A&A 375, 319–327 (2001)
DOI: 10.1051/0004-6361:20010738

Astronomy & Astrophysics

Successsive resonance-enhanced two-photon ionization of elements abundant in nebulae

I. Atoms and ions of C, N, and O

S. Johansson[1] and V. S. Letokhov[2,1]

[1] Department of Physics, Lund University, PO Box 118, 22100 Lund, Sweden
e-mail: Vladilen.Letokhov@fysik.lu.se
[2] Institute of Spectroscopy, Russian Academy of Sciences, Troitsk, Moscow region, 142190, Russia

Received 8 January 2001 / Accepted 21 May 2001

Abstract. We discuss resonance-enhanced two-photon ionization (RETPI) and present schemes of successive RETPI of the elements C, N, and O in nebulae. RETPI is activated by intrinsic radiation stored in the form of trapped spectral lines of HI, HeI, and HeII in the optically thick nebula. The rate of this two-step photoionization is comparable with or exceeds the low recombination rate of the photoions formed in the process. This leads to an accumulation of photoions and subsequent RETPI until such highly charged ions are formed that they cannot further be ionized in this way by the intrinsic radiation from the strong spectral lines of HI, HeI, and HeII.

Key words. atomic processes – radiation mechanisms – planetary nebulae: general

1. Introduction

Photoinduced processes have been known to play an important part in the physics of planetary and gaseous nebulae (Aller 1956, 1984; Pottasch 1984; Osterbrock 1989), as they are responsible for the inflow of radiative energy from the central star (stars). The short-wavelength radiation provides the photoionization absorption of hydrogen and helium, whose atoms repeatedly take part in this process. The H and He ions recombine with electrons, thus providing for the permanent conversion of the absorbed stellar radiation into kinetic energy of the photoelectrons produced and radiative energy of the recombination transitions, especially the high-intensity EUV lines in the Lyman series of HI and HeII and the resonance lines of HeI. This line radiation suffers from resonance scattering due to diffusion trapping in the optically dense nebula, which results in Doppler diffusive broadening of the spectral lines and an increase of their intensity inside the nebula. Both these effects are important for the radiative cooling of the nebula because of the escape of trapped photons through the wings of the spectral lines (the Zanstra effect) (Zanstra 1949) and the escape of optically thin fluorescence lines of other elements produced in a Bowen mechanism (Bowen 1935). The latter is caused by an accidental wavelength coincidence between a H or He line and an absorption line (1→2) of these elements (see Fig. 1a).

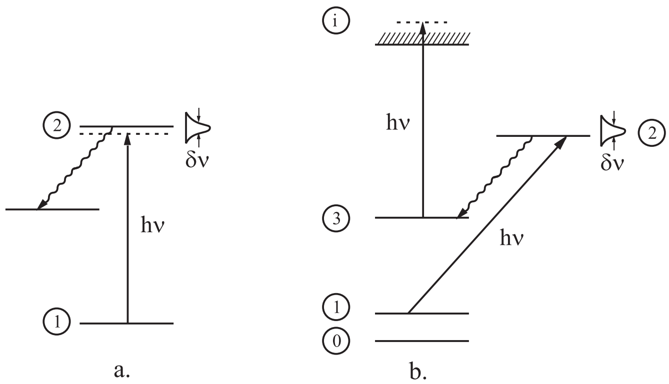

Fig. 1. a) Accidental resonance "bound-bound state" excitation $1 \rightarrow 2$ of Bowen fluorescence and **b)** possible subsequent "bound-free" photoionization transition $3 \rightarrow i$ via the metastable state 3.

The efficiency of the resonant excitation by the Bowen mechanism is especially illustrated by two UV Fe II lines around 2508 Å, observed with the Hubble Space Telescope in high spectral and spatial resolution of gaseous condensations in the vicinity of η Car (Johansson et al. 1996). The abnormally high brightness of these lines is associated with the Bowen excitation of FeII from the low-lying metastable state $a^4D_{7/2}$ (state 1 in Fig. 1b) by HLyα (Johansson & Hamann 1993). To explain the high brightness of these lines and the abnormal intensity ratios between them and their satellite lines, we have considered in (Johansson & Letokhov 2001a) the possibility of a subsequent photoionization of Fe II from the long-lived states

Send offprint requests to: S. Johansson,
e-mail: sveneric.johansson@fysik.lu.se

$c^4F_{7/2,9/2}$ (state 3 in Fig. 1b), in which FeII gets accumulated as a result of the fast UV 2→3 decay. It is important to note that the rate of the photoionization process involving diluted radiation from the central star is much slower than the photoionization rate associated with the intense radiation from the trapped HLyα line, whose wavelength at 1215 Å is shorter than the threshold wavelength required to photoionize FeII from the state $c^4F_{7/2}$ (λ = 1244 Å). This mechanism of biphotonic stepwise photoionization (Johansson & Letokhov 2001b) of FeII by HLyα explains the appearance of the FeIII lines observed in the same region of η Car as the bright FeII lines (Johansson et al. 2000)

The slow recombination rate of the photoions produced in rarefied nebular plasmas allows their accumulation at a comparatively low photoionization rate. We have, therefore, considered the possibility of resonance-enhanced two-photon ionization (RETPI) even in the absence of the exact intermediate resonance (Johansson & Letokhov 2001c), necessary for a photoselective excitation mechanism. In the RETPI case, the intrinsic EUV radiation of the very intense lines of HI, HeI, and HeII, trapped in the nebula, can stimulate RETPI to occur in several ionization stages of various elements in spite of its low rate compared to laboratory laser experiments (Letokhov 1987). In the present paper, we consider RETPI schemes for C, N, and O, which are of interest in the understanding of the mechanisms that convert the hydrogen and helium radiation trapped in the nebula into spectral lines of other elements. The implementation of this new elementary mechanism for particular problems (shift in the ionization balance, anomalous intensity ratios of spectral lines, etc.) in specific planetary nebulae is the subject for future considerations and are out of the scope of the present paper.

2. The rate of RETPI in a dichromatic radiation field

Let an isolated atomic particle X (a neutral atom or an ion in a given ionization state N) reside in an isotropic radiation field with a spectral intensity of $P_1(\nu)$ [in photons/cm^2 s Hz]. Let also the spectral width of the radiation field $\delta\nu_1$ exceed the Doppler width $\Delta\nu_{\rm D}$ of the allowed transition $1 \rightarrow 2$, and the shift $\Delta\nu$ (hereafter called "detuning") of the central frequency ν_1 of the radiation field from the allowed transition ν_{12} exceed the spectral width $\delta\nu_1$ (Figs. 2a,b). In that case, the one-quantum resonant excitation to level 2 is impossible because the energy defect $h\Delta\nu$ cannot be transferred to a third partner in a rarefied nebular medium. But, resonance scattering of the radiation is possible here, and its probability is determined by the Lorentzian wing of the natural (radiative) broadening $\gamma = 2A_{21}$ of the transition $1 \rightarrow 2$, whose amplitude is proportional to $(\frac{\gamma/2}{2\pi\Delta\nu})^2$ (Weisskopf 1933). The direct excitation of level 2 is only possible as a result of the two-quantum excitation of a virtual level with an energy of $2h\nu_1$ and the subsequent spontaneous emission of a photon with a frequency of $h(2\nu_1{-}\nu_{12})$. According to Makarov (1983), the probability of this process is proportional to $(\frac{\gamma/2}{2\pi\Delta\nu})^4$, and it is substantially lower than the resonance scattering probability.

However, the atomic particle XN can virtually be in the excited state 2 with a probability of W_2. It is not difficult to calculate W_2 and express it in terms of the Einstein coefficient A_{21}.

For the sake of simplicity we assume that the radiation field with the amplitude E is linearly polarized and acts in a coherent fashion at the frequency ν_1 detuned by $\Delta\nu$ from the frequency ν_{12}. Actually, the quantity W_2 is independent of the polarization of the radiation. In the case when the frequency detuning $\Delta\nu \gg \delta\nu_1$, the oscillations of the wave function amplitudes are faster than the oscillations of the field amplitude. Under such a condition the interaction of the radiation field with a two-level quantum system can be treated in a coherent way. According to perturbation theory we have

$$W_2 = \frac{1}{g_1}\sum_{M_1}\left(\frac{d^{(M_1M_2)}E}{2\hbar\cdot 2\pi\Delta\nu}\right)^2 \qquad (1)$$

where the summation is extended over all magnetic sublevels of the state 1, and $\Delta\nu$ is expressed in Hz. The quantity $d^{(M_1M_2)}$ is the dipole moment of the transition between sublevels M_1 and M_2.

Next we want to express the sum of the squared dipole moments on the right-hand side of formula (1) in terms of A_{21}. We can use the well-known formula relating the matrix element of the dipole moment operator for a pair of non-degenerate sublevels to the rate $\gamma_{21}^{(M_2,M_1)}$ of spontaneous transitions between these sublevels (see, for example, Sobel'man 1979):

$$\gamma_{21}^{(M_2,M_1)} = \frac{4\omega_{21}^3}{3\hbar c^3}\left(d_{21}^{(M_2M_1)}\right)^2, \qquad (2)$$

where $\omega_{21} = 2\pi\nu_{21}$. The total decay rate of any magnetic sublevel M_2 is

$$A_{21} = \sum_{M_1}\gamma_{21}^{(M_2,M_1)}. \qquad (3)$$

We use another well-known formula for the matrix elements of vectors (see, for example, Landau & Lifshitz 1958) and obtain

$$\sum_{M_1}\left(d_{21}^{(M_2,M_1)}\right)^2 = \frac{\hbar c^3}{4\omega_{21}^3}g_2A_{21}. \qquad (4)$$

We express E^2 on the right-hand side of formula (1) in terms of the radiation intensity I_1 (in photons/cm^2 s) at the frequency $\omega_1 = 2\pi\nu_1$:

$$E^2 = \frac{8\pi}{c}\hbar\nu_1 I_1. \qquad (5)$$

Finally, by substituting formulae (4) and (5) into Eq. (1), we get

$$W_2 = \frac{1}{32\pi^3}\frac{g_2}{g_1}\frac{\lambda_{21}^3}{\lambda_1}\frac{A_{21}}{(\Delta\nu)^2}I_1 \approx \frac{1}{32\pi^3}\frac{g_2}{g_1}\frac{\lambda_{21}^2A_{21}}{(\Delta\nu)^2}I_1, \qquad (6)$$

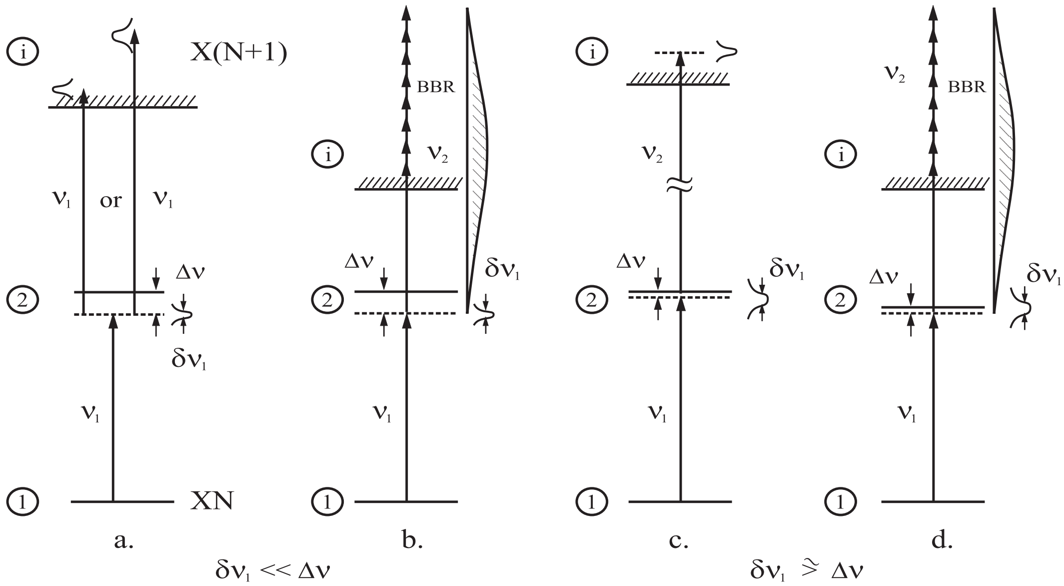

Fig. 2. Various schemes of resonance-enhanced two-photon ionization (RETPI) by two monochromatic spectral lines with various detuning $\Delta\nu$ of the exciting spectral line ν_1: **a)** and **b)** $\delta\nu_1 \ll \Delta\nu$; **c)** and **d)** $\delta\nu_1 \geq \Delta\nu$; and with a small binding energy of the excited state 2 in **b)** and **d)**.

where g_i is the statistical weight of the ith level, and the lower level 1 is the ground state and $\lambda_1 \approx \lambda_{21}$. Note that there is no need here to account for the Doppler broadening $\Delta\nu_{\mathrm{D}}$ of the transition 1→2. The shift $\Delta\nu$ in the transition frequency of all atomic particles is much greater than $\Delta\nu_{\mathrm{D}}$, and they are practically equally excited to the virtual level, no matter what velocity they have.

The ionization energy of the atomic particle in the virtually excited state is (IP − $h\nu_1$), where IP is the ionization potential of XN. It is lower than either of the photon energy $h\nu_1$ of the same radiation field and the photon energy of a second intense field with an average frequency ν_2, for which $h\nu_2 >$ (IP − $h\nu_1$). In that case, there exists a certain probability that the atomic particle will make a two-quantum transition through the virtual level to the ionization continuum at a rate $W_{\mathrm{ph}}^{(2)}$ (s^{-1}) given by:

$$W_{\mathrm{ph}}^{(2)} = W_2 W_{\mathrm{ph}}^{2i} \tag{7}$$

where W_{ph}^{2i} is the photoionization rate of the excited level 2 expressed in terms of the overlapping integral of the spectral radiation intensity $P_2(\nu)$ and the spectral dependence of the photoionization cross section of the excited level, $\sigma_{2i}(\nu)$:

$$W_{\mathrm{ph}}^{2i} = \int_{\mathrm{IP}-h\nu}^{\infty} \sigma_{2i}(\nu) P_2(\nu)\mathrm{d}\nu. \tag{8}$$

The radiation in the first-step excitation can also serve as a radiation in the second step, provided that (IP−$h\nu_1$) $< h\nu_1$. In the case of relatively monochromatic intense trapped lines inside the nebula, the spectrum width $\delta\nu_2$ is much smaller than the spectral resonances for photoionization from the excited state (Fig. 2a), so that the photoionization rate from this state can be written as

$$W_{\mathrm{ph}}^{2i} \approx \sigma_{21}(\nu_2) I_2 \tag{9}$$

where I_2 is the intensity of the narrow-band radiation at the frequency ν_2 (in photons/cm^2 s). Based on these considerations, one obtains from formulae (6)–(9) a simple expression for the rate of the resonance-enhanced two-photon ionization of an atomic particle in a two-frequency isotropic radiation field (subject to the condition ($|\Delta\nu| \gg \delta\nu_1$):

$$W_{\mathrm{ph}}^{(2)} = \frac{\lambda_{21}^2}{32\pi^3} \frac{g_2}{g_1} \frac{A_{21}}{(\Delta\nu)^2} \sigma_{2i} I_1 I_2 \tag{10}$$

where $\Delta\nu$ is expressed in Hz.

In the case of intense non-monochromatic photoionizing radiation (for example, diluted, but spectrally wide, black body radiation from the central star as shown in Fig. 2b) it is necessary to use expression (8) for an estimation of W_{ph}^{2i}.

Before we discuss the less frequent case of exact resonance ($|\Delta\nu| \leq \delta\nu_1$), where the Bowen resonance excitation mechanism proves effective, let us estimate the rate of the RETPI process under nebular conditions. A specific feature of planetary nebulae is the huge optical thickness in the spectral lines of transitions to the ground state or low metastable states. Because of this, the photons resulting from the recombination of the photoions HII (or HeII), which are being constantly produced under the effect of the ionizing radiation of the star, repeatedly undergo resonance scattering before they leave the nebular medium and thus become observable. This diffusive confinement (or trapping) is accompanied by a broadening of the nebular emission lines and a large increase of the radiation intensity inside the nebula. The broadening is limited either by the decay of the photons as a result of absorption or by their escape from the spectral line wings due to a Doppler frequency redistribution induced by their scattering on moving resonant particles. All these effects are essential for spectral lines having a huge optical thickness $\tau_0 \approx 10^4$–10^8. As a result, the intensity P of the trapped spectral lines reaches its maximum, i.e., the intensity of the black-body radiation at the corresponding wavelength λ, and can be described in terms of the effective temperature $T_{\rm eff}$:

$$P(\lambda, T) = \frac{8\pi}{\lambda^2} \frac{\delta\nu}{{\rm e}^{h\nu/kT_{\rm eff}} - 1} [\text{photons/cm}^2\,\text{s}], \tag{11}$$

where $\delta\nu$ is the spectral line width [in Hz]. According to numerous measurements (Aller 1956, 1984; Osterbrock 1989) the effective temperature $T_{\rm eff}$ of the electrons and the trapped radiation for most nebulae reaches values as high as (10–15) $\times$ 10^3 K, and the width $\delta\nu_1$ of the trapped spectral lines at an optical density of $\tau_0 \approx 10^8$ can reach a few hundreds of cm^{-1}. Under such conditions, the intensity of the strong trapped spectral lines is very high, a fact noted in many works (Aller 1956, 1984; Osterbrock 1989). For example, for HLyα (λ = 1215 Å) at $T_{\rm eff}$ = 15 $\times$ 10^3 K and $\delta\nu \approx 250$ cm^{-1}, the intensity $P \approx 6 \times 10^{20}$ photons/cm^2 s, which corresponds to 10^3 W/cm^2 of CW(continuous wave) EUV radiation, a figure as yet unattainable in spite of the status of present-day advanced laser technology. Such high radiation intensity at a comparatively small occupation number of the photon quantum state, $< n > \approx 5 \times 10^{-4}$, is explained by the large number of free space states within the limits of a solid angle of 4π steradians and the large spectral width $\delta\nu$.

In terms of the effective temperature $T_{\rm eff}$ of the equivalent Planck radiation (11), the RETPI rate (10) may be represented in the form

$$W_{1i} = \frac{2}{\pi} \frac{g_2}{g_1} \frac{\delta\nu_1 \delta\nu_2}{(\Delta\nu)^2} \frac{\sigma_{2i}}{\lambda_2^2} A_{21} \left[\exp\left(\frac{h\nu_1}{kT^1_{\rm eff}}\right) - 1 \right]^{-1} \times \left[\exp\left(\frac{h\nu_2}{kT^2_{\rm eff}}\right) - 1 \right]^{-1}. \tag{12}$$

Since $h\nu_1$ in the case of HLyα and other, short-wavelength intense lines (HI, HeI, HeII) is much greater than $kT^{1,2}_{\rm eff}$, Eq. (12) may be represented in the more simple form

$$W_{1i} \approx \frac{2}{\pi} \frac{g_2}{g_1} \frac{\delta\nu_1 \delta\nu_2}{(\Delta\nu)^2} \frac{\sigma_{2i}}{\lambda_2^2} A_{21} \exp\left(-\frac{h\nu_1}{kT^1_{\rm eff}} - \frac{h\nu_2}{kT^2_{\rm eff}} \right). \tag{13}$$

As an example let us consider a sequence of ions having $\sigma_{2i} \approx 10^{-17}$ cm^2 and $A_{21} \approx 10^9$ s^{-1}. For HLyα (λ_{12} = 1215 Å) with a spectral width $\delta\nu \approx 300$ cm^{-1} and a detuning of $\Delta\nu \approx 1000$ cm^{-1}, the rate $W^{(2)}_{\rm ph}$ reaches values of 10^{-7}–10^{-6} s^{-1} at a temperature $T_{\rm eff} \approx (13$–$15) \times 10^3$ K. This value of $W^{(2)}_{\rm ph}$ is very small for laboratory laser experiments (Letokhov 1987), but it is quite substantial under nebular conditions, for which the recombination rate $W_{\rm rec}$ is very low. At $W^{(2)}_{\rm ph} > W_{\rm rec}$ there will occur an accumulation of photoions in the next higher ionization state.

To understand the efficiency of the RETPI process, let us compare its rate W_{1i} with the rate of ionization by electrons. In both cases, there exists the exponential factor $\exp(-\frac{E}{kT})$, where $E = h(\nu_1 + \nu_2) >$ IP, IP is the ionization potential, and T is the corresponding radiation or electron temperature. Thus, the main difference between collisional ionization by electrons and collisionless RETPI by photons lies in the difference between the non-exponential factor denoted by $A_{\rm e}$ and $A_{\rm ph}$, respectively. According to (13), in the case of RETPI

$$A_{\rm ph} \approx \frac{2}{\pi} \frac{g_2}{g_1} \frac{\delta\nu_1 \delta\nu_2}{(\Delta\nu)^2} \frac{\sigma_{2i}}{\lambda_2^2} A_{21}, \tag{14}$$

and for the case considered above, $A_{\rm ph} \approx (1$–$10)$ s^{-1}. In the case of electronic excitation, the corresponding non-exponential factor is $A_{\rm e} \approx \langle n_{\rm e} \upsilon_{\rm e} \sigma_{\rm e} \rangle$, where $\upsilon_{\rm e}$ is the velocity of electrons with the energy $E >$ IP, $n_{\rm e}$ is the electron concentration, $\sigma_{\rm e}$ is the electronic ionization cross section. For $n_{\rm e} \approx 10^5$ cm^{-3}, $A_{\rm e} \approx 10^{-2}$–$10^{-3} \ll A_{\rm ph}$. Therefore, the rate of formation of ions by RETPI is higher than that through collisions with energetic electrons, and it is quite comparable with the recombination rate of the photoions produced. For this reason, the concentration of photoions (for example, $X(N+1)$) is much higher than that provided by collisions with the electrons having kinetic energy $E_{\rm kin} >$ IP, where IP is the ionization potential of XN. Accordingly, the recombination emission lines of XN become brighter. In addition, the concentration of the ions produced (for example, $X(N+1)$) becomes sufficient for a subsequent RETPI and as a result of that the appearance of bright emission lines from $X(N+1)$.

In the text below, we consider a situation where the most intense trapped spectral lines of HI, HeI, and HeII (lines in the range 232–1215 Å) in the EUV region are in close resonance with the allowed transitions in a number of atoms and their ions. These are thus capable of undergoing RETPI at a rate $W^{2i}_{\rm ph}$ comparable to the recombination rate $W_{\rm rec}$ in nebular media. We will refrain from considering any values of the actual effective temperature and spectral width, as these vary for different nebulae and within the same nebula as well. We restrict ourselves to analysis of specific cases, where we find close resonances of trapped H and He lines with the strongest transitions in carbon, nitrogen, and oxygen atoms and their ions in successive stages of ionization, i.e. we consider various chains of successive RETPI occurring in these elements.

As will be evident below, in some cases the frequency detuning $\Delta\nu$ is only a few tens of cm^{-1}, i.e., it can be smaller than the spectral line width $\delta\nu_1$ (Fig. 2c). In such an exact accidental resonance, as is the case in the Bowen mechanism, the excited level is actually populated. The maximum excitation probability W_2 would be limited by the effective temperature $T^1_{\rm eff}$ of the exciting radiation, provided that the radiation is in equilibrium with the resonant atoms:

$$W_2 = \left[\exp\left(\frac{h\nu_1}{kT^1_{\rm eff}}\right) - 1 \right]^{-1}. \tag{15}$$

In this case of exact resonance the biphotonic ionization rate increases and reaches, according to (7) and (15), its maximum

$$W_{\rm ph}^{(2)}(1-i) = \sigma_{2i}(\nu_2) I_2 \left[\exp\left(\frac{h\nu_1}{kT_{\rm eff}^1}\right) - 1\right]^{-1} \qquad (16)$$

or

$$W_{\rm ph}^{(2)}(1-i) \;=\; 8\pi \frac{\sigma_{2i}(\nu_2)}{\lambda_{12}^2} \delta\nu_2 \left[\exp\left(\frac{h\nu_1}{kT_{\rm eff}^1}\right) - 1\right]^{-1} \times \left[\exp\left(\frac{h\nu_2}{kT_{\rm eff}^2}\right) - 1\right]^{-1}. \qquad (17)$$

As will be demonstrated below, in some cases the photon energy in the photoionization process might be very small (in the visible region) as shown in Figs. 2b,d, and the diluted black body radiation of the central star can contribute to the rate of step-wise photoionization. The rate of the RETPI process, which involves black body radiation in the second step, depends upon the condition of the nebula, and in particular upon the distance from the central star.

The detailed calculation of the rates of resonance-enhanced two-photon and of resonance stepwise two-photon ionization by dichromatic isotropic noncoherent radiation will be published elsewhere.

3. Successive RETPI schemes for C, N, and O

Analysis of the allowed quantum transitions in the carbon atom and subsequent carbon ions, on the basis of the data published in Bashkin & Stoner (1975), immediately points to the possibility of RETPI with suitable pairs of intense trapped lines of HI, HeI, and HeII in the EUV region extending from HLyα (1215 Å) to HeIILyγ (243 Å). Out of a large number of suitable pairs of lines, we included only those which have close coincidences with allowed transitions in the carbon atom or the carbon ions. We have not considered transitions characterized by either a large frequency detuning $\Delta\nu$ or a low transition probability A_{21} as the RETPI rate could be low according to (10) and (13).

The RETPI schemes for carbon (Figs. 3, 4) provide a path to CV. However, there is no suitable quasi-resonance for C V with any of the EUV spectral lines of H and He, which means that this ion is the end ion in the successive RETPI chain CI→CII→CIII →CIV→CV. For each step in the chain, the relative frequency detuning ($\Delta\nu/\nu_{21}$) does not exceed 1.5×10^{-2}. At an effective radiative temperature of the spectral lines in the range (15–20) $\times 10^3$ K the total RETPI chain achieves a relatively high rate of 10^{-6}–10^{-4} s^{-1}. In some steps, CI→CII, for example, the photoionization of the virtually excited state 2 can be carried out by radiation over a very broad spectral region: $\lambda_2 < 1.2$ μm. This means that the rate $W_{\rm ph}^{2i}$ defined by Eq. (8) can be high enough even for photoionization caused by the diluted black body radiation from the central star. However, this rate depends on the radiation dilution factor, i.e., the distance from the central star.

Analogous to the carbon scheme, RETPI of nitrogen can lead to N V. Moreover, for the RETPI cases NI→NII and NII→NIII, the detuning from the exact resonance frequency is very small, 300–400 cm^{-1} and $\pm$100 cm^{-1}, respectively. Therefore, one can expect that the relative concentration of the NI and NII ions will be low because of the significant shift of the photoionization balance toward multiply ionized nitrogen ions. As for carbon, the virtually excited states in the low ionization stages of nitrogen are close enough to the ionization limit to be entirely photoionized by the black body radiation from the central star.

Similar RETPI schemes for oxygen (Figs. 7, 8) can terminate with the OV photoion. Specific features of oxygen are the existence of several very close resonances between OI and the lines HLy$_{\beta,\gamma,\delta}$, and the close proximity of virtually excited oxygen states to the ionization limit. Obviously the ionization balance OI→OII should be substantially shifted toward OII.

All the excitation and ionization steps for the RETPI process in carbon, nitrogen and oxygen are summarized in Table 1, which shows the chains for the atoms and successive ions of these elements. The atoms and ions are exposed to pairs of spectral lines with progressively increasing photon energies: HI + HI, HI + HeI, HI + HeII, HeI + HeI, and finally HeII + HeII.

4. Conclusion

It seems reasonable that the proposed elementary process of photoionization by the intrinsic EUV radiation trapped in nebulae should be taken into consideration in calculating the ionization balance of those elements, whose observed spectral lines are essentially the only information available on the processes occurring inside nebulae. Since the RETPI rate $W_{\rm ph}^{(2)}$ apparently can exceed the collisional electron ionization rate $W_{\rm e}$, RETPI will enhance the corresponding recombination lines, which are frequently interpreted in the framework of ionization by electron collisions as being the result of an anomalous abundance effect (Aller 1984; Pottasch 1984).

From this standpoint, one can understand that the existing determinations of the chemical abundance from spectral data frequently disagree. For example, an anomalous ratio often observed for the abundances of CIV and CII can be explained by the RETPI of CIII, carried out jointly by the HLyγ line (972 Å) and HeIILyα at 303 Å or directly by HeIILyβ at 256 Å (Fig. 4). As another example, one can refer to the "puzzle" of the anomalous oxygen recombination lines in planetary nebulae (Dinerstein et al. 2000). The authors of that paper suggest the existence of a "line enhancement effect due to an unidentified physical mechanism". We believe the anomay can be the result of the RETPI process according to the scheme in Fig. 8.

Moreover, in the case of exact resonance in some link of the successive RETPI chain, the ionization balance between two successive ions of an element can be primarily

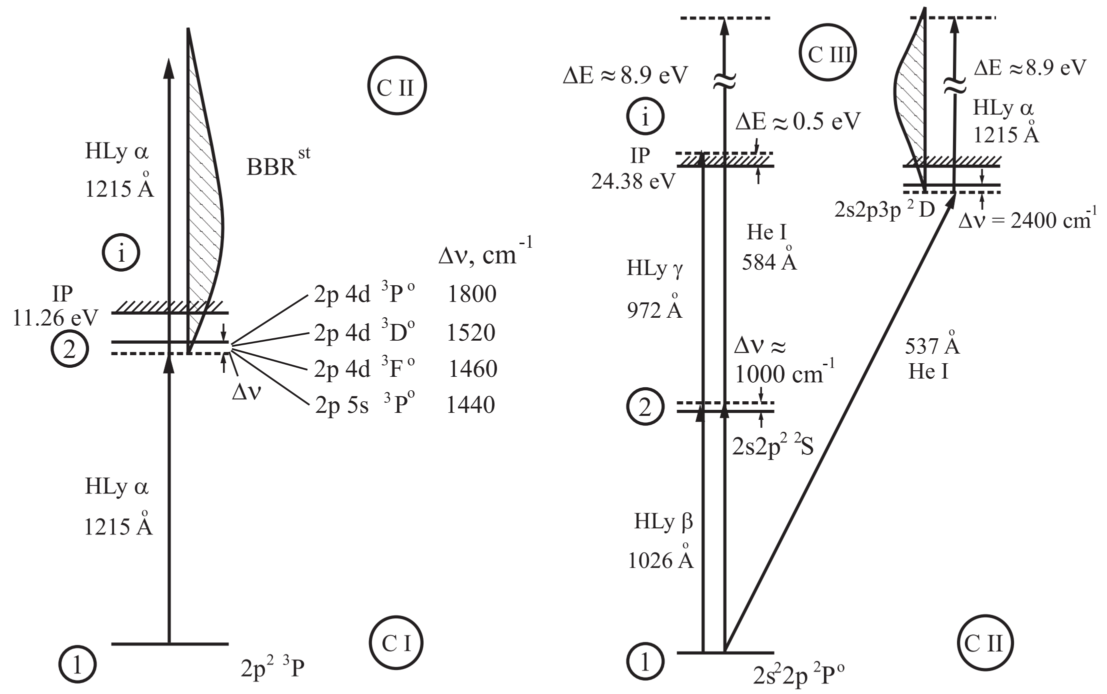

Fig. 3. Chain of successive RETPI schemes for carbon: CI→CII (left) and CII→CIII (right); For CIII and CIV, see Fig. 4.

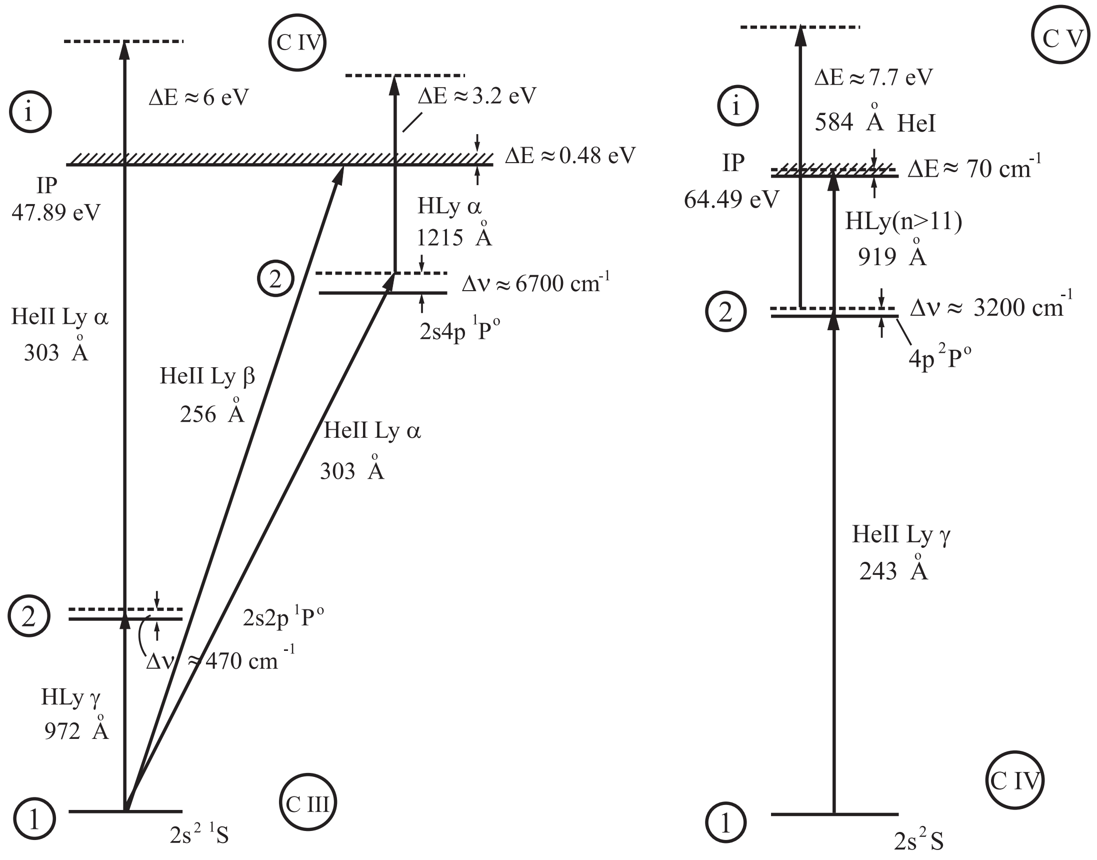

Fig. 4. Chain of successive RETPI schemes for carbon: CIII→CIV (left) and CIV→CV (right); For CI and CII, see Fig. 3.

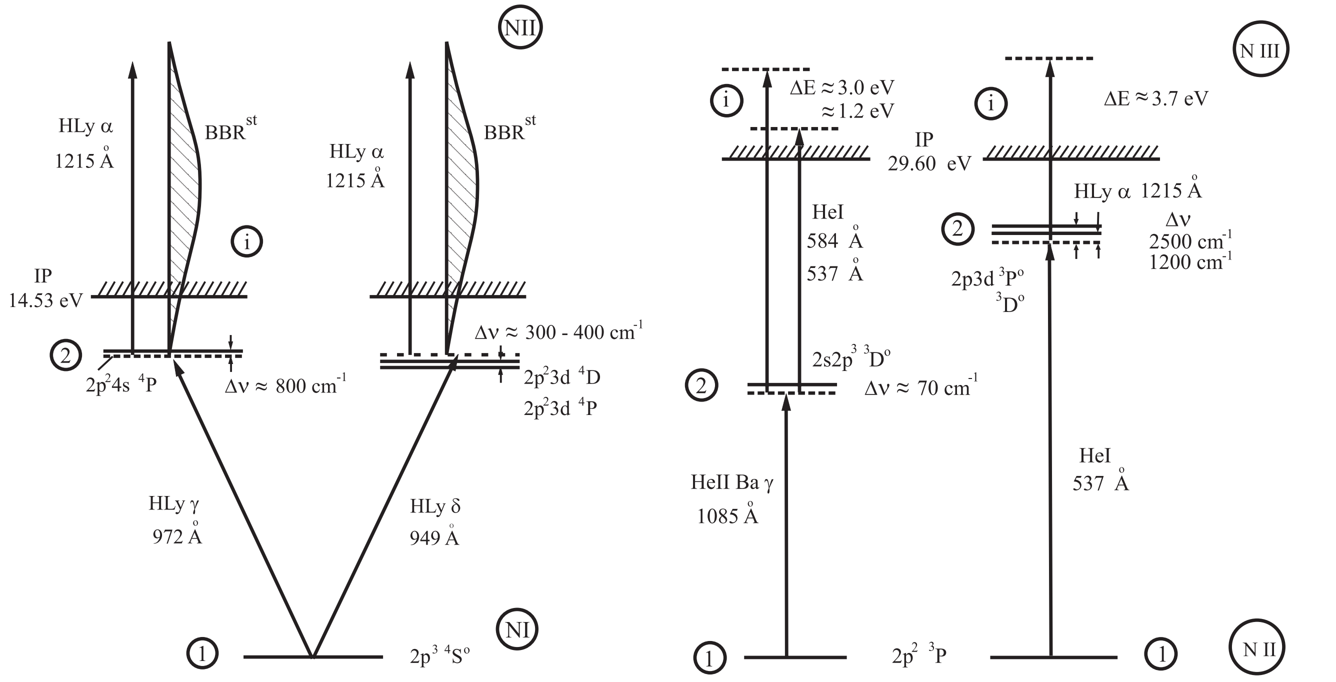

Fig. 5. Chain of successive RETPI schemes for nitrogen: NI→NII (left) and NII→NIII (right); For NIII and NIV, see Fig. 6.

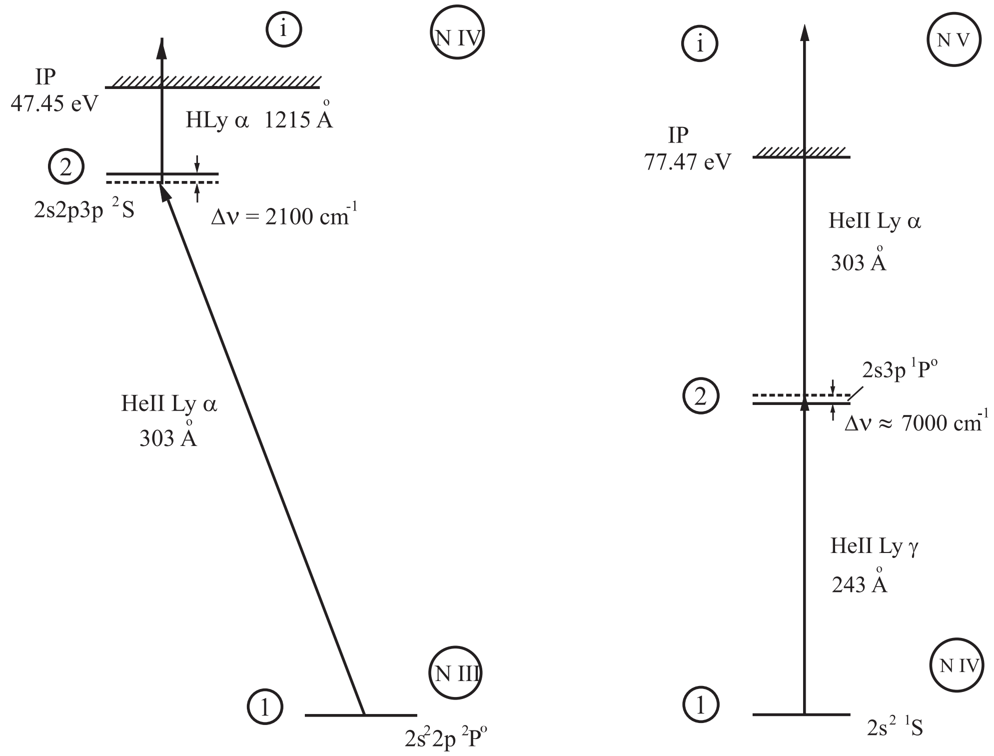

Fig. 6. Chain of successive RETPI schemes for nitrogen: NIII→NIV (left) and NIV→NV (right); For NI and NII, see Fig. 5.

governed by the rate of the RETPI process in an appropriate excitation scheme and not by electron collisions. In that case, the determination of the electron temperature from a comparison between the spectral line intensities of such neighbouring ions may prove incorrect.

In conclusion, we would like to emphasize that RETPI is a photonic mechanism with no requirement of any collisions. For this reason, it is especially effective in nebulae with low electron densities, and with sizes and concentrations of HI, HII, HeI, and HeII that are great enough

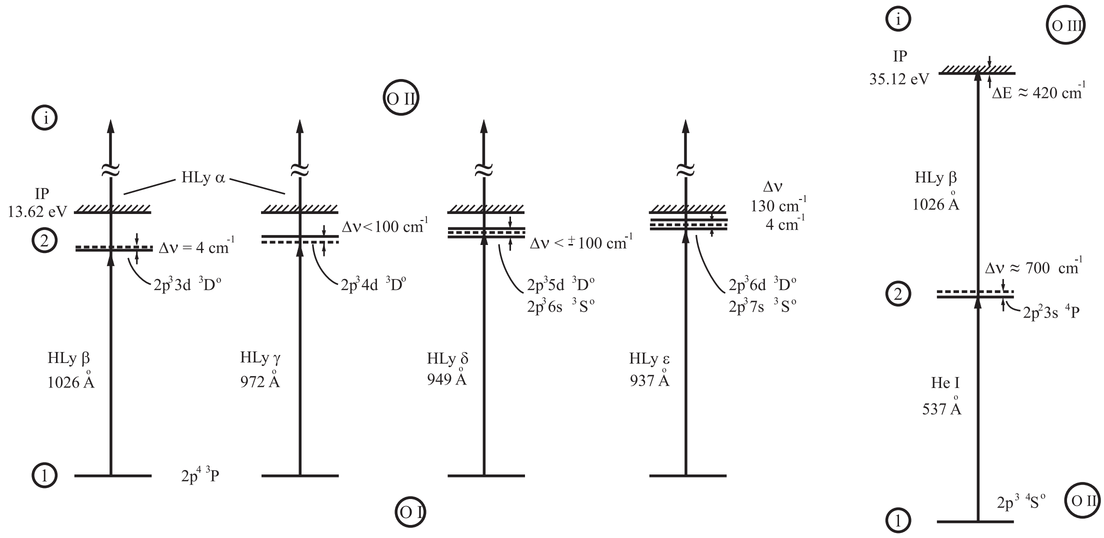

Fig. 7. Chain of successive RETPI schemes for oxygen: OI→OII (left) and OII→OIII (right); For OIII and OIV, see Fig. 8.

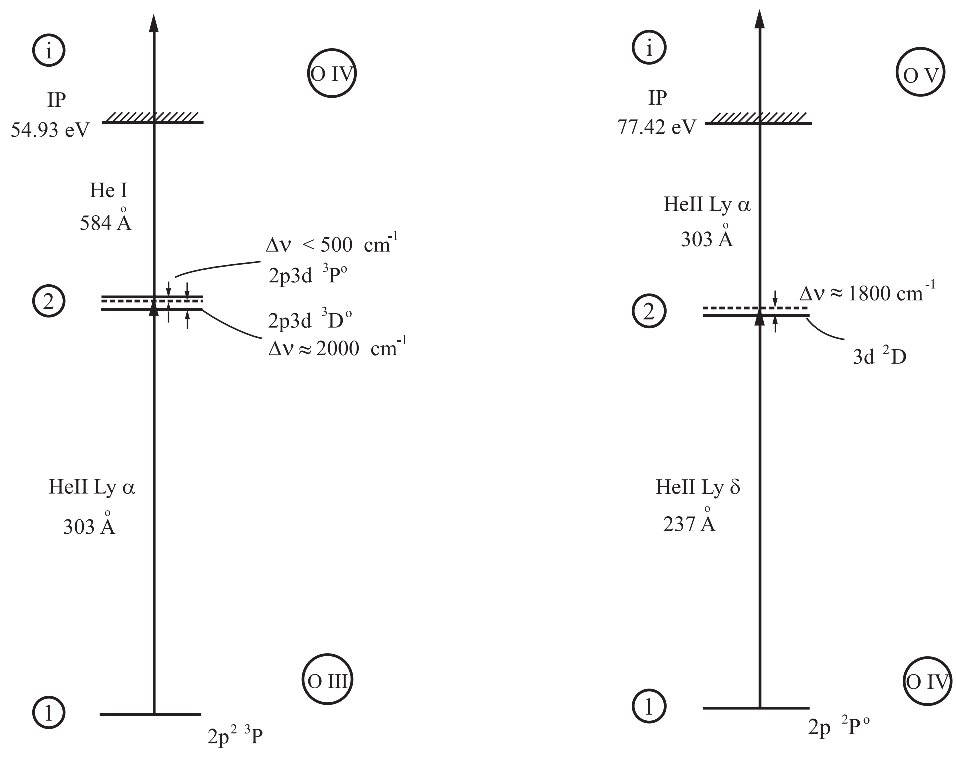

Fig. 8. Chain of successive RETPI schemes for oxygen: OIII→OIV (left) and OIV→OV (right); For OI and OII, see Fig. 7.

to form strong optically thick emission lines of HI, HeI, and HeII. The necessary combinations of parameters (volume, concentration, etc.) that can provide the very high intensity of these spectral lines for efficient RETPI *inside* nebulae will be the subject for future research, as well as the effect that RETPI causes on the ionization balance in individual nebulae. Such problems will be natural applications of the newly proposed mechanism in planetary nebulae. The high efficiency of RETPI by the intrinsic EUV radiation trapped in nebulae compared to that of electronic collision ionization is a result of the higher density and velocity of photons compared to electrons.

Acknowledgements. One of the authors (V. S. L.) is thankful to Dr. A. A. Makarov for valuable discussions of the problem. We are also grateful to Dr. T. Gull for useful comments on the manuscript. This project is carried out when V. S. L. is holding the Tage Erlander Guest Professorship at Lund University, and it is supported by grants (S. J.) from the Swedish Natural Science Research Council and the Swedish National Space Board.

Table 1. Chains of subsequent schemes of RETPI for C, N, and O activated by strong EUV lines of HI, HeI and HeII.

Ele-ment	**Pairs of Photoionizing Spectral Lines of HI, HeI and HeII**				
	HI + HI	HI + HeI	HI + HeII	HeI + HeII	HeII + HeII
C	CI→CII (11.26 eV) 1215 Å + 1215 Å CII→CIII 1026 Å + 972 Å	CII→CIII (24.38 eV) 1026 Å + 584 Å 537 Å + 1215 Å	CIII→CIV (47.89 eV) 972 Å + 303 Å 303 Å + 1215 Å CIV→CV 243 Å + 919 Å	CIV→CV (64.49 eV) 243 Å + 584 Å	
N	NI→NII (14.53 eV) 972 Å + 1215 Å 949 Å + 1215 Å	NII→NIII (29.60 eV) 537 Å + 1215 Å	NIII→NIV (47.45 eV) 303 Å + 1215 Å	NII→NIII (29.60 eV) 1085 Å + 584 Å	NIV→NV (77.47 eV) 243 Å + 303 Å
O	OI→OII (13.62 eV) 1026 Å + 1215 Å (972 Å) + 1215 Å (949 Å) + 1215 Å (937 Å) + 1215 Å	OII→OIII (35.12 eV) 537 Å + 1026 Å		OIII→OIV (54.93 eV) 303 Å + 584 Å	OIV→OV (77.41 eV) 237 Å + 303 Å

References

Aller, L. H. 1956, Gaseous Nebulae (Chapman and Hall, Ltd., London)

Aller, L. H. 1984, Physics of Thermal Gaseous Nebulae (D. Riedel Publ. Co., Dordrecht)

Bashkin, S., & Stoner Jr, J. O. 1975, Atomic Energy Levels and Grotrian Diagrams, vol. 1 (North-Holland, Amsterdam)

Bowen, I. S. 1935, ApJ, 81, 1

Dinerstein, H. L., Lafon, C. E., & Garnett, D. R. 2000, in Assymetrical Planetary Nebulae II: From Origins to Microstructure, ed. T. H. Kastner, N. Soker, & S. Rappaport, ASP Conf. Ser., 199, 301

Johansson, S., & Hamann, F. W. 1993, Phys. Scr., T47, 157

Johansson, S., & Letokhov, V. S. 2001a, in Eta Carinae and other Mysterious Stars, ed. T. R. Gull, S. Johansson, & K. Davidson, ASP Conf. Ser., 242, 309

Johansson, S., & Letokhov, V. S. 2001b, submitted to A&A

Johansson, S., & Letokhov, V. S. 2001c, Science, 291, 625

Johansson, S., Davidson, K., Ebbets, D., et al. 1996, in Science with Hubble Telescope-II, ed. P. Benvenuti, F. D. Macehetto, & F. J. Schreier, Space Telescope Science Institute, 361

Johansson, S., Zethson, T., Hartman, H., et al. 2000, A&A, 361, 977

Landau, L. D., & Lifshitz, E. M. 1958, Quantum Mechanics (Addison-Wesley, Reading, Mass)

Letokhov, V. S. 1987, Laser Photoionization Spectroscopy (Orlando, Academic Press Inc)

Makarov, A. A. 1983, Zh. Eksp. Teor. Fiz., 85, 1192

Osterbrock, D. E. 1989, Astrophysics of Gaseous Nebulae and Active Galactic Nuclei (Univ. Science Books, Sausalito, California)

Pottasch, S. R. 1984, Planetary Nebulae (D. Reidel Publ. Co., Dordrecht)

Sobel'man, I. I. 1979, Atomic Spectra and Radiative Transitions (Springer, Berlin)

Weisskopf, V. 1933, Z. Phys., 85, 451

Zanstra, H. 1949, Bull. Astr. Insts. Neth., 11, 1