remark that the presence of the magnetic field allows one to control the free motion energy of the ejected electron E_k, which, in turn, gives the possibility of controlling the oscillations of the cross section. In particular, one can define two regimes, characterized by values greater or smaller than one of the dimensional parameter $\gamma = \omega_c / 2\Delta_2$. When $\gamma > 1$, the cross sections exhibit the three regions commented on above, while for $\gamma < 1$, the curves showing the cross sections as a function of ω_r exhibit oscillations for any value of the uv photon energy. When the uv photon is polarized along **B** the presence of the LF field yields enhancements in the cross section, which are a peculiarity of this geometry and originate by the oscillating moment imparted to the ejected electron by the LF field. If, instead, the uv photon is polarized perpendicularly to **B**, the LF field replicates the enhancements the cross section exhibits also in its absence.

In conclusion, the reported results lend support to the physical picture explaining the oscillations observed in the experiment of Ref. [1] as due to reflections by the barrier created by the LF field. While the QS approximation holds insofar as the kinetic energy E_k associated with the motion along the LF field is considerably smaller than the height barrier, the barrier effects persist until E_k approximately reaches the value of the barrier height. This result is in agreement with that reported in Ref. [7]. We conclude by stressing that the presence of a magnetic field enhances the interference effects and that experiments such as those of Ref. [1], carried out with the additional presence of a magnetic field, could provide more precise information on the role the LF field plays in photodetachment.

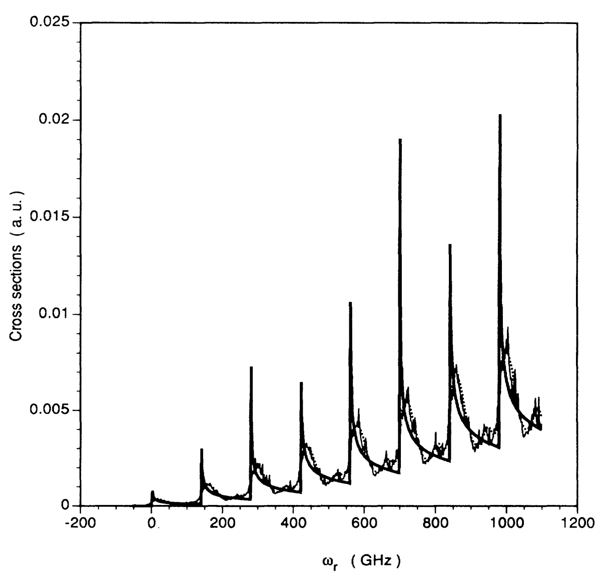

FIG. 6. Same as in Fig. 5 with $B = 5 \times 10^4$ G.

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APPENDIX

Here we outline the derivation of the cross section of photodetachment of the negative ion H^- in the presence of a quasistatic electric field E_s and a magnetic field B using Cartesian coordinates to maintain the same gauge used in the main text to derive the cross section with a LF field. The starting point is the S matrix

$$S_{fi}=-i\int_{-\infty}^{+\infty}\langle\Psi_f(r,t)|V(r)|\Psi_i^+(r,t)\rangle dt\ , \tag{A1}$$

where Ψ_i^+ is the bound-state wave function in the presence of two static fields and Ψ_f is the free electron wave function. Taking the static electric field E_s along $\mathbf{B}$ and $\omega_c<1$, Ψ_f in Cartesian coordinates is written as

$$\Psi_f=c_n\frac{1}{\sqrt{2\pi}}\exp(ip_xx)\exp\left[-\frac{\omega_c}{2}(y-y_0)^2\right]H_n[\sqrt{\omega_c}(y-y_0)] \times\left[\frac{4}{E_s}\right]^{1/6}\mathrm{Ai}\left[(2E_s)^{1/3}\left[z-\frac{E_k}{2E_s}\right]\right]\exp\{-i[(n+\tfrac{1}{2})\omega_c+\Delta_1+E_k]t\}\exp(-i\lambda_1\sin\omega_1t-i\rho_1\sin2\omega_1t) \tag{A2}$$

with $\lambda_1=(\mathbf{E}_1\mathbf{P})/\omega^2$, $\Delta_1=E_1^2/4\omega^2$; $\rho_1=\Delta_1/2\omega_1$, $c_n=(\omega_c/\pi)^{1/2}/\sqrt{2^nn!}$; and $y_0=-p_x/\omega_c$, E_k is the electron free motion energy along $\mathbf{B}$. Ai is the Airy function. Neglecting the effects of the static fields on the initial state, Ψ_i is approximated by the free field bound-state wave function Eq. (5). Keeping only terms linear in the uv field E_1, the S matrix for the uv photon polarized along B reads

$$S_{fi}=\frac{\sqrt{2\pi}}{\omega_1}E_1c_nN2^{-1/3}E_s^{1/6}\mathrm{Ai}\left[-\frac{2E_k}{(2E_s)^{2/3}}\right]\exp\left[-\frac{p_x^2}{2\omega_c}\right]H_n\left[\frac{p_x}{\sqrt{\omega_c}}\right]2\pi\delta[E_k+(n+\tfrac{1}{2})\omega_c-I_0-\omega_1)]\ . \tag{A3}$$

Proceeding in the usual way, the cross section is obtained as

$$(\sigma_\parallel)_{\mathrm{dc}}=\frac{32\pi^3}{c\omega_1^3}\frac{N^2\omega_c}{2^{2/3}}E_s^{1/3}\sum_{n=0}^{n_{\max}}\left[\mathrm{Ai}\left[-\frac{2E_k}{(2E)^{2/3}}\right]\right]^2 \tag{A4}$$

with Ai the derivative of the Airy function and $n_{\max}=(\omega_1-I_0)/\omega_c-\frac{1}{2}$. Now we change E_s into $E_2\sin\alpha$ and average Eq. (A4) over α in the interval $0<\alpha<\pi/2$:

$$\langle\sigma_\parallel\rangle_{\mathrm{dc}}=\frac{2}{\pi}\int_0^{\pi/2}(\sigma_\parallel)_{\mathrm{dc}}d\alpha\ . \tag{A5}$$

This procedure gives the QS cross section to compare to the LF cross section, Eq. (7) of the main text. If the uv photon is polarized perpendicularly to $\mathbf{B}$, the static cross section is

$$(\sigma_\perp)_{\mathrm{dc}}=\frac{8\pi^3}{c\omega_1^2}\omega_c^2N^22^{2/3}E_s^{-1/3}\sum_{n=0}^{n_{\max}}\left[\mathrm{Ai}\left[-\frac{2E_k}{(2E)^{2/3}}\right]\right]^2(n+\tfrac{1}{2}) \tag{A6}$$

and the corresponding QS cross section is obtained with the same recipe as above:

$$\langle\sigma_\perp\rangle_{\mathrm{dc}}=\frac{2}{\pi}\int_0^{\pi/2}(\sigma_\perp)_{\mathrm{dc}}d\alpha\ . \tag{A7}$$

Equation (A7) is to be compared with Eq. (8).

[1] M. C. Baruch, T. F. Gallagher, and D. J. Larson, Phys. Rev. Lett. **65**, 1336 (1990).
[2] C. Bryant, A. Mohagheghi, J. E. Stewart, J. B. Donaue, C. R. Quick, R. A. Reeder, V. Yuan, C. R. Hummer, W. W. Smith, S. Cohen, W. P. Reinhardt, and L. Overman, Phys. Rev. Lett. **58**, 2412 (1987).
[3] I. I. Fabrikant, Zh. Eksp. Teor. Fiz. **79**, 2070 (1980) [Sov. Phys. JETP **52**, 1045 (1980)].
[4] D. J. Larson, P. S. Armstrong, M. C. Baruch, T. F. Gallagher, and T. Olsen, in *Abstracts of the Fifteenth International Conference on the Physics of Electronic and Atomic Collisions, New York, 1989*, edited by A. Dalgarno, R. S. Freund, P. M. Koch, M. S. Lubell, and T. B. Lucatorto (North-Holland, Amsterdam, 1989), p. 513.

[5] S. Bivona, R. Burlon, and C. Leone, Phys. Rev. A **45**, 3268 (1992).
[6] A. Bugacov, B. Piraux, M. Pont, and R. Shakeshaft, Phys. Rev. A **45**, 3041 (1992).
[7] S. Bivona, R. Burlon, and C. Leone, Phys. Rev. A **48**, R3441 (1993).
[8] I. I. Fabrikant, Phys. Rev. A **43**, 258 (1991).
[9] See, for example, M. L. Goldberger and K. M. Watson, *Collision Theory* (Wiley, New York, 1965).
[10] S. Bivona, R. Burlon, G. Ferrante, and C. Leone, Phys. Rev. A **50**, 1984 (1994).
[11] B. H. Armstrong, Phys. Rev. **131**, 1132 (1963); M. L. Du and J. B. Delos, *ibid.* **38**, 5609 (1988).
[12] M. L. Du, Phys. Rev. A **40**, 1330 (1989).
[13] C. Leubner, Phys. Rev. A **23**, 2877 (1981).
[14] R. Reiss, Phys. Rev. A **22**, 1786 (1980).
[15] V. Canuto and J. Ventura, Fundam. Cosmic Phys. **2**, 203 (1977).
[16] S. Nuzzo and M. Zarcone, Nuovo Cimento D **13**, 741 (1991).