Ion Collisions in Very Strong Electric Fields

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Abstract

A Classical Trajectory Monte Carlo (CTMC) simulation has been made of processes of charge exchange and ionization between an hydrogen atom and fully stripped ions embedded in very strong static electric fields \(O(10^{10} \text{ V/m})\), which are thought to exist in cosmic and laser–produced plasmas. Calculations show that the presence of the field affects absolute values of the cross sections, enhancing ionization and reducing charge exchange. Moreover, the overall effect depends upon the relative orientation between the field and the nuclear motion. Other features of a null-field situation, such as scaling laws, are revisited.

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I. Introduction

The study of the effects of externally imposed electric or magnetic fields on atomic systems has been since a long time an active area of research in physics, both when dealing with isolated atoms (Stark and Zeeman effects, photoionization, chaotic spectra, modifications of energy levels and wave, see e.g. Seipp and Taylor 1994, Wang and Greene 1994, Delande and Gay 1986, Friedrich 1990, Zang and Rustgi 1994, Delande et al 1994, Bylicki et al 1994, Buchleitner et al 1994), as when the attention is focused onto many-particle processes (electron–atom or atom–atom collisions).

In order to influence the electronic dynamics it is necessary that the force exerted by the field is comparable with the Coulomb intra-atomic forces, which are quite strong if the atom is not in a highly excited state: for magnetic and electric fields the natural units of measure are $B_0 = \frac{m_e^2 e^3 c}{\hbar^3} = 2.35 \cdot 10^5 \, \text{T}$, and $F_0 = \frac{m_e^2 e^5}{\hbar^4} = 5.14 \cdot 10^{11} \, \text{V/m}$. Fields of such strengths may be generated only under rather exotic situations: in white dwarfs magnetic fields of $10^5 \, \text{T}$, and in neutron stars up to $10^9 \, \text{T}$, have been discovered (Ruder et al 1994). The huge unipolar induction currents which generate the magnetic field are sustained by differences of potential of order $O(10^{13} \, \text{V/m})$ (Ruder et al 1994). By a different mechanism, electric fields may be generated even independently by magnetic fields: it is the double layer (DL) phenomenon which is expected to be found, for example, on the surface of neutron stars (Williams et al 1986, Raadu 1989). A DL is a local discontinuity surface consisting microscopically of two oppositely charged layers of plasma. The resultant electric field within this region is very strong and acts to re-establish a neutrality condition which is violated, for example, as a consequence of temperature gradients on the surface of the star or because of charge density gradients produced by the different rates at which electrons and ions, falling on the star surface from the outer space, are decelerated (Williams et al 1986). It is thought that fields $> 10^{10} \, \text{V/m}$ may be generated through this mechanism.

DLs of considerable strength are created also in Earth laboratories during laser experiments devoted, for example, to inertial confinement fusion studies: in these experiments a small target of hydrogen isotopes is irradiated with high power ($\approx 10^{17} \div 10^{20} \, \text{W/m}^2$) laser pulses. The sudden heating, which converts into a hot plasma most of the target, drives a fast dilatation of the gaseous outer corona where, because of their higher velocity, the electrons...
will tend to lead the expansion. The result is a DL which slows down the electrons and accelerates the ions (for a review of the subject see Eliezer and Hora 1989). Numerical simulations show that, inside these DLs which have a width of about $10^{-7} \div 10^{-5}$ m, electric fields of $10^8 \div 10^{11}$ V/m may be reached (Eliezer and Hora 1989, Eliezer et al 1988).

Charged particles within a DL usually have great kinetic energies because of the heating, the electrostatic force, and—in stars—the gravitational attraction: they may be $O(10\text{KeV}/\text{amu})$ (Williams et al 1986, Eliezer and Hora 1989). DLs thus provide us with rather extreme conditions where to examine interatomic processes, particularly collisions. In this paper we aim to perform a numerical investigation of the processes of electron capture and ionization between a neutral atom and a bare ion. This kind of study is not new in literature: a similar work has been performed by Grosdanov and McDowell (1985) and McDowell and Zardona (1985) for very strong, astrophysical magnetic fields. Electric fields have been instead studied by Olson and MacKellar (1981). There, however, the attention was focused on fields attainable in ordinary laboratory conditions, i.e. much weaker than those considered here; to compensate for the smallness of the field only Rydberg atoms had to be considered. Here we carry on the analysis on low lying electronic states instead that highly excited ones; some topics not covered by Olson and MacKellar (1981) are touched: the results of the scattering process, expressed in terms of cross sections, are studied with respect to the charge of the incident particle and the initial state of the system. Also, particular attention is paid to the geometrical features (direction and spatial extension) of the external field.

II. Theory

The collision process has been studied using the CTMC method, in which both nuclei and the electron are considered classical particles obeying Newton laws but the initial conditions of the electron are randomly chosen within a distribution which simulates the quantum mechanical behavior. The study of the process is then reduced to numerically integrate a set of coupled ordinary differential equations.

The CTMC method, developed by Abrines and Percival (1966a), has been frequently used in studies of atomic scattering, being quite easy to implement and rather accurate in results. The method works at its best in the region
of intermediate impact energies, i.e. when the relative impact internuclear speed \( v_p \) is greater than the classical orbital electron velocity \( v_e \): \( v_p \geq v_e \); in this work we will consider the range \( 1 \leq v_p/v_e \leq 3.5 \) (corresponding to energies \( E \leq 300 \text{ KeV/amu} \)), which is quite in accordance with the velocities expected to be found in DLs.

The equations of motion for each of the particles are derived from the Hamiltonian

\[
H = \frac{3}{2} \sum_{i=1}^{3} \frac{\vec{p}_i^2}{2} + \sum_{i<j=1}^{3} \frac{Z_i Z_j}{|\vec{q}_i - \vec{q}_j|} + \sum_{i=1}^{3} Z_i \vec{q}_i \cdot \vec{F},
\]  

(1)

where \( \vec{q}_i, \vec{p}_i \ (i = 1, 2, 3) \) are the coordinates and the momenta conjugate for the two nuclei and the electron; \( Z_i \) the charge of the \( i \)-th particle, and \( \vec{F} \) the electric field (here and in the following, for the sake of easiness, the electric field will be expressed in atomic units; in these units \( F_0 \) of section one is equal to 1).

An atom embedded in a strong electric field is an unstable system: classically, a bound electron is field–ionized if the strength \( F \) of the field is greater than \( E_b^2/4 \), with \( E_b \) binding energy. Thus, any classical calculation loses meanings at strengths beyond this value (for an hydrogen atom in the ground state, \( E_b^2/4 = F_0/16 = 0.0625 \text{ au} \)). However, quantum mechanically, there is not a ionization threshold, instead all bound states turn into resonances of finite lifetime, whatever small \( F \) be. Several calculations of lifetimes exist in literature; here we shall refer to Damburg and Kolosov (1976): in that work it is estimated that, in a field of \( F = 0.04 \text{ au} \) the mean life of an electron in the ground state is \( \tau_e \approx 2 \cdot 10^{-11} \text{ s} \). A particle moving at a speed \( v_p \geq 10^6 \text{ m/s} \) (corresponding to a kinetic energy \( \geq 25 \text{ KeV/amu} \) travels during this time a distance \( l_p = v_p \tau_e \geq 10^{-5} \text{ m} \). If we refer, for example, to a laser–produced DL, its characteristic spatial length is given by the Debye length \( \lambda_D \); usually, this value is \( \leq 10^{-5} \text{ m} \) (see Eliezer and Hora 1989, sect. 4).

Further, in a laser–produced DL the electric field is not at all static, the timescale needed to establish it and over which its variations are not negligible being in the range of \( 10^{-13} \div 10^{-12} \text{ s} \) (Eliezer and Hora 1989, sects. 4,5). Even within such short times, a particle at the speed \( v_p \) travels for several hundreds of \( \AA \), which is well beyond the typical length over which charge transfer processes take place. So we see that, even under rather extreme conditions, a classical description of the collision process does not lose its meaning (we remark however that for cosmic DLs the picture is not so satisfactory, because
the width of the DL is supposed to be greater, thus in this case the validity of our hypotheses may be questionable).

Two further remarks need to be made when dealing with excited atoms (we will consider only \( n = 2 \)): first, the validity of the classical approximation still holds because we shall consider weaker fields, \textit{ad hoc} rescaled to compensate for the smaller binding energy. Further, because of the Stark mixing between levels due to \( F \), and in particular between the levels \( 2s - 2p_0 \), all the electrons are allowed to decay by spontaneous emission of radiation into the ground state within short times (\( \approx 10^{-9} \) s) which, however, are always longer than the timescales of the collision process.

The electric field exerts also an influence on electronic energies (Stark shift) and wave functions. As far as the first are concerned, we see from Damburg and Kolosov (1976) that the correction to the null-field values is about 1 percent and thus may be neglected in the following.

Using a bit more care instead is needed when dealing with position and momentum distributions; in CTMC method quantum wave functions are replaced with a random sampling of the initial coordinates from statistical distributions. The older and more used choice is the microcanonical distribution (Abrines and Percival 1966b), where momentum \( \vec{p}_e \) and position \( \vec{q}_e \) are picked up from an uniform distribution subject to the bound that the total energy is equal to its quantum mechanical value,

\[
E = \frac{|\vec{p}_e|^2}{2} - \frac{1}{|\vec{q}_e|} = -\frac{1}{2n^2} \tag{2}
\]

This method provides a statistical momentum distribution which is equal to the quantum mechanical one, while the radial distribution is rather bad, showing a cut–off beyond \( r = |\vec{q}_e| = 2n^2 \text{ au} \). This causes a wrong estimate of total cross section at low impact energies, underestimating the contribution from collisions at large impact parameters. Several methods have been devised to overwhelm this deficiency, with good results (see Hardie and Olson 1983, Cohen 1985, and also Eichenauer et al. 1981, Kunc 1988, Reinhold and Falcon 1988a,b, Schmidt et al. 1990 for a general discussion over the CTMC and other classical approximations). However, any classical method which attempts to simulate a correct radial quantum distribution has to give up the assumption that electron energy is a well fixed quantity. The electron will thus have some probability to be given an energy beyond the ionization threshold when the field is turned on, and to be field ionized. This spurious
contribution will add to the correct probability of ionization for ionic impact. As an example, in Cohen's model (Cohen 1985), with a field strength of \( F_0/2 \), about 30 per cent of the electrons would have an energy beyond the ionization threshold because of the choice of the initial conditions.

Initial position and momentum have been chosen as if the electron were not under the influence of the field. In Fig. 1 we report the evolution of spatial distribution; it is clear that already after a short period, \( \leq 5 \) au, the spatial distribution has reached the equilibrium value compatible with the field.

Target and projectile nuclei have been set at an initial distance of \( 20Z \) au, where \( Z \) is the projectile charge. The equations of motion have been integrated until the nuclei were well far apart. For each test a number of runs ranging between \( 3 \cdot 10^4 \) and \( 26 \cdot 10^4 \) was performed. We did not attempt to minimize the statistical error for every point, however for all of the data the error goes from 1 to a maximum of 20 per cent.

When considering the final state of the electron it is essential to remember that an electron captured in an high enough quantum number may be subsequently field ionized. We have followed the convention of Olson and MacKellar (1981) of considering as ionized such electrons.

Three different field geometries have been considered: adopting the system of reference where the neutral atom is initially at rest and the ion is moving along the \( z \) axis towards the positive direction, the electric field has been chosen either parallel to \( z \), antiparallel or perpendicular. In the following we will adopt the convention of designate a parallel or antiparallel field according to its projection on the \( z \) axis, so \( F > 0 \) means \( \vec{F} \parallel \vec{v}_p \).

III. Hydrogen–proton collisions

The first series of calculations was about \( H - H^+ \) collisions: in Fig. 2 we plotted electron capture cross section, \( \sigma_{\text{cx}} \), and ionization, \( \sigma_{\text{ion}} \), versus impact velocity for two different values of the field, chosen parallel. It may be stated, alike shown in Olson and MacKellar (1981), that \( F \) enhances ionization, but at the same time opposes to the process of recombination on the projectile, working against charge exchange. The overall effect is quite clearly increasing with the strength of \( F \) as far as \( \sigma_{\text{ion}} \) is concerned, while for \( \sigma_{\text{cx}} \) the effect is less marked, even though visible. We note further that the field is effective on electron capture only at low energies, its influence being negligible above
v = 2 au (i.e. \( E = 100 \text{ KeV/amu} \)). Instead, \( \sigma\text{ion} \) is affected in the same manner along all the energy range.

In Fig. 3 the same calculations have been performed by varying the field direction. We may state that \( \sigma_{\text{cx}} \) is dependent from the sign of \( F \), the difference being small but regular. In Olson and MacKellar (1981) it was already suggested that when the field pulled the electron in the direction of the outgoing projectile, electron capture and loss were augmented. A less definite effect is visible on \( \sigma\text{ion} \). There are not appreciable differences when the field is set orthogonal to the ion direction or parallel to it.

The calculations were then repeated with a change in the initial condition, the electron being now set in the \( n = 2 \) state. From Fig. 4, where we have plotted the data corresponding to two opposite values of the field, one argues that the difference between \( \sigma_{\text{cx}}(F < 0) \) and \( \sigma_{\text{cx}}(F > 0) \) is clearly enhanced, while not a clear behavior along the entire energy range appears for \( \sigma\text{ion} \).

It is well known that, in field–free conditions, the collision process has certain symmetries, basing upon which some scaling laws can be inferred (see, for example, Janev 1991): all collisions between an hydrogen atom in the state \( n \) and an ion of charge \( Z \) lie on the same curve when one plots cross sections in terms of the reduced quantities

\[
\tilde{\sigma}(\tilde{E}) = \frac{\sigma(E)}{n^4Z}, \quad \tilde{E} = \frac{n^2E}{\sqrt{Z}}.
\]

Under non zero–field conditions the same relations should hold, provided the external field is scaled according

\[
\tilde{F} = \frac{F}{n^4}.
\]

In Fig. 5 we have plotted the same quantities of Figs. 3 and 4 in terms of these new reduced variables; both for \( \sigma_{\text{cx}} \) and \( \sigma\text{ion} \) a single trend is discernible, but the spreading of the points is not negligible.

As far as differential cross sections are concerned, we have plotted in Fig. 6 the quantities \( d\sigma/db \) for two scatterings at the same impact velocity and opposite values of \( F \). The effect of \( F \), in both cases, is to reduce the effective range within which capture takes place, while ionization is now made possible at larger impact parameters.

**IV. Hydrogen–multiply charged ions scattering**
In order to study features of non–resonant scattering we made simulations of collisions of hydrogen with $He^{2+}$, $Li^{3+}$, $C^{6+}$, and $O^{8+}$ ions, whose results are reported in Fig. 7. We may observe that a different behavior exist between the lightest ions ($H$, $He$), and the others. This is also clearly revealed in Fig. 8 where the same quantities have been plotted in terms of the reduced variables of eq. 3; however we observe that $Z$–scaling of eq. 3 still holds for the heaviest ions and on the whole this scaling law is obeyed quite well.

A further subject is the distribution of $\sigma_{cx}$ over quantum states of the projectile. In this work we only considered distributions over the principal quantum number $n$. Classically, a bound electron may have any value of the energy, below the ionization potential. Capture into discrete energy levels is approximated by a procedure developed, for example, in Olson (1981) and Salop (1979). It is well known that there exist a proportionality relation between the ion charge $Z$ and the quantum state $n$ into which the electron is preferentially captured, $n \approx Z^{3/4}$ (Olson 1981). In Fig. 9 we plotted these distributions for $C^{6+}$ and $O^{8+}$ ions: $\vec{F}$ has no effect when capture takes place in low lying states while, as one could foresee, capture in higher states, which are less bound to the nucleus, is strongly suppressed. The position of the maximum is unaltered with respect to the null-field case (Olson 1981, Salop 1979). We remark that $n = 7$ and $n = 9$ are, respectively for carbon and oxygen, the maximum non field–ionized quantum numbers at this value of $F$.

V. Shrinking the extension of the field

Until now all the calculations have been performed under the hypothesis of an uniform external field, extending in all directions up to infinity. As a final issue we modified this configuration by limiting the field within a finite interval: it is abruptly set to zero outside a distance $L$ from the target nucleus along the $z$–axis in both directions. After this modification we have a more symmetrical situation, in which both Coulomb interparticle forces and $\vec{F}$ can exercise an influence over only a limited region. Ideally, the length $L$ should be taken great enough to include most of the interparticle interaction within it: we chose $L = 10 \, Z \, \text{au}$. Repeating previous calculations for $H^-H^+$ collisions gives some modifications: $\sigma_{cx}$ is no longer uniformly suppressed and may become even greater than in the zero field case, and the opposite is true for $\sigma_{\text{ion}}$ (Fig. 10). Another non negligible difference stems when dealing with
high charged ions: in Fig. 11 we have plotted n−distributions for $H - C^{6+}$ impacts at a velocity $v_p = 2$. The absolute values of $\sigma_{cx}$ corresponding to $F = 0, +0.03, -0.03$ au are, respectively, $9.5 \cdot 10^{-16} \text{ cm}^2$, $5.4 \cdot 10^{-16} \text{ cm}^2$, $17.4 \cdot 10^{-16} \text{ cm}^2$, and the auxiliary contribution to the $F < 0$ case comes from electrons which are bound in high lying states around $n = 10, 11$. We suggest that this result may be explained by a mechanism similar to capture–via–ionization (see McCartney and Crothers 1994): the incoming projectile strikes the electron, raising it in the continuum with an initial velocity $\vec{v}_0$, which we will assume to be parallel to the ion velocity $\vec{v}_p$. The electron is then accelerated by the field. On $z = L$, it reaches the maximum speed which, with the position $F = -|F|$, may be written $v_f \approx \sqrt{v_0^2 + 2L|F|}$. The relative distance between the ion and the electron is now

$$D = v_p \frac{v_f - v_0}{|F|} - L. \quad (5)$$

Capture takes place with preference when $v_p$ and $v_f$ closely match, $v_p \approx v_f$. The binding energy is then

$$E \approx -Z \frac{D}{2n^2}, \quad (6)$$

from which one gets $n \approx [\sqrt{ZD/2}]$ ([ ] means the integer part). The condition $v_p = v_f$ gives $v_0 = \sqrt{v_p^2 - 2L|F|}$. The substitutions $v_p = 2$, $Z = 6$, $L = 60$, $|F| = 0.03$ give $n \approx 10$ in good accordance with Fig. 11. A similar, but much weaker, feature appears in the $F > 0$ case; the maximum at $n = 7$ may be partially explained by the same mechanism, provided that now the field has a braking effect on the electron, whose initial speed must thus be greater than $v_p$.

VI. Summary

In this work we have performed some numerical simulations with the aim to investigate how a strong static electric field may modify an energetic collision between an ion and an atom. We studied total and partial cross sections for charge–exchange and ionization versus some meaningful parameters as impact velocity $v_p$, ion charge $Z$, initial target quantum state $n$ and spatial extension of the field. It has been found that the field deeply affects
the results of the scattering; detailed effects depend upon the strength of the field and its orientation with respect the relative nuclear motion. Scaling laws, true in null-field situations, seem now more doubtful, and are surely true only within strict ranges of the parameters.

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References

Abrines R A and Percival I C 1966a *Proc. Phys. Soc.* **88** 862
— 1966b *Proc. Phys. Soc.* **88** 873
Bezchastnov V G and Potekhin A J 1994 *J. Phys. B: At. Mol. Opt. Phys.* **27** 3349
Buchleitner A, Gremaud B and Delande D 1994 *J. Phys. B: At. Mol. Opt. Phys.* **27** 2663
Bylicki M, Themelis S I and Nicolaides C A 1994 *J. Phys. B: At. Mol. Opt. Phys.* **27** 2741
Cohen J S 1985 *J. Phys. B: At. Mol. Phys.* **18** 1759
Damburg R J and Kolosov V V 1976 *J. Phys. B: At. Mol. Phys.* **9** 3149
Delande D and Gay J C 1986 *Phys. Rev. Lett.* **57** 2006
Delande D et al 1994 *J. Phys. B: At. Mol. Opt. Phys.* **27** 2771
Eichenauer D, Grün N and Scheid W 1981 *J. Phys. B: At. Mol. Phys.* **14** 3929
Eliezer S et al 1988 Laser Interaction and Related Plasma Phenomena (eds. Hora H and Miley P G, Plenum, New York, vol. 8) p. 279
Eliezer S and Hora H 1989 *Phys. Rep.* **172** 339
Friedrich H 1990 Atoms in Strong Fields (eds. C.A. Nicolaides, C. Clark and M. Nayfeh) p. 247
Grosdanov T P and McDowell M R C 1985 *J. Phys. B: At. Mol. Phys.* **18** 921
Hardie D J W and Olson R E 1983 *J. Phys. B: At. Mol. Phys.* **16** 1983
Janev R K 1991 *Phys. Lett. A* **160** 67
Kunc J A 1988 *J. Phys. B: At. Mol. Phys.* **21** 3619
McCartyney M and Crothers D S F 1994 *J. Phys. B: At. Mol. Opt. Phys.* **27** L485
McDowell M R C and Zardona M 1985 *Adv. At. Mol. Phys.* **21** 255
Olson R E 1981 *Phys. Rev. A* **24** 1726
Olson R E and MacKellar A D 1981 *Phys. Rev. Lett.* **46** 1451
Raadu M A 1989 *Phys. Rep.* **178** 25
Reinhold C O and Falcon C A 1988a *J. Phys. B: At. Mol. Phys.* **21** 1829
— 1988b *J. Phys. B: At. Mol. Phys.* **21** 2473
Ruder H, Wunner G, Herold H and Geyer F 1994 Atoms in Strong Magnetic Fields (Springer–Verlag)
Salop A 1979 *J. Phys. B: At. Mol. Phys.* **12** 919
Schmidt A, Horbatsch M and Dreizler R M 1990 *J. Phys. B: At. Mol. Phys.* **23** 2327
Seipp I and Taylor K T 1994 *J. Phys. B: At. Mol. Opt. Phys.* **27** 2785
Wang Q and Greene C H 1991 *Phys. Rev. A* **44** 7448
Williams A C *et al* 1986 *Astrophys. J.* **305** 759
Zang J X and Rustgi M L 1994 *Phys. Rev. A* **50** 861
Figure Captions

**Figure 1**: upper, radial distribution; lower, polar angle distribution. Full line represents microcanonical distribution at \( F = 0 \). We set up the field \( F = 0.03 \) au at \( t = 0.0 \). Our plots show distributions at \( t = 0.0 \) (full circles), \( t = 5.0 \) au (empty circles), \( t = 10.0 \) au (full squares). A total of \( 10^5 \) runs have been performed.

**Figure 2**: upper, \( \sigma_{cx} \); lower \( \sigma_{ion} \) versus impact velocity. Full diamonds, \( F = 0.0 \); open circles, \( F = 0.03 \) au; full circles, \( F = 0.04 \) au.

**Figure 3**: upper, \( \sigma_{cx} \); lower \( \sigma_{ion} \) versus impact velocity. Full diamonds, \( F = 0.0 \); open circles, \( F = 0.03 \) au; full circles, \( F = -0.03 \) au; open squares, \( F = 0.03 \) au (orthogonal).

**Figure 4**: upper, \( \sigma_{cx} \); lower \( \sigma_{ion} \) versus impact velocity. Open circles, \( F = 0.03/2^4 \) au; full circles, \( F = -0.03/2^4 \) au. The electron is in the initial status \( n = 2 \).

**Figure 5**: upper, charge exchange; lower, ionization. In ordinate appear scaled cross sections \( \sigma/n^4 \); in abscissae scaled velocity \( v/n \). Open circles, \( F = 0.03 \) au, \( n = 1 \); full circles, \( F = 0.03/2^4 \) au, \( n = 2 \); open squares, \( F = -0.03 \) au, \( n = 1 \); full squares, \( F = -0.03/2^4 \) au, \( n = 2 \).

**Figure 6**: upper, \( d\sigma_{cx}/db \); lower, \( d\sigma_{ion}/db \) versus impact parameter \( b \). Impact energy is kept fixed at \( E = 50 \) KeV/amu. Full diamonds, \( F = 0.0 \); open circles, \( F = 0.03 \) au; full circles, \( F = -0.03 \) au.

**Figure 7**: upper, \( \sigma_{cx} \); lower \( \sigma_{ion} \) versus impact velocity. Field strength is kept fixed at \( 0.03 \) au. Full triangles \( H^+ \); full circles \( He^{2+} \); open circles \( Li^{3+} \); full squares \( C^{6+} \).

**Figure 8**: upper, charge exchange; lower, ionization. In ordinate appear scaled cross sections \( \sigma/Z \); in abscissae scaled velocity \( v/Z^{1/4} \). Open triangles, \( O^{8+} \); the other symbols are the same as in Figure 7.

**Figure 9**: distribution of \( \sigma_{cx} \) over \( n \). Upper, \( C^{6+} \); lower, \( O^{8+} \). Open squares \( E = 50 \) KeV/amu; full squares \( E = 100 \) keV/amu.

**Figure 10**: upper, \( \sigma_{cx} \); lower \( \sigma_{ion} \) versus impact velocity. Full diamonds, \( F = 0.0 \); open circles, \( F = 0.03 \) au; full circles, \( F = -0.03 \) au; open squares, \( F = 0.03 \) au (orthogonal). The length \( L \) has been kept fixed at \( L = 10 \) au.

**Figure 11**: distribution of \( \sigma_{cx} \) over \( n \) for \( C^{6+} - H \) impacts at an energy \( E = 100 \) KeV/amu. Diamonds, \( F = 0.0 \); open circles, \( F = 0.03 \) au; full circles, \( F = -0.03 \) au. The length \( L \) is kept fixed, \( L = 60 \) au.