Effective collision strengths for ground-state and 2s<sup>2</sup>2p–2s2p<sup>2</sup>
fine-structure transitions in C II

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Abstract

The ab initio R-matrix method is used to calculate effective collision strengths for electron-impact excitation of the boron-like ion C II. We consider the fine-structure forbidden \(^2P_{1/2}^o - 2P_{3/2}^o\) transition within the 2s<sup>2</sup>2p ground configuration as well as the 16 fine-structure transitions arising from excitation from the 2s<sup>2</sup>2p to the 2s2p<sup>2</sup> configuration. The 16 lowest \(L_S\) target states are included in the calculation. Effective collision strengths are obtained by averaging the electron collision strengths over a Maxwellian distribution of electron velocities. Results are presented for electron temperatures in the range \(\log T(K) = 3.0 - 5.5\), appropriate for astrophysical applications. Good agreement with previous evaluations is obtained.

Key words: atomic processes – line: formation.

1 INTRODUCTION

Carbon is one of the most astrophysically important elements, occurring as both neutral and ionized species. Singly ionized carbon is the predominant form of carbon throughout \(\text{HII}\) regions and also in the outer regions of molecular clouds. The 2s<sup>2</sup>2p–2s2p<sup>2</sup> transition has been observed along the ridge of star formation in NGC 6334 (Boreiko & Betz 1995). Lines arising from the 2s<sup>2</sup>2p–2s2p<sup>2</sup> transition were included. Effective collision strengths for a range of temperatures were presented for transitions among the 2s<sup>2</sup>2p–2s2p<sup>2</sup> configuration. The 16 lowest \(L_S\) target states were included. However, results reported were restricted to \(L_S\) transitions. The fine-structure transition within the ground state was considered by Blum & Pradhan (1991). The target states employed were the same as those used by Luo & Pradhan but some refinements were applied, notably the use of observed target state energies to improve resonance positioning, and the use of a more complete partial-wave expansion. Effective collision strengths for all fine-structure transitions arising among the 10 \(L_S\) target states are reported by Blum & Pradhan (1992) for a range of electron temperatures.

The availability of improved computing capabilities enables more sophisticated calculations to be undertaken. The aim of the present work is to perform such a calculation, by including an appropriate set of target states and considering carefully any resonance structure. We will thus be able to calculate accurate and reliable collision strengths and hence effective collision strengths for a range of temperatures of interest to astrophysicists.

2 THE METHOD

We used the general configuration-interaction code \texttt{civ3} (Hibbert 1975) to calculate wavefunctions for the C II ion in \(L_S\) coupling. This is achieved by expressing each state \(\Psi\) as a linear combination of single-configuration functions \(\Phi_i\) possessing the same total \(L_S\pi\) symmetry:

\[
\Psi(LS) = \sum_{i=1}^{m} a_i \Phi_i(\alpha_i, LS).
\]

The mixing coefficients \(a_i\) are the eigenvector components of the Hamiltonian matrix with particular \(L_S\pi\) symmetry, which has elements

\[a = \text{ai}\]
where $H$ denotes the Hamiltonian operator. The corresponding eigenvalues are upper bounds to exact energies of the ionic states. $\alpha_i$ represent the coupling of the angular momenta associated with the one-electron spin orbitals to form a total $L$ and $S$ and, if necessary, additional quantum numbers required to characterize the configuration state function $\Phi_i$. We build the $\{\Phi_i\}$ from a set of one-electron functions, each consisting of a product of a radial function, a spherical harmonic and a spin function:

$$u_{nlm\mu}(r, \sigma) = \frac{1}{r} P_{nl}(r) Y^m_l(\theta, \phi) \chi_{nl\mu}(\sigma).$$

(3)

We choose these orbitals to be analytic in our calculations and express the radial part as a linear combination of Slater-type orbitals

$$P_{nl}(r) = \sum_{jnl} c_{jnl} \chi_{jnl}(r),$$

(4)

where

$$\chi_{jnl}(r) = \frac{[2\zeta_{jnl}]^{1/2}}{[2(I_j n_l)^{1/2}]} \exp(-\zeta_{jnl} r).$$

(5)

The powers of $r$, $|I_{jnl}|$, are held constant, but the coefficients $\{c_{jnl}\}$ and the exponents $\{\zeta_{jnl}\}$ are treated as variational parameters when determining the radial functions. We include 14 orbitals in the present calculation, 10 'spectroscopic' (1s, 2s, 2p, 3s, 3p, 3d, 4s, 4p, 4d, 4f) and four pseudo-orbitals ($\bar{3}s$, $\bar{3}p$, $\bar{3}d$, $\bar{3}f$). The $c_{jnl}$, $I_{jnl}$ and $\zeta_{jnl}$ parameters for the 1s orbital were taken to be the Hartree–Fock values calculated by Clementi & Roetti (1974) for the $2s^22p^2$ ground state. The 2s and 2p functions were also taken from this ground state. The $2s$ and $2p$ functions were also taken from this ground state but were re-optimized on the energy of the $2s 2p^2 2P^o$ state. The largest difference between this calculation and the observed values is excellent and so we consider the wavefunctions to be sufficiently sophisticated. We also compare with the energies obtained by Luo & Pradhan and Lennon et al. In particular, we note that the present energy levels are in much better accord with the observed values than the work of Luo & Pradhan or Lennon et al. The largest difference between this calculation and the observed values is less than 1.5 per cent; for many of the states the agreement is considerably better than this. We also note that in addition to being a larger calculation (16 states compared with ten and eight by Luo & Pradhan and Lennon et al., respectively), we include all 16 lowest LS states of C II. This ensures that the collision strength spectrum to be obtained contains all the autoionizing resonances converging to these thresholds.

<table>
<thead>
<tr>
<th>Function</th>
<th>$c_{jnl}$</th>
<th>$I_{jnl}$</th>
<th>$\zeta_{jnl}$</th>
<th>$c_{jnl}$</th>
<th>$I_{jnl}$</th>
<th>$\zeta_{jnl}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>1s</td>
<td>0.06002</td>
<td>1</td>
<td>9.68937</td>
<td>3s</td>
<td>2.85189</td>
<td>1</td>
</tr>
<tr>
<td></td>
<td>0.92321</td>
<td>1</td>
<td>5.54099</td>
<td>-15.33621</td>
<td>2</td>
<td>1.32664</td>
</tr>
<tr>
<td></td>
<td>0.02758</td>
<td>2</td>
<td>4.96979</td>
<td>30.44600</td>
<td>3</td>
<td>1.32664</td>
</tr>
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<td>2.14468</td>
<td>-26.63725</td>
<td>4</td>
<td>1.32664</td>
</tr>
<tr>
<td></td>
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<td>1.56842</td>
<td>8.69986</td>
<td>5</td>
<td>1.32664</td>
</tr>
<tr>
<td></td>
<td>-0.01585</td>
<td>1</td>
<td>9.31849</td>
<td>2p</td>
<td>0.05211</td>
<td>2</td>
</tr>
<tr>
<td></td>
<td>-0.21133</td>
<td>1</td>
<td>5.55500</td>
<td>0.10385</td>
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<td>2.41548</td>
</tr>
<tr>
<td></td>
<td>-0.10868</td>
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<td></td>
<td>0.48786</td>
<td>2</td>
<td>2.15137</td>
<td>0.60002</td>
<td>2</td>
<td>1.26824</td>
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<td></td>
<td>0.60703</td>
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<td>1.61952</td>
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<td>1.26824</td>
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<tr>
<td>3s</td>
<td>0.12878</td>
<td>1</td>
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<td>3p</td>
<td>0.03907</td>
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<tr>
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<td>2</td>
<td>1.72327</td>
<td>0.31060</td>
<td>3</td>
<td>1.06969</td>
</tr>
<tr>
<td></td>
<td>-0.15499</td>
<td>3</td>
<td>0.89922</td>
<td>-1.33030</td>
<td>3</td>
<td>0.79918</td>
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<tr>
<td></td>
<td>0.07421</td>
<td>1</td>
<td>4.68036</td>
<td>4p</td>
<td>0.27656</td>
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<td></td>
<td>-0.33778</td>
<td>2</td>
<td>1.72817</td>
<td>-2.16498</td>
<td>3</td>
<td>0.55866</td>
</tr>
<tr>
<td></td>
<td>2.43878</td>
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<td>2.65828</td>
<td>4</td>
<td>0.55245</td>
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<td>-2.89178</td>
<td>4</td>
<td>0.64168</td>
<td>0.49985</td>
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<td>2.35168</td>
</tr>
</tbody>
</table>

Table 1. The orbital parameters ($c_{jnl}$, $I_{jnl}$, $\zeta_{jnl}$) of the radial wavefunctions.
The \( R \)-matrix method, as described by Burke & Robb (1975) and Seaton (1987) was employed in the electron-scattering calculations, using the associated computer codes described by Berrington et al. (1987). We included 40 Schmidt-orthogonalized continuum orbitals for each continuum orbital \( \alpha \), ensuring a converged \( R \)-matrix was obtained up to the highest electron energy considered (10 Ryd). The \( R \)-matrix radius was taken to be 28.8 au. The \( LS \)-coupled Hamiltonian matrices were adjusted so that their theoretical term energies matched the observed values from the NIST database. The adjustment of energies in this way ensures the correct positioning of resonances relative to the thresholds included in the calculation. We note that the energy order of the two highest states included in this work differs from that given by the NIST data base. The energy levels for these two states were thus both adjusted to an energy of 1.5384 Ryd, the arithmetic mean of the two NIST values.

All total angular momenta up to and including \( L = 12 \) were included to form \( (N + 1) \)-electron symmetries with singlet, triplet and quintet multiplicities, for both even and odd parities. Consequently, for the optically forbidden transitions considered in this work, convergence of the collision strengths was fully achieved. However, for optically allowed transitions, it is necessary to consider the effect of higher partial waves with \( L \geq 12 \) since they significantly affect the collision strengths. We assume that the partial collision strengths form a geometric series expansion, with a geometric scaling factor given by the ratio of two adjacent terms. This ‘topping-up’ procedure has been used successfully in similar calculations (see, for example, Bell & Ramsbottom 2000; Ramsbottom, Bell & Keenan 2001).

The configurations describing these \( (N + 1) \)-electron symmetries were generated by the addition of one electron to those configurations generated by the replacement of two electrons from the \( n = 2 \) complex in the electron distributions \( 2s^22p \) and \( 2s2p^2 \). Having obtained the \( K \)-matrices within the framework of the \( LS \)-coupling scheme, we utilize the \textsc{jaom} program of Saraph (1978) to transform to \( LSJ \) intermediate-coupling including fine-structure mixing of the target terms and thus produce collision strengths, \( \Omega \), between the \( J \)-resolved levels (McLaughlin & Bell 2000) provide fuller details of the transformation.

For excitation from level \( i \) to level \( f \), a (dimensionless) effective collision strength \( \gamma_i \), in terms of \( \Omega \), at electron temperature \( T_e \) (in kelvin) is given by

\[
\gamma_i(T_e) = \int_0^\infty \Omega_i(E) \exp(-E_i/kT_e) \, dE_i / kT_e, \tag{6}
\]

where \( E_i \) is the final free electron energy after excitation and \( k \) is Boltzmann’s constant.

In order for the Maxwellian-averaged effective collision strengths to be evaluated accurately, it is essential that the energy dependence of the collision strength is known over a wide energy range and for a large number of energy points. It is imperative that the complex autoionizing resonances, which converge to the target-state thresholds are fully delineated, since the presence of these structures can have a significant effect upon the effective collision strengths (see, for example, Ramsbottom, Bell & Stafford 1996). In consideration of this, we have utilized a fine mesh of incident impact energies, typically 0.0002 Ryd across the energy range 0–10 Ryd. Pseudo-resonances, which arise as a result of the inclusion of pseudo-orbitals in the wavefunction representation (Burke, Sukumar & Berrington 1981) have been smoothed out for energies above the highest-lying threshold included in this work \([2s^24f^25p^0]\) using a cubic spline fit. This procedure ensures that there is no distortion of the effective collision strength results in the high-energy (and hence high-temperature) region.

### 3 Results and Discussion

#### 3.1 The \( 2s^22p_1^23P_{1/2} \rightarrow 2s^22p^23P_{3/2} \) transition

In Fig. 1(a) we present the collision strength as a function of the incident electron energy (in Ryd) for the \( 2s^22p_1^23P_{1/2} \rightarrow 2s^22p^23P_{3/2} \) ground-state fine-structure transition. The resonance structure has been fully delineated by the fine energy mesh used. Of particular note is the broad resonance, which has a peak at 0.156 Ryd. Using the \( R \)-matrix eigenvectors, this resonance is caused by the \( 2s2p^31D^0 \) state of \( \text{Cl} \). The position of this resonance has been the subject of some debate in previous studies. The prediction by Edlén (1934) that this state should lie at 0.06 Ryd above the ionization limit has led to

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Table 2. Target state energies (in Ryd) relative to the \( 3s^23p^43P \) ground state of \( \text{Cl} \) compared with values from the NIST data base, Luo & Pradhan (1990) and Lennon et al. (1985).

<table>
<thead>
<tr>
<th>( \text{Cl} ) state</th>
<th>Present ( LS ) energy</th>
<th>NIST energy</th>
<th>Luo &amp; Pradhan</th>
<th>Lennon et al.</th>
<th>Number of configurations</th>
</tr>
</thead>
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<tr>
<td>( 2s^22p_2^12P_0 )</td>
<td>0.0000</td>
<td>0.0000</td>
<td>0.0000</td>
<td>0.0000</td>
<td>181</td>
</tr>
<tr>
<td>( 2s^22p_2^32p_1 )</td>
<td>0.3873</td>
<td>0.3918</td>
<td>0.3790</td>
<td>0.3670</td>
<td>64</td>
</tr>
<tr>
<td>( 2s^22p_2^2D_0 )</td>
<td>0.6885</td>
<td>0.6828</td>
<td>0.6917</td>
<td>0.6954</td>
<td>198</td>
</tr>
<tr>
<td>( 2s^22p_2^3D_0 )</td>
<td>0.8892</td>
<td>0.8789</td>
<td>0.9011</td>
<td>0.9004</td>
<td>95</td>
</tr>
<tr>
<td>( 2s^22p_2^3P_0 )</td>
<td>1.0228</td>
<td>1.0080</td>
<td>1.0284</td>
<td>1.0426</td>
<td>116</td>
</tr>
<tr>
<td>( 2s^23p^23P_0 )</td>
<td>1.0627</td>
<td>1.0616</td>
<td>1.0679</td>
<td>1.0706</td>
<td>95</td>
</tr>
<tr>
<td>( 2s^23p^23P_1 )</td>
<td>1.1991</td>
<td>1.2000</td>
<td>1.2060</td>
<td>1.2104</td>
<td>118</td>
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<tr>
<td>( 2p^34S_0 )</td>
<td>1.2972</td>
<td>1.2939</td>
<td>1.3244</td>
<td>–</td>
<td>19</td>
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<tr>
<td>( 2p^34D_0 )</td>
<td>1.3269</td>
<td>1.3260</td>
<td>–</td>
<td>1.3532</td>
<td>108</td>
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<tr>
<td>( 2p^34F_0 )</td>
<td>1.4310</td>
<td>1.4324</td>
<td>–</td>
<td>–</td>
<td>95</td>
</tr>
<tr>
<td>( 2s^24s^22S_0 )</td>
<td>1.4801</td>
<td>1.4806</td>
<td>–</td>
<td>–</td>
<td>181</td>
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<tr>
<td>( 2s^24p^23P_0 )</td>
<td>1.5214</td>
<td>1.5215</td>
<td>–</td>
<td>–</td>
<td>80</td>
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<tr>
<td>( 2s^24f^21D_1 )</td>
<td>1.5298</td>
<td>1.5317</td>
<td>–</td>
<td>–</td>
<td>198</td>
</tr>
<tr>
<td>( 2p^33P_0 )</td>
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<td>1.5373</td>
<td>1.5785</td>
<td>–</td>
<td>180</td>
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<tr>
<td>( 2s^24f^23P_0 )</td>
<td>1.5364</td>
<td>1.5395</td>
<td>–</td>
<td>–</td>
<td>161</td>
</tr>
</tbody>
</table>
some authors recommending that the resonances should be shifted so that the strongest resonance lies at its observationally predicted position. Such a procedure leads to the calculation of effective collision strengths, which are very different from the values obtained from an unadjusted spectrum (Hayes & Nussbaumer 1984a; Keenan et al. 1986).

In the absence of the confirmation of the position predicted by Edlen by direct spectroscopic observation, we have performed a superposition-of-configurations calculation using the \textsc{civ3} computer program for the $^1\text{D}^\pi$ states of C$_1$. The ionization potential for C$_1$ is known, and so we can calculate the energy above the ionization potential at which the $2s2p^3$$^1\text{D}^\pi$ state of C$_1$ lies, since the energy difference between the ionization threshold and the lowest $^1\text{D}^\pi$ state ($2s^22p^3d$) is also known.

An appropriate initial set of orbitals was chosen. The 1s and 2s functions were taken to be the Hartree–Fock functions for the C$_1$ ground state of Clementi & Roetti. The real (2p, 3d, 4s) and corrector (3s, 3p, 5d) orbitals were obtained using the \textsc{civ3} code by optimizing on appropriate states of C$_1$. We considered replacement of two electrons in the $2s2p^3$ configuration with any two from the orbital set to generate the configuration-interaction wavefunctions. The energies of the $^1\text{D}^\pi$ states are given in Table 3. The energy differences for the $2s^22p^4d^1\text{D}^\pi$ state and the ionization threshold, relative to the $2s^22p^3d^1\text{D}^\pi$ state from the NIST data base are 0.0530 and 0.1201 Ryd, respectively.

Further pseudo-orbitals were now included in the calculation and their effect on the energies is shown in Table 3. The use of this large orbital set has enabled us to try to replicate the main channel-coupling effects that are in the R-matrix expansion. This ‘final’ energy of 0.1219 Ryd falls within the broad resonance structure observed in Fig. 1(a) and would clearly indicate that the Edlen energy value of 0.06 Ryd is in error. We therefore conclude that the unadjusted collision strength spectrum should be used to calculate the Maxwellian-averaged collision strengths.

In Table 4 we present effective collision strengths for the ground-state fine-structure transition at temperatures of 2000, 10 000 and 20 000 K and compare with previous determinations. Whilst Hayes & Nussbaumer (1984a) and Keenan et al. advised that the effective collision strengths obtained following adjustment (see above for more details) are more accurate than those obtained without adjustment, Lennon et al. and Blum & Pradhan (1991) consider no

\begin{table}[h]
\centering
\begin{tabular}{lcccc}
\hline
State & $2s^22p^3d^1\text{D}^\pi$ & $2s^22p^4d^1\text{D}^\pi$ & $2s^22p^3_{\text{f}}^1\text{D}^\pi$ & $2s^22p^3^3\text{D}^\pi$ \\
\hline
Initial orbital set & -74.8319 & 0.05171 & 0.2862 & 0.1661 \\
+6d & -74.8323 & 0.05167 & 0.2787 & 0.1586 \\
+7d & -74.8324 & 0.05167 & 0.2776 & 0.1575 \\
+8d & -74.8336 & 0.05168 & 0.2769 & 0.1568 \\
+4s & -74.8327 & 0.05167 & 0.2759 & 0.1558 \\
+5s & -74.8330 & 0.05177 & 0.2627 & 0.1426 \\
+6s & -74.8332 & 0.05177 & 0.2624 & 0.1423 \\
+4f & -74.8338 & 0.05184 & 0.2439 & 0.1238 \\
+5f & -74.8339 & 0.05187 & 0.2420 & 0.1219 \\
\hline
\end{tabular}
\caption{Energies in Ryd of $^1\text{D}^\pi$ states of C$_1$ for different orbital sets.}
\footnotesize{aAbsolute energy.  \\
bRelative to $2s^22p^3^1\text{D}^\pi$.  \\
cRelative to ionization threshold.}
\label{tab:energies}
\end{table}

\begin{table}[h]
\centering
\begin{tabular}{lccc}
\hline
State & $2s^22p^3d^1\text{D}^\pi$ & $2s^22p^4d^1\text{D}^\pi$ & $2s^22p^3_{\text{f}}^1\text{D}^\pi$ \\
\hline
&  &  &  \\
\hline
Present work & 1.79 & 2.24 & 2.42 \\
Blum & Pradhan & 1.64 & 2.15 & 2.28 \\
Keenan et al.$^a$ & 1.63 & 2.11 & 2.28 \\
Keenan et al.$^b$ & 2.30 & 2.76 & 2.64 \\
Lennon et al. & 1.63 & 2.11 & 2.28 \\
Hayes & Nussbaumer$^a$ & 1.71 & 2.44 & 2.57 \\
Hayes & Nussbaumer$^b$ & 2.03 & 2.90 & 2.82 \\
\hline
\end{tabular}
\caption{A comparison of the present calculated effective collision strengths for the $2s^22p^3^3\text{D}^\pi$–$2s^22p^3_{\text{f}}^1\text{D}^\pi$ fine-structure transition with the previous evaluations of Blum & Pradhan (1991), Keenan et al. (1986), Lennon et al. (1985) and Hayes & Nussbaumer (1984a).}
\footnotesize{$^a$Ab initio calculation.  \\
$^b$Resonances shifted to Edlen (1934).}
\label{tab:collision_strengths}
\end{table}

Figure 1. (a) Collision strength as a function of incident electron energy in Ryd and (b) the effective collision strengths as a function of log temperature in K for the $2s^22p^2P^1_{1/2}$–$2s^22p^2P^1_{3/2}$ fine-structure transition [solid line: present results, dashed line: Blum & Pradhan (1992)].
adjustment to be necessary. The present authors concur with this assessment and therefore consider a comparison with the unadjusted values to be more appropriate.

For each of the three temperatures considered, the present effective collision strengths are, in the main, higher than the unadjusted results of other authors. At 2000 K the effective collision strength is more than 9 per cent higher than that of Blum & Pradhan. This is likely to be caused by the inclusion of a greater number of target states and the consequent increased amount of resonant structure. At higher temperatures, the difference is smaller. We note that at 10 000 and 20 000 K the present results are lower than the unadjusted results of Hayes & Nussbaumer. However, their work included a smaller number of configurations and used less accurate wavefunctions (energy separations within 3 per cent compared with 1.5 per cent in the present work). We are thus confident that the present results are more accurate.

In Fig. 1(b) we compare the effective collision strength graphically for temperatures in the range given by log $T$ (K) = 3.0–5.5 with the results of Blum & Pradhan. Their results are the most recent and are in good agreement with the previous (unadjusted) calculations of Keenan et al. and Lennon et al. For the temperature range for which a comparison is possible, the present effective collision strengths are higher than those of Blum & Pradhan. The graph shows a gradual increase up to a temperature given by log $T$ (K) = 4.25, after which the effective collision strength begins to decrease rapidly.

### 3.2 The 2s2p$^3$P$^o$–2s2p$^1$P$^o$ fine-structure transitions

The 2s$^2$2p$^3$P$^o$ and 2s2p$^1$P$^o$ states give rise to six fine-structure forbidden transitions. As examples, we present in Figs 2 and 3 the collision strength and corresponding Maxwellian-averaged collision strength for the $^2\Pi_{1/2}$–$^2\Pi_{1/2}$ and $^2\Pi_{3/2}$–$^2\Pi_{5/2}$ transitions, respectively.

![Figure 2](https://example.com/figure2.png)

**Figure 2.** (a) Collision strength as a function of incident electron energy in Ryd and (b) the effective collision strengths as a function of log temperature in K for the 2s$^2$2p$^3$P$^o$–2s2p$^1$P$^o$ fine-structure transition [solid line: present results, dashed line: Blum & Pradhan (1992)].

For each of the transitions, rather broad resonance structures can be observed in the incident electron energy range 0.4–0.6 Ryd. Beyond this energy, however, narrower autoionizing resonances converging to the thresholds included in this work appear in the spectrum. Care has been taken to ensure that these resonances have been fully resolved, as they have the effect of enhancing the effective collision strength.

The present effective collision strength data are presented, together with the results of Blum & Pradhan (1992) in Figs 2(b) and 3(b). For these two transitions, the effective collision strength decreases only very slightly in the temperature range given by log $T$ (K) = 3.0–4.5 before falling off rapidly towards higher temperatures. This sharp decrease in the effective collision strength is characteristic of forbidden transitions. A comparison with the work of Blum & Pradhan reveals that for the temperature range considered by them, the present results are up to 17 per cent higher. We also note that the shallow dip in effective collision strength followed by a small increase is not observed in this work.

A numerical comparison for all six transitions of the present Maxwellian-averaged collision strengths at temperatures of 2000, 10 000 and 20 000 K with previous evaluations is given in rows 1–6 of Table 5. This reveals that the results obtained in the present work are higher than those from previous calculations at temperatures of 2000 and 10 000 K. We note, though, that the results are slightly lower than those of Hayes & Nussbaumer (1984b) at 20 000 K. However, given the accuracy of the wavefunctions employed and the careful consideration of resonances, we are confident in the accuracy of the present results.

### 3.3 The 2s$^2$2p$^3$P$^o$–2s2p$^1$D$^o$ fine-structure transitions

Examples of the collision strength data for the fine-structure transitions between the 2s$^2$2p$^3$P$^o$ and 2s2p$^1$D$^o$ states are presented in...
Lennon et al. (1985, L) and Hayes & Nussbaumer (1984b, HN) at specific strength is given in Fig. 4(b) and is compared with the work the target states lie (0

\begin{table}
\centering
\begin{tabular}{|c|c|c|c|c|c|c|c|c|c|}
\hline
Transition & \multicolumn{3}{|c|}{This} & \multicolumn{3}{|c|}{BP} & \multicolumn{3}{|c|}{HN} \\
\hline & \multicolumn{1}{|c|}{T = 2000 K} & \multicolumn{1}{|c|}{T = 10000 K} & \multicolumn{1}{|c|}{T = 20000 K} \\
\hline \(2p^3P_{1/2} - 2p^3P_{1/2}\) & 0.291 & 0.258 & 0.288 & 0.280 & 0.282 & 0.243 & 0.280 & 0.276 & 0.278 \\
\hline \(2p^3P_{1/2} - 2p^3P_{1/2}\) & 0.429 & 0.382 & 0.424 & 0.412 & 0.419 & 0.362 & 0.413 & 0.409 & 0.414 \\
\hline \(2p^3P_{1/2} - 2p^3P_{1/2}\) & 0.263 & 0.240 & 0.253 & 0.249 & 0.265 & 0.235 & 0.253 & 0.260 & 0.267 \\
\hline \(2p^3P_{1/2} - 2p^3P_{1/2}\) & 0.200 & 0.182 & 0.194 & 0.191 & 0.201 & 0.177 & 0.193 & 0.197 & 0.202 \\
\hline \(2p^3P_{1/2} - 2p^3P_{1/2}\) & 0.553 & 0.497 & 0.542 & 0.529 & 0.547 & 0.477 & 0.533 & 0.536 & 0.545 \\
\hline \(2p^3P_{1/2} - 2p^3P_{1/2}\) & 1.21 & 1.08 & 1.20 & 1.16 & 1.18 & 1.02 & 1.17 & 1.16 & 1.17 \\
\hline \(2p^3P_{1/2} - 2s_1^3/2\) & 1.15 & 1.29 & 1.12 & 1.23 & 1.37 & 1.40 & 1.27 & 1.36 & 1.51 \\
\hline \(2p^3P_{1/2} - 2s_1^3/2\) & 0.577 & 0.514 & 0.659 & 0.556 & 0.610 & 0.543 & 0.623 & 0.561 & 0.608 \\
\hline \(2p^3P_{1/2} - 2s_1^3/2\) & 0.922 & 0.874 & 1.03 & 0.913 & 1.01 & 0.931 & 1.01 & 0.942 & 1.03 \\
\hline \(2p^3P_{1/2} - 2s_1^3/2\) & 2.54 & 2.73 & 2.17 & 2.66 & 2.95 & 2.95 & 2.48 & 2.89 & 3.20 \\
\hline \(2p^3P_{1/2} - 2s_1^3/2\) & 0.661 & 0.715 & -- & -- & 0.660 & 0.692 & -- & -- & 0.672 \\
\hline \(2p^3P_{1/2} - 2s_1^3/2\) & 0.624 & 0.680 & -- & -- & 0.748 & 0.764 & -- & -- & 0.735 \\
\hline \(2p^3P_{1/2} - 2s_1^3/2\) & 2.69 & 3.04 & -- & -- & 3.26 & 3.41 & -- & -- & 3.22 \\
\hline
\end{tabular}
\caption{A comparison of present C II \(2s^22p^2 - 2s2p^2\) fine-structure effective collision strengths with the previous evaluations of Blum & Pradhan (1992, BP), Lennon et al. (1985, L) and Hayes & Nussbaumer (1984b, HN) at specific temperatures.}
\end{table}

Figs 4 and 5. For \(2p^3P_{1/2} - 2s_1^3/2\) in Fig. 4, the background collision strength increases gradually over the energy range within which the target states lie (0–1.538 Ryd). Some weak resonances, mostly narrow in shape, can be observed in the spectra. The effective collision strength is given in Fig. 4(b) and is compared with the work of Blum & Pradhan (1992). For this transition, the effective collision strength rises gently with increasing temperature, up to a peak at \(\log T(K) = 4.5\), after which it begins to fall rapidly. Within the temperature range for which a comparison is possible, the effective collision strength is generally lower than that calculated by Blum & Pradhan. The background of the collision strength for the \(3P_{3/2} - 3D_{5/2}\) transition shown in Fig. 5(a), shows no general increase. Resonances are stronger and will therefore have a greater influence upon the size of the effective collision strength. The corresponding effective collision strength is displayed in Fig. 5(b) together with the
determination of Blum & Pradhan (1992). For this transition, the present calculated effective collision strength is higher than that of Blum & Pradhan by an average of 8 per cent over the temperature range for which a comparison is possible.

The numerical values of the effective collision strength for the four fine-structure \(^2\text{P}^o\rightarrow^2\text{D}^o\) transitions are compared in Table 5 with the results of Blum & Pradhan (1992), Lennon et al. (1985) and Hayes & Nussbaumer (1984b). The present evaluations are in good agreement with the previous calculations.

### 3.4 The \(2s^22p^2^3\text{P}^e\rightarrow2s2p^2^1\text{S}^e\) fine-structure transitions

As an example, we present the collision strength for \(2s^22p^2^3\text{P}_{3/2}^o\rightarrow2s2p^2^3\text{S}_{1/2}^e\) in Fig. 6(a). Following a rather broad resonance close to the threshold energy, the background collision strength increases gradually for the incident electron energy displayed. Many resonant structures can be observed in this energy range and have been fully delineated through the use of a fine energy mesh.

The Maxwellian-averaged collision strength corresponding to this transition is presented graphically in Fig. 6(b) together with the values calculated by Blum & Pradhan (1992). The effective collision strength from the present calculation decreases slightly with increasing temperature before log \(T(K)\) begins to rise. At first the increase is gentle but beyond log \(T(K)\), a steep rise in the effective collision strength can be observed. For lower temperatures [below log \(T(K)\) = 4.25] the present results are up to 11 per cent less than those of Blum and Pradhan. However, at higher temperatures, the present evaluations become larger than the previous work.

We compare the present numerical values of the effective collision strength for the two \(^3\text{P}^e\rightarrow^2\text{S}^e\) fine-structure transitions at temperatures of 2000, 10 000 and 20 000 K with the evaluations of Blum & Pradhan (1992) in rows 11 and 12 of Table 5. The present results are lower.

### 3.5 The \(2s^22p^2^3\text{P}^o\rightarrow2s2p^2^3\text{P}^e\) fine-structure transitions

Four fine-structure transitions arise between the \(2s^22p^2^3\text{P}^o\) and \(2s2p^2^3\text{P}^e\) states. Example collision strengths and corresponding effective collision strengths are displayed in Figs 7 and 8. For these excitations, the background collision strength is observed to increase gradually to approximately twice its initial value in the energy range 0–1.54 Ryd. Resonant structures appear in the spectrum, and have been considered fully for these transitions.

The present effective collision strength for \(^3\text{P}^o_{3/2}\rightarrow^2\text{P}^e_{3/2}\) is compared with the work of Blum & Pradhan (1992) in Fig. 7(b). It shows a very shallow increase in effective collision strength with an increase in temperature up to log \(T(K)\) = 4.25, after which the effective collision strength begins to fall rapidly. The results for \(^3\text{P}^o_{1/2}\rightarrow^2\text{P}^e_{1/2}\) are similar. The present results are lower than those obtained by Blum & Pradhan, for which the data show only a gradual rise. There are thus larger differences at temperatures higher than log \(T(K)\) = 4.25.

For \(^3\text{P}^o_{1/2}\rightarrow^2\text{P}^e_{1/2}\) (Fig. 8b), the effective collision strength can be seen to rise gradually with increasing temperature up to log \(T(K)\) = 4.5, after which a shallower decrease in effective collision strength can be observed. Below log \(T(K)\) = 3.75 the present results are seen to be lower than those of Blum & Pradhan (1992); above this temperature they become higher than the previous evaluations.

In the last four rows of Table 5, the present evaluations of the effective collision strengths for the four \(^2\text{P}^o\rightarrow^2\text{P}^e\) transitions are compared numerically with the work of Blum & Pradhan (1992). In general, good agreement is noted.

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**Figure 6.** (a) Collision strength as a function of incident electron energy in Ryd and (b) the effective collision strengths as a function of log temperature in K for the \(2s^22p^2^3\text{P}_{3/2}^o\rightarrow2s2p^2^3\text{S}_{1/2}^e\) fine-structure transition [solid line: present results, dashed line: Blum & Pradhan (1992)].

**Figure 7.** (a) Collision strength as a function of incident electron energy in Ryd and (b) the effective collision strengths as a function of log temperature in K for the \(2s^22p^2^3\text{P}_{3/2}^o\rightarrow2s2p^2^3\text{P}_{1/2}^e\) fine-structure transition [solid line: present results, dashed line: Blum & Pradhan (1992)].
4 CONCLUSIONS

We have presented collision strengths and the corresponding Maxwellian-averaged effective collision strengths for the ground-state fine-structure forbidden transition and for the $2s^22p^2\frac{1}{2}S - 2s2p^2\frac{1}{2}S$ fine-structure transition [solid line: present results, dashed line: Blum & Pradhan (1992)].

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