main line emitter discovered to date. The peak flux densities $J$ are $5.3 \times 10^{-20}$ W m$^{-2}$ and $1.7 \times 10^{-20}$ W m$^{-2}$ at 1665 and 1667 MHz respectively. There is weak emission at 1612 MHz and weak absorption at 1720 MHz.

OH338.9+0.6 is the only other source in Table 1 to have emission at 1612 and/or 1720 MHz. The 1612 MHz emission is slightly stronger than that at 1665 MHz, while the 1667 MHz transition is seen in absorption.

Our search is largely complementary in sky coverage to recent searches from the northern hemisphere. The detailed comparison of results in the small region of overlap is given elsewhere.

The longitude distribution of known sources of emission on the main 18 cm lines shows a concentration to the inner parts of the Galaxy. Half the OH emitters are found between longitudes $333^\circ$ and $26^\circ$. They are found close to the plane, half of them lying between $b = \pm 0.28$. The distributions show that we are dealing with distant objects in the galactic disk. The longitude distribution shows marked concentrations in some ranges, e.g. for $328^\circ < l < 342^\circ$, $345^\circ < l < 356^\circ$, close to $l = 20^\circ$, for $30^\circ < l < 36^\circ$ and $43^\circ < l < 50^\circ$. These are possibly directions tangential to spiral arms.

For 13 of the more intense continuum sources we found strong absorption at 1665 MHz. Nearly 90 sources have since been observed at 1667 MHz, and 9 more cases of OH absorption have been found. These results will be reported elsewhere.

For any atom perturbed by neutral particles at low pressures, the frequency distribution of emitted line radiation has a Lorentz profile, and measurements can thus only determine two parameters for the line—width and shift. Only two parameters for the interaction potential can then be inferred. Hindmarsh, Petford, and Smith used experimental widths and shifts to evaluate the constants $C_{12}$ and $C_6$ for a Lennard-Jones potential. Despite the inclusion of the repulsive part, the van der Waals constant was still, in many cases, a factor of four larger than the upper limit set by the Unsold value. Hindmarsh, du Plessis, and Farr included a further term $-C_8 r^{-8}$ and found the ‘experimental’ values of $C_6$ showed improved agreement with theory.

From quantum mechanical perturbation theory, the van der Waals constant can be written

$$C_6 = -\frac{2e^2}{3h} \sum_{j, j'} \frac{\left| \langle i | j \rangle \right|^2 \left| \langle i' | j' \rangle \right|^2}{E_i - E_j + E_{ii} - E_{jj}},$$

where $E_i$, $E_j$, and $E_{ii}$ are the initial energies and $E_{ii}$, $E_{jj}$ are the intermediate energies for the radiating atom and perturber respectively, and $\langle i | j \rangle$ and $\langle i' | j' \rangle$ are the corresponding dipole matrix elements. If the perturbing particles are hydrogen or the inert gases, these atoms are usually in their ground state. Since the excited states of these atoms are more than 10 eV above the ground state, the term $E_i - E_j$ for the radiator is much smaller, and may be neglected. The van der Waals constant was then expressed by Unsold as

$$C_6 = \frac{\alpha e^2}{\hbar} \langle r^2 \rangle,$$

where $\alpha$ is the static polarizability of the perturber and $\langle r^2 \rangle$ denotes the mean square radius of the radiator in the particular state. The Unsold expression for the van der Waals constant must thus be an upper limit to the value of $C_6$.

For a van der Waals type interaction, the theory of line broadening predicts for all spectral lines a red shift and a constant shift/width ratio. However, experiments have observed blue shifts and variable shift/width ratios. The existence of blue shifts suggests the presence of a repulsive potential. An inverse power form, $C_n r^{-n}$, is usually chosen to represent the repulsive part, and it has been found that $n = 12$ provides the best fit for experimental results; such a potential is known as the Lennard-Jones potential. From quantum mechanical considerations of the repulsive forces occurring when charge distributions overlap, a potential of the form $e^{-\lambda r}$ has been found to have closer agreement with the calculations than in inverse power form.

For any atom perturbed by neutral particles at low pressures with closer agreement with the calculations than in inverse power form. For collisions between neutral atoms, the broadening can be treated in the impact approximation and only adiabatic effects of the collisions considered. The resulting width and shift of the line are

$$\frac{\hbar}{e} = h/2 \pi.$$
\[ w = 2\pi N \int_0^\infty \int_0^\infty \rho [1 - \cos \phi(\rho, v)] \, d\rho \, dv \]
\[ d = 2\pi N \int_0^\infty \int_0^\infty \rho [\sin \phi(\rho, v)] \, d\rho \, dv , \]
where \( \phi(\rho, v) \) is the total phase shift produced in the atom by a collision with impact parameter \( \rho \) and velocity \( v \); \( N \) is the perturber particle density. The phase shift is given by
\[ \phi(\rho, v) = \frac{1}{\hbar} \int_0^\infty V(r) \, dt = \frac{1}{\hbar} \int_0^\infty V(\rho, v, t) \, dt . \]

In this paper, the interaction potential has been chosen as a combination of an exponential repulsive part and an \( r^{-6} \) attractive potential:
\[ V(r) = A e^{-\lambda r} - C_6 r^{-6} . \]

Since the Unsöld expression, equation (2), should give a reasonable estimate of \( C_6 \), this value will be used in equation (5), thus allowing \( A \) and \( \lambda \) to be determined. A disadvantage of this form of the potential is that it does not represent the true potential at very small interatomic distances; it predicts an attractive force rather than repulsive. However, collisions with such small separations make negligible contribution to the width and shift.

For an exp-6 potential, equation (5), the width and shift can be expressed
\[ w = 2k B(\alpha, \beta) , \quad d = k S(\alpha, \beta) , \]
where
\[ k = \frac{N \pi \bar{v} a_0^2}{8\hbar \tilde{v}} \left[ \frac{3\pi C_6 e^{2\lambda}}{8\hbar} \right]^{2/5} \]
\[ \bar{v} = \text{mean relative velocity of collision} \]
and
\[ B(\alpha, \beta) = \beta^{-2/5} \int_0^\infty x \sin^2 \left\{ \frac{\phi(x)}{2} \right\} \, dx , \]
\[ S(\alpha, \beta) = \beta^{-2/5} \int_0^\infty x \sin \{\phi(x)\} \, dx . \]

The total phase shift is
\[ \phi(x) = \alpha x K_1(x) - \beta x^{-5} \]
with \( \alpha = \frac{2Ae^2}{\hbar \bar{v} \tilde{\lambda}} \), \( \beta = \frac{3\pi C_6 e^{2\lambda}}{8\hbar \tilde{v}} \), and \( x = \frac{\lambda \rho}{a_0} \);
\( K_1(x) \) is the modified Bessel function of the second kind, first order.

From experimental values of the broadening of krypton lines by the inert gases, equations (7a) and (7b) have been inverted to determine the parameters of the potential. The results obtained for krypton-helium and krypton-neon interactions are shown in Table 1. From the form of the integrands of equations (7a) and (7b), shown in Figures 1 and 2, it can be seen that there is a region of maximum contribution to the width and shift. It is in this region that the potential will be best determined.

Results obtained by an iterative technique have converged rapidly for blue-shifted lines, but are inconclusive for red-shifted lines. This is not surprising since it is only the parameters for the repulsive component of the potential

![Figure 1. Width and shift integrands for the λ7601 line of krypton perturbed by neon. Exp-6 and Lennard-Jones potentials also shown.](https://www.cambridge.org/core/terms).

![Figure 2. Temperature dependence of width and shift integrands for the λ7601 line of krypton perturbed by helium.](https://www.cambridge.org/core/terms).

**Table I**

<table>
<thead>
<tr>
<th>Interaction</th>
<th>( A ) ((e^2 , a_0^{-1}))</th>
<th>( \lambda ) ((a_0^{-1}))</th>
<th>( C_6 ) ((e^2 , a_0^{-6}))</th>
</tr>
</thead>
<tbody>
<tr>
<td>Kr-He</td>
<td>1.336</td>
<td>0.74</td>
<td>48.43</td>
</tr>
<tr>
<td>Kr-Ne</td>
<td>1447.2</td>
<td>1.312</td>
<td>89.86</td>
</tr>
</tbody>
</table>
which are determined. Since $C_4$ has been fixed, there can be no increase in the depth of the potential well by variation of $A$ or $\lambda$. Figure 1 shows the width and shift integrands for the krypton-neon interaction and the corresponding $\exp(-1)$ potential. The Lennard-Jones potential which yields the same width and shift is also shown. For krypton-neon, the sum of the radii of the atoms is 9.2 bohrs, so that even at a separation of 14 bohrs there will be considerable interaction from higher order terms in the multipole expansion.

While Hindmarsh, du Plessis, and Farr found improved agreement for $C_6$ by inclusion of an attractive $r^{-8}$ term, the constants which they determined bear no real relation to the theoretical values. They must be regarded as parameters which are evaluated by requiring that the potential be of the form $C_6 r^{-12} - C_8 r^{-8} - C_8 r^{-6}$.

A satisfactory potential must also be able to predict the temperature dependence of the width and shift. From experimental values at 80°K, the potential for the krypton-helium interaction was used to estimate the values at 295°K; agreement with experiment was within 10%. The integrands at both temperatures are shown in Figure 2. It can be seen that the increase in temperature causes the region of maximum contribution to occur at smaller impact parameters. This suggests that a more complete representation of the potential could be obtained if a series of measurements at different temperatures are performed for each radiator-perturber interaction.

Scattering experiments between neutral atoms also provide a method of determining empirical interatomic potentials. However, in such experiments, both atoms are in the ground state, or at least in a metastable state. The inversion of spectral line widths and shifts provides the only method of estimating empirical potentials for excited atomic states. Such empirical potentials can then be compared with theoretical potentials such as those calculated by Lewis, McNamara and Michels.


Broadening of the Solar Sodium D-Lines by Atomic Hydrogen

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It is well known 1-3 that the wings of the solar sodium D-lines cannot be predicted successfully using present line-broadening theories and elemental abundances derived from an equivalent-width analysis of the weak lines. By far the most important cause of broadening of the D-lines is collision broadening by neutral hydrogen. Because of the complexity of the interaction between an emitting sodium atom and a colliding hydrogen atom, it has been customary to express the interatomic potential in the form of an inverse power series in the interatomic distance $R$. Such an approach is convenient when the cross section for broadening collisions is much greater than the extent of the excited state wave function, since the leading term in the interaction series is sufficient. A simple expression for the line broadening can thus be derived, on the basis of a van der Waal type interaction.

The importance of repulsive contributions even at laboratory temperatures has been emphasized, however, by several authors.5-8 Also, the usual approximation of a single set of parameters to describe the interatomic potential for all orientations of the colliding atoms is clearly invalid for all but S-states, since the real situation involves many angle-dependent terms and the variation of the coupling of the spin and orbital angular momenta with interatomic separation.

The difficulties of separately computing many terms in the series expansion for the potential, in addition to the overlap contributions, may be overcome by using molecular potentials and treating the colliding atoms as a molecule during the collision. Molecular energies for the system Na-H have therefore been calculated by the valence configuration-interaction (VCI) method9,10 and are illustrated in Figure 1.

The calculated potential energy curves for the excited states have been uniformly adjusted by the small calculated error (.031 eV) in the $1^2P-2^3S$ term value in order to ensure proper atomic limiting behaviour. An average term value is employed since the $II$ molecular states are degenerate in our approximate theory. A parallel Hartree-Fock analysis of the lowest $3\Sigma^+$ and $3\Pi$ states indicated that the region of validity of the potential curves derived from the VCI wave functions extends down to about 3A, which is adequate for our purposes.

The variation with temperature of the calculated (half)