Electron detachment from negative ions in a bichromatic laser field

M Yu Kuchiev† and V N Ostrovsky‡

School of Physics, University of New South Wales, Sydney 2052, Australia

Received 14 January 1998

Abstract. Negative-ion detachment in a two-colour laser field is considered within the recent modification of the Keldysh model which makes it quantitatively reliable. The general approach is illustrated by calculation of angular differential detachment rates, partial rates for particular above threshold detachment channels and total detachment rates for the H\(^-\) ion in a bichromatic field with a 1:2 frequency ratio. Both perturbative and strong-field regimes are examined. Polar asymmetry and phase effects are quantitatively characterized with some new features revealed.

1. Introduction

Photoionization of atoms in a bichromatic laser field have recently received considerable attention in both theory (see, for instance, Baranova et al 1990, 1992, 1993, Baranova and Zel’dovich 1991, Szöke et al 1991, Anderson et al 1992, Schafer and Kulander 1992, Potvliege and Smith 1991, 1992a, b, 1994, Pazdzersky and Yurovsky 1994, Protopapas et al 1994, Véniard et al 1995, Pazdzersky and Usachenko 1995a, b, 1997, Fifirig et al 1997) and experiment (Muller et al 1990, Ce Chen and Elliott 1990, Baranova et al 1991, 1992, Yin et al 1992). One of the principal points of interest seems to be dependence of the observables on the difference of field phases \(\phi\), i.e. the problem of phase control. Another important aspect is the angular (polar) asymmetry of photoionization rate. These effects are interrelated since polar asymmetry is \(\phi\)-dependent and vanishes for some particular value of phase (see more detail below). The calculations have been carried out previously for the ionization of a hydrogen atom in two laser fields with a frequency ratio 1:2 (Schafer and Kulander 1992), 1:3 (Potvliege and Smith 1991) and 2:3 (Potvliege and Smith 1994). Potvliege and Smith (1992) presented results for various frequency ratios and initial states. Different schemes were employed, but all of them implied numerically intensive work.

For the multiphoton electron detachment from negative ions some analytical treatment exists (Baranova et al 1993, Pazdzersky and Yurovsky 1994, Pazdzersky and Usachenko 1995a, b, 1997) being limited mostly to the case when one or both fields are weak. The presence of a large number of parameters in the problem sometimes makes results of analytical studies not directly transparent. The case of fields with comparable (and large) intensities is also of interest bearing in mind both possible experimental realizations and the theoretical description of the transition between the multiphoton and tunnelling regimes.

† E-mail address: kuchiev@newt.phys.unsw.edu.au
‡ Permanent address: Institute of Physics, The University of St Petersburg, 198904 St Petersburg, Russia. E-mail address: Valentin.Ostrovsky@pobox.spbu.ru
The multiphoton electron detachment from negative ions presents a unique situation when quantitatively reliable results can be obtained by analytical methods in a broad range of parameters characteristic to the problem. Indeed, it has been demonstrated recently (Gribakin and Kuchiev 1997a, b) that judicious modification of the Keldysh (1964) model† ensures a very good quantitative agreement with results of numerically intensive developments, being much more simple. In many cases it also provides good agreement with available experimental data. It works unexpectedly well even for a small number \( n \) of photons absorbed. In addition to the numerous examples considered previously (Gribakin and Kuchiev 1997a, b), here we briefly comment on the most recent experiments by Zhao \textit{et al} (1997) on non-resonant excess photon detachment of negative hydrogen ions. After absorption of two photons, the electron is ejected in superposition of \( S \)- and \( D \)-waves due to selection rules. The experiment demonstrates the prevalence of \( D \)-wave contribution (90\% \pm 10\%). Our calculations give for \( D \)- and \( S \)-waves population 86.2 and 13.6\% respectively‡ for experimental conditions (light frequency \( \omega = 0.0428 \), field intensity \( I = 10^{10} \text{ W cm}^{-2} \)). The elaborate numerical calculations by Telnov and Chu (1996) and by Nikolopoulos and Lambropoulos (1997) give for the \( D \)-wave population 91 and 89\% respectively. The difference between these values is almost the same as the difference between our result and that of Nikolopoulos and Lambropoulos (1997), all three theoretical predictions lying within experimental error bars. This, together with the cases considered earlier allows us to suggest that even for \( n = 2 \) an approach (Gribakin and Kuchiev 1997a, b) provides an accuracy comparable with that of the most elaborate numerical developments.

This paper extends the approach of Gribakin and Kuchiev (1997a, b) to the case of a bichromatic field. Its objective is to provide the scheme and some representative quantitative results for two-colour electron detachment from negative ions. In particular, the phase effects and the polar asymmetry are studied. The number of parameters in the problem is quite large and at the moment they cannot be fixed experimentally. Nevertheless it seems to be worthwhile to carry out some pivoting calculations in order to obtain insight into the possible magnitude of effects specific for negative ions in bichromatic fields. We consider angular differential detachment rates, heights of above threshold detachment (ATD) peaks and total detachment rates.

2. Scheme of calculations

The previously developed scheme (Gribakin and Kuchiev 1997a, b) needs some modifications to the incorporate bichromatic problem when the electric field in the light wave

\[
\vec{F}(t) = \vec{F}_1 \cos \omega_1 t + \vec{F}_2 \cos(\omega_2 t + \varphi)
\]

is a superposition of two harmonic components with frequencies \( \omega_1, \omega_2 \) and amplitude vectors \( \vec{F}_1, \vec{F}_2 \) respectively; \( \varphi \) is the difference of field phases. Atomic units are used throughout the paper unless stated otherwise.

We consider a case of commensurable field frequencies§ which implies that the common

† Subsequent development of this model was due to Perelomov \textit{et al} (1966), Faisal (1973) and Reiss (1980); Perelomov and Popov (1967) were the first to consider multicolour process within this framework in terms of influence of higher harmonics on ionization probabilities.

‡ The model shows also some population of \( G \) and higher partial waves. However, this unphysical effect proves to be less than 0.2\% thus sustaining the model applicability.

§ The general treatment of incommensurable frequencies case was considered by Ho \textit{et al} (1983), Delone \textit{et al} (1984), Ho and Chu (1984), Manakov \textit{et al} (1986), Potvliege and Smith (1992).
Electron detachment from negative ions

period $T$ exists such that

$$T = \frac{2\pi}{\omega_1} M_1 = \frac{2\pi}{\omega_2} M_2$$

(2)

for some mutually simple integers $M_1$ and $M_2$.

The exact expression for the differential transition rate $d\omega_\lambda$ is derived following appendix A of the paper by Gribakin and Kuchiev (1997a) with the result

$$d\omega_\lambda = \frac{2\pi}{T} \sum_{\epsilon_f} |A_{\lambda\epsilon_f}|^2 \delta(E_\lambda - E_0 - \epsilon_f) \rho_\lambda,$$

(3)

$$A_{\lambda\epsilon_f} = \frac{1}{T} \int_0^T \langle \psi_\lambda(t)|V(t)|\psi_0(t) \rangle \, dt.$$  

(4)

Here $\psi_0(t) = \exp(-iE_0 t)\phi_0$ describes an initial state for the time-independent Hamiltonian $H_0$, and $\psi_\lambda(t)$ is a quasienergy state for the total Hamiltonian $H = H_0 + V(t)$, which includes interaction with the periodic field $V(t) = -e\vec{r} \cdot \vec{F}(t)$, $V(t) = V(t + T)$:

$$\frac{i}{\hbar} \frac{\partial \psi_\lambda}{\partial t} = [H_0 + V(t)]\psi_\lambda,$$

(5)

$$\psi_\lambda(t) = \exp(-iE_\lambda t) \phi_\lambda, \quad \phi_\lambda(t) = \phi_\lambda(t + T),$$

(6)

$E_\lambda$ is the quasienergy, $\rho_\lambda$ is the density of states, $\vec{r}$ is the active electron vector. The energy $\epsilon_f$ absorbed from the electromagnetic field could be presented as $\epsilon_f = n_1\omega_1 + n_2\omega_2$ with some integers $n_1$ and $n_2$. However, this representation (i.e. the choice of $n_1$ and $n_2$) is generally non-unique which reflects the existence of different coherent interfering paths (with a different number of absorbed and emitted photons at each frequency) leading to the same final state.

If the interaction of a light wave with an electron is described in the dipole-length form, as presumed above, then a long-range contribution to the matrix elements is emphasized. Therefore it is sufficient to employ an asymptotic form of the initial-state wavefunction (Gribakin and Kuchiev 1997a):

$$\phi_0(\vec{r}) \approx A r^{-\nu-1} \exp(-k r) Y_{\nu m}(\hat{r}) \quad (r \gg 1),$$

(7)

where $E_0 = -\frac{1}{2} k^2$, $\nu = Z/\kappa$, $Z$ is the charge of the atomic residual core ($\nu = Z = 0$ for a negative ion), and $\hat{r}$ is the unit vector. The coefficients $A$ are tabulated for many negative ions (Radzig and Smirnov 1985).

The amplitude $A_{\lambda\epsilon_f}$ (4) is evaluated neglecting the influence of atomic potential on the photoelectron in the final state. Further on, the integral over time in (4) is calculated within the stationary phase approximation which presumes a large magnitude of the classical action

$$S(t) = \frac{1}{2} \int_0^T (\vec{p} + \vec{k}_t)^2 \, dt' - E_0 t,$$

(8)

where $\vec{k}_t$ is the classical electron momentum due to the field

$$\vec{k}_t = e \int_0^T \vec{F}(t') \, dt'.$$

(9)

The photoelectron translational momentum $\vec{p}$ plays the role of quantum number $\lambda$ above; in particular, the quasienergy $E_\lambda = E_{\vec{p}}$,

$$E_{\vec{p}} = \frac{1}{2} \vec{p}^2 + \frac{e^2}{4\omega_1^2} F_1^2 + \frac{e^2}{4\omega_2^2} F_2^2$$

(10)

includes a contribution from the electron quiver energy due to the field.
The result of calculations of the amplitude (4) can be written as a modification of the formula (25) in the paper by Gribakin and Kuchiev (1997a):

\[
A_{\tilde{p}e_f} = -\frac{(2\pi)^2}{T A_0} \Gamma(1 + v/2) \frac{2^{v/2} \kappa^2}{\sqrt{2\pi(-iS'_v)^{v+1}}} \sum_{\mu} Y_{lm}(\hat{p}_\mu) \exp(iS_\mu) \sqrt{2\pi(S''_\mu + 1)}.
\] (11)

A corresponding expression for the detachment rate for the negative-ion case (\(v = 0\)) reads,

\[
\frac{d\omega_{ef}}{d\Omega_{\tilde{p}}} = \frac{1}{2\pi^2} \frac{1}{\tilde{p}} |A_{\tilde{p}e_f}|^2 = \frac{(2\pi)^2}{T^2} 2^{v/2} \kappa^2 \sum_{\mu} Y_{lm}(\hat{p}_\mu) \frac{2^{v/2} \kappa^2}{\sqrt{2\pi(S''_\mu)}} \left| \exp(iS_\mu) \right|^2.
\] (12)

Here the subscript \(\mu\) indicates that the function is calculated at the saddle point \(t = t_\mu\) which satisfies equation

\[
S'(t_\mu) = 0.
\] (13)

In the plane of the complex-valued time the saddle points \(t_\mu\) lie symmetrically with respect to the real axis. Summation in (11) includes the points lying in the upper half-plane (\(\text{Im} t_\mu > 0\)) with \(0 \leq \text{Re} t_\mu \leq T\); \(\hat{p}_\mu\) is a unit vector in the direction of \(\tilde{p} + \hat{k}_\mu\). The magnitude of the final electron translational momentum \(p\) is governed by energy conservation

\[
\frac{1}{2} \kappa^2 + E_{\tilde{p}} = \epsilon_f,
\] (14)

which ensures that the momentum is real for open ATD channels.

According to (8), (9), (1) we have

\[
S'(t) = \frac{1}{2}(\tilde{p} + \hat{k}_t)^2 + \frac{1}{2} \kappa^2 = \frac{1}{2} \tilde{p}^2 + \frac{\epsilon^2}{2\omega_1^2} F_1^2 \sin^2 \omega_1 t + \frac{\epsilon^2}{2\omega_2^2} F_2^2 \sin^2(\omega_2 t + \varphi)
\]

\[
+ \tilde{p} \cdot \tilde{F}_1 \frac{\epsilon}{\omega_1} \sin \omega_1 t + \tilde{p} \cdot \tilde{F}_2 \frac{\epsilon}{\omega_2} \sin(\omega_2 t + \varphi)
\]

\[
+ \tilde{F}_1 \cdot \tilde{F}_2 \frac{\epsilon^2}{\omega_1 \omega_2} \sin \omega_1 t \sin(\omega_2 t + \varphi) + \frac{1}{2} \kappa^2.
\] (15)

The frequencies \(\omega_1\) and \(\omega_2\) are integer multiples of the basic frequency \(\omega = 2\pi/T\)

\[
\omega_1 = M_1 \omega, \quad \omega_2 = M_2 \omega.
\] (16)

Assuming for definiteness that \(M_2 > M_1\) and introducing \(\zeta = \exp(i\omega t)\) we notice that the function

\[
\mathcal{P}(\zeta) = \zeta^{2M_2} S'(\zeta)
\] (17)

is a polynomial of the power \(4M_2\) in \(\zeta\). This observation is of practical importance since equation (13) defining the saddle point is equivalent to

\[
\mathcal{P}(\zeta) = 0.
\] (18)

The efficient numerical procedures for finding roots of polynomials are available, and, in particular, one can be confident that all the roots are found.

The practical calculations are conveniently carried out using the Mathematica program (Wolfram 1991). Starting from the expression for \(S'(t)\) one derives \(S(t)\) and \(S''(t)\) by analytical integration and differentiation respectively. The saddle points are found using equation (18), and the roots \(t_\mu\) lying in the upper half-plane are selected. Finite summation over \(\mu\) in (11) or (12) completes the calculation.

The roots \(t_\mu\) and hence the photoionization amplitude (11) and rate (12) depend on the orientation of electron translational momentum \(\tilde{p}\) with respect to the field amplitudes \(\tilde{F}_1\) and \(\tilde{F}_2\). It is not difficult to consider the fields of various relative orientation and polarization,
but for simplicity we limit our further calculations to linear polarized fields with $\vec{F}_1 \parallel \vec{F}_2$. Then the differential photoionization rate depends only on the single angle $\theta$ between the vectors $\vec{p}$ and $\vec{F}_1 \parallel \vec{F}_2$.

3. Results

Our calculations for $\text{H}^-\text{ detachment}$ are carried out for the parameters $\kappa = 0.2354$, $A = 0.75$ (Radzig and Smirnov 1985). The frequencies ratio $\omega_1/\omega_2 = 1:2$ is considered. In this case the field (1) has zero mean value, but possesses polar asymmetry (i.e. asymmetry under inversion of the $z$-axis directed along $\vec{F}_1 \parallel \vec{F}_2$) which could be conveniently characterized by the time-average value (Baranova et al 1990)

$$\langle F^3 \rangle = \frac{3}{4} F_1^2 F_2 \cos \varphi. \quad (19)$$

Presuming that $F_1, F_2 > 0$, from this expression one infers, for instance, that for the phase $\varphi \in [0, \frac{\pi}{2}]$ the electric field effectively attains larger values in the positive-$z$ direction than in the negative-$z$ one. This is illustrated, for example, by figure 1 in the paper by Schafer and Kulander (1992), or by figure 2 in the paper by Baranova et al (1993). Note that our definition of the phase $\varphi$ is the same as in the papers by Baranova et al (1990), Muller et al (1990), Pazdzersky and Yurovsky (1995), but differs from that chosen by Schafer and Kulander (1992) who describe the electric field as $\vec{F}(t) = \vec{F}_1 \sin \omega_1 t + \vec{F}_2 \sin (\omega_2 t + \varphi_{KS})$.

Namely, the phases are related by $\varphi_{KS} = \varphi - \frac{\pi}{4}$.

Although the formula (19) shows that the field has polar symmetry for $\varphi = \frac{\pi}{2}$ and maximal polar asymmetry for $\varphi = 0$, quite paradoxically, the differential detachment rate (12) possesses polar symmetry for $\varphi = 0$ (i.e. is invariant under the transformation $\theta \Rightarrow \pi - \theta$), and is asymmetrical for other values of phase (see discussion by Baranova et al (1990), Schafer and Kulander (1992), Pazdzersky and Yurovsky (1995)).

The other features of the phase effects are as follows.

(i) The system Hamiltonian is a $2\pi$-periodic function of the parameter $\varphi$.

(ii) The system Hamiltonian is not changed by simultaneous transformation $\varphi \Rightarrow -\varphi$, $\theta \Rightarrow \pi - \theta$. The same applies to the differential detachment rate (12).

(iii) The transformation $\varphi \Rightarrow \pi - \varphi$ leaves the Hamiltonian invariant only if $t$ is replaced by $-t$.

As stressed by Baranova et al (1990), Baranova and Zeldovich (1991), Anderson et al (1992), Baranova et al (1993), the problem is invariant under the time inversion operation provided the final-state electron interaction with the atomic core is neglected. This is the case in the present model. Thereby our differential ionization rates are the same for $\varphi$ and $(\pi - \varphi)$; hence it is sufficient for us to consider phases only from the interval $\varphi \in [0, \frac{\pi}{4}]$.

The calculations by Baranova et al (1990) within the perturbation theory and by Schafer and Kulander (1992) within the wave propagation technique took into account the final state electron–core interaction. Therefore they have found some deviations from the symmetry under $\varphi \Rightarrow (\pi - \varphi)$ transformation. However, for the negative-ion photodetachment this effect could be anticipated to have only a minor influence.

At first we consider two fields with frequencies $\omega = 0.0043$ and $2\omega$ and equal intensities $I_1 = I_2 = 10^9$ W cm$^{-2}$. It is well known that the regime of detachment process is governed by the Keldysh parameter $\gamma = \omega \kappa / F$ ($\gamma \gg 1$ for multiphoton detachment in perturbative regime; $\gamma \ll 1$ for strong field, or tunnelling regime). In the present case for the first field we have $\gamma_1 = \omega_1 \kappa / F_1 = 6$, and for the second field $\gamma_2 = 2\gamma_1$, which corresponds to the multiphoton regime.
Figure 1. Detachment of the H\(^{−}\) ion in a bichromatic field with the frequencies \(\omega = 0.0043\) and \(2\omega\) and equal intensities \(I_1 = I_2 = 10^9\) W cm\(^{-2}\). Differential detachment rate (see equation (20)) (in units \(10^{-12}\) au) as a function of the angle \(\theta\) is shown for the first (a), absorption of \(n = 7\) photons of frequency \(\omega\) and second (b, \(n = 8\)) ATD peaks and various values of the field phase difference: full curve—\(\varphi = 0\); short-broken curve—\(\varphi = \frac{\pi}{4}\); chain curve—\(\varphi = \frac{\pi}{2}\); dotted curve—\(\varphi = \frac{3\pi}{4}\); long-broken curve—\(\varphi = \frac{\pi}{2}\).

Figure 1 (as well as figures 2–4 below) shows the differential detachment rate as a function of the angle \(\theta\). The abscissas of the plots give a magnitude of

\[
\frac{1}{2} \frac{1}{2\pi} \frac{d\omega_{e_i}}{d\cos \theta} = \frac{1}{(2\pi)^2} p |A_{\vec{r}_e}|^2, \tag{20}
\]

where the right-hand side was calculated using the right-hand side of formula (12). The left-hand side of equation (20) has the factor \(\frac{1}{2}\). This means that the true detachment rate
in the case of the H$^-$ ion is twice as large as that given by equation (12). This accounts for the two possible spin states of the residual H atom (i.e. for the presence of two equivalent electrons in H$^-$).

In figure 1 we show the differential detachment rate for the first and second ATD peaks which correspond to absorption of $n = 7$ and $n = 8$ photons of frequency $\omega$ respectively. In figure 2 the same results are shown but for double the value of the field amplitude $F_2$. In figure 3 the amplitude $F_2$ is the same as in figure 1, but the amplitude $F_1$ is doubled.

In all cases the angular distributions exhibit strong dependence on the field phase difference $\varphi$. This is well expected since the angular patterns are formed by interference of contributions coming from different complex-valued moments of time $t_\mu$. For $\varphi = 0$ the

Figure 2. Same as in figure 1, but for unequal field intensities $I_1 = 10^6$ W cm$^{-2}$, $I_1 = 4 \times 10^6$ W cm$^{-2}$. The differential detachment rates is shown in units $10^{-10}$ au.
distribution is rather flat, with $\varphi$ increasing it becomes more oscillatory. An interesting and not obvious feature is that for $\varphi = \frac{1}{2} \pi$ the rate goes to zero at the values of angle $\theta$ where it has minimum.

In figure 4 we present the results for the same frequencies as before and equal field intensities $I_1 = I_2 = 10^{13}$ W cm$^{-2}$. Here the Keldysh parameter for the first field is $\gamma_1 = \omega_1 \kappa / F_1 = 0.6$. Bearing in mind the presence of the second field one can suppose that the situation corresponds to the onset of strong-field domain. The first open ATD channel corresponds to absorption of $n = 18$ photons of frequency $\omega$. The patterns in differential rates become more oscillatory than in the weak filed case.

The partial detachment rate for each ATD channel is obtained by integration of (12)
Electron detachment from negative ions

Figure 4. Same as in figure 1, but in the strong field regime: \( I_1 = I_2 = 10^{11} \text{ W cm}^{-2} \). The detachment rate for the first (a, absorption of \( n = 18 \) photons of frequency \( \omega \)) and second (b, \( n = 19 \)) ATD peaks is shown in units \( 10^{-6} \text{ au} \)

over angles

\[
    w_{\psi_j} = \int \frac{dw_{\psi_j}}{d\Omega_{\vec{p}}} d\Omega_{\vec{p}} = \int_{\theta=\pi}^{\theta=0} \frac{dw_{\psi_j}}{d\cos \theta} d\cos \theta. \tag{21}
\]

We present separately the result \( w_{\psi_j}^{(u)} \) of integration over the upper half-space of the electron ejection (\( 0 < \theta < \frac{\pi}{2} \)) and its counterpart \( w_{\psi_j}^{(l)} \) for the lower half-space (\( \frac{\pi}{2} < \theta < \pi \)). These magnitudes provide a bulk characterization for the partial rate polar asymmetry. As discussed above, the polar asymmetry disappears (i.e. \( w_N^{(u)} = w_N^{(l)} \)) for \( \varphi = 0 \) (the open
Figure 5. Partial detachment rates for various ATD channels for the H⁻ ion in bichromatic field with the same parameters as in figure 1 (perturbative regime) and various values of the field phase difference ϕ. N labels ATD peaks with the lowest N = 1 peak corresponding to absorption of seven photons of frequency ω = 0.0043. The rate w_N^{(a)} of electron emission in the upper half-space is shown by circles, its counterpart w_N^{(l)} for the lower half-space is depicted by triangles. The plot for ϕ = 1/8π additionally includes the rate for ϕ = 0 (crosses) when emission is polar symmetrical (w_N^{(a)} = w_N^{(l)}). The symbols are joined by lines to help the eye.

ATD channels are labelled below by the number N = 1, 2, ... in the order of increasing emitted electron energy. In the perturbative regime (figure 5) for the same conditions as in figure 1 the bulk polar asymmetry parameter a_N(ϕ) = w_N^{(a)}/w_N^{(l)} for the lowest ATD peak (N = 1) grows monotonously from 1 for ϕ = 0 to 2.27 for ϕ = 1/2π. For higher ATD peaks the asymmetry effect is much more prominent, for instance, for N = 5 asymmetry parameter a_5(ϕ) varies from 1 for ϕ = 0 to 63.1 for ϕ = 1/2π.

The partial detachment rates integrated over all ejection angles w_N = w_N^{(a)} + w_N^{(l)} are shown in figure 6 for three representative values of ϕ. This magnitude manifests low sensitivity to the variation of phase. For example, for N = 1 it changes by about 4% when the phase ϕ varies from 0 to 1/2π.

In the tunnelling regime (figure 7, field parameters as in figure 4) the bulk polar asymmetry generally is somewhat less substantial. For instance, for the lowest ATD peak we obtain a_{N=1}(ϕ = 1/2π) = 1.18 and for N = 5 a_{N=5}(ϕ = 1/2π) = 2.76. Note that in both figures 5 and 7 the electron emission in the upper half-space (0 < θ < 1/2π) is more probable for all N, i.e. w_N^{(a)} > w_N^{(l)} in the interval of phases considered (ϕ ∈ [0, 1/2π]). As discussed above (see the property (ii)), the situation is reversed for ϕ ∈ [−1/2π, 0]. For small value of the phase ϕ = 1/2π in figure 7 there is a clear tendency to swap the relation between w_N^{(a)} and w_N^{(l)} for higher values of N which is prevented by a kind of ‘pseudocrossing’. For
Electron detachment from negative ions

Figure 6. Same as in figure 5, but for the detachment rate integrated over all angles \( w_N = w_N^{(o)} + w_N^{(l)} \). The results for three values of the phase \( \varphi \) are shown: circles—\( \varphi = 0 \); blocks—\( \varphi = \frac{1}{4} \pi \), triangles—\( \varphi = \frac{1}{2} \pi \).

Figure 7. Same as in figure 5, but for the field parameters chosen as in figure 4 (strong-field regime). The lowest \( N = 1 \) peak corresponds to absorption of 18 photons of frequency \( \omega = 0.0043 \). The detachment rate is shown in units \( 10^{-6} \) au.

larger \( \varphi \) the partial rates \( w_N^{(l)} \) are strongly suppressed when \( N \) increases (similar behaviour is also observed in the perturbative regime). The phase effects in the partial rates \( w_N \) are
Figure 8. Same as in figure 6, but for the field parameters chosen as in figure 7 (strong-field regime). The detachment rate is shown in units $10^{-6}$ au.

Table 1. Total rates $w$ (summed over all ATD channels) for detachment of the $\text{H}^-$ ion in a bichromatic field with the frequencies $\omega = 0.0043$ and $2\omega$, equal intensities $I_1 = I_2$ and some representative values of the phase difference $\varphi$. The detachment rates $w^{(u)}$ and $w^{(l)}$ for electron ejection into the upper and lower half-spaces respectively are also shown.

| $\varphi$ | $w$ (in units $10^{-11}$ au) | $w^{(u)}$ (in units $10^{-6}$ au) | $w^{(l)}$ (in units $10^{-6}$ au) | $w^{(l)}$ (in units $10^{-6}$ au) |
|------------|-------------------------------|---------------------------------|---------------------------------|---------------------------------|
| 0          | 8.89 4.45                     | 4.45                            | 4.45                            | 163.6 81.8                       |
| $\frac{\pi}{4}$ | 8.91 5.15                     | 3.76                            | 3.76                            | 167.2 102.3                      |
| $\frac{\pi}{2}$ | 8.95 5.79                     | 3.16                            | 3.16                            | 169.3 117.9                      |
| $\frac{3\pi}{4}$ | 8.98 6.24                     | 2.73                            | 2.73                            | 172.3 130.9                      |
| $\pi$      | 8.99 6.41                     | 2.58                            | 2.58                            | 173.0 136.0                      |

less significant in the tunnelling regime (figure 8).

The total detachment rates are obtained by summation over all open ATD channels:

$$w = \sum_N w_N. \quad (22)$$

The results for the total rates as well as $w^{(l,u)} = \sum_N w^{(l,u)}_N$ are presented in table 1.

In the perturbative regime the asymmetry parameter $w^{(l)}/w^{(u)}$ for $\varphi = \frac{\pi}{2}$ is equal to 2.49 which is close to the values 2.27 obtained above for the lowest $N = 1$ ATD peak. This is easily understandable, since in the perturbative regime the lowest peak gives a dominant contribution to the total rate. In the strong-field regime the situation is less straightforward since several ATD peaks give a comparable contribution to the total rate (figure 8). The asymmetry parameter is $w^{(l)}/w^{(u)} = 3.46$ for $\varphi = \frac{3\pi}{2}$. Interestingly, the total rate $w$ is practically insensitive to the phase change both in the perturbative regime (where it varies only within 1%) and in the strong field regime (variation within 2%). In the latter case the partial rates $w_N$ (figure 8) exhibit some oscillatory structure as functions of $N$ with the positions of extrema depending on the phase $\varphi$. This $\varphi$-dependence almost completely disappears after summation over $N$ as table 1 shows.
4. Conclusion

In summary, the approach of Gribakin and Kuchiev (1997a, b) provides convenient tool for investigating two-colour photodetachment of negative ions. The bichromatic electron detachment for $H^-$ ion in the fields with 1:2 frequency ratio is examined in the perturbative and tunnelling regimes. Polar asymmetry and phase effects are quantitatively characterized and some new features revealed. It should be noted that via the recoil mechanism the predicted effect also leads to acceleration of the core thus creating anisotropic flux of neutral $H$ atoms.

Acknowledgments

We appreciate fruitful discussion with G F Gribakin. MYuK thanks the Australian Research Council for support. This work was supported by the Australian Bilateral Science and Technology Collaboration Program. VNO acknowledges the hospitality of the staff of the School of Physics of UNSW where this work was carried out.

References

Anderson D Z, Baranova N B, Green K and Zel’dovich B Ya 1992 Zh. Eksp. Teor. Fiz. 102 397–405 (Engl. transl. 1992 Sov. Phys.–JETP 75 210–14)
Baranova N B, Beterov I M, Zel’dovich B Ya, Ryabtsev I I, Chudinov A N and Shul’ginov A A 1992 Pis. Zh. Eksp. Teor. Fiz. 55 431–4 (Engl. transl. 1992 JETP Lett. 55 439–44)
Baranova N B and Zel’dovich B Ya 1991 J. Opt. Soc. Am. B 8 27–32
Baranova N B, Zel’dovich B Ya, Chu’dinov A N and Shul’ginov A A 1990 Zh. Eksp. Teor. Fiz. 98 1857–68 (Engl. transl. 1990 Sov. Phys.–JETP 71 1043–9)
Ce Chen and Elliott D S 1990 Phys. Rev. A 48 1497–505
Delone N B, Beterov I M, Zel’dovich B Ya, Ryabtsev I I, Chudinov A N and Shul’ginov A A 1992 Pis. Zh. Eksp. Teor. Fiz. 55 431–5 (Engl. transl. 1992 JETP Lett. 55 439–44)
Baranova N B, Reiss H R and Zel’dovich B Ya 1993 Phys. Rev. A 48 1497–505
——1997b J. Phys. B: At. Mol. Opt. Phys. 30 L657–64
Ho T S and Chu S I 1984 J. Phys. B: At. Mol. Opt. Phys. 17 2101–28
——1985 J. Phys. B: At. Mol. Opt. Phys. 18 2197–2207
——1990 Sov. Phys.–JETP 71 1043–9
Keldysh L V 1964 Zh. Eksp. Teor. Fiz. 47 1945–57 (Engl. transl. 1965 Sov. Phys.–JETP 20 1307–14)
Manakov N L, Ovsianikov V D and Rapoport L P 1986 Phys. Rep. 141 319–433
Muller H G, Bucksbaum P H, Schumacher D W and Zav’ev A 1990 J. Phys. B: At. Mol. Opt. Phys. 23 2761–9
Nikolopoulos L A A and Lambropoulos P 1997 Phys. Rev. A 56 3106–15
Nikolopoulos L A A and Lambropoulos P 1997 Phys. Rev. A 56 3106–15
Pazdzersky V A and Yurovsky V A 1994 Phys. Rev. A 51 632–40
——1995 Laser Phys. 5 1141
Manakov N L, Ovsianikov V D and Rapoport L P 1986 Phys. Rep. 141 319–433
Muller H G, Bucksbaum P H, Schumacher D W and Zav’ev A 1990 J. Phys. B: At. Mol. Opt. Phys. 23 2761–9
Pazdzersky V A and Yurovsky V A 1994 Phys. Rev. A 51 632–40
——1995 Laser Phys. 5 1141
——1997 J. Phys. B: At. Mol. Opt. Phys. 30 3387–402
Perelomov A M and Popov V S 1967 Zh. Eksp. Teor. Fiz. 52 514–26 (Engl. transl. 1967 Sov. Phys.–JETP 25 336–43)
——1968 J. Phys. B: At. Mol. Opt. Phys. 21 L297–310
Reiss H R 1980 Phys. Rev. A 22 1786–813
Schafer K J and Kulander K C 1992 Phys. Rev. A 45 8026–33
Szőke A, Kulander K C and Bardsley J N 1991 J. Phys. B: At. Mol. Opt. Phys. 24 3165–71
Telnov D A and Chu S I 1996 J. Phys. B: At. Mol. Opt. Phys. 29 4401–10
Véniard V, Taleb R and Maquet A 1995 Phys. Rev. Lett. 74 4161–4
Wolfram S 1991 Mathematica: A System for Doing Mathematics by Computer 2nd edn (Palo Alto: Addison-Wesley)
Yin Y-Y, Ce Chen and Elliott D S 1992 Phys. Rev. Lett. 69 2353–6
Zhao X M, Gulley M S, Bryant H C, Strauss C E M, Funk D J, Stintz A, Rislove D C, Kyrala G A, Ingalls W B and Miller W A 1997 Phys. Rev. Lett. 78 1656–9