Ab initio computed diabatic potential energy surfaces of OH–HCl

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The two four-dimensional diabatic potential energy surfaces (DPESs) for OH–HCl are computed that correlate with the twofold degenerate \(^2\Sigma^+\) ground state of the free OH radical. About 20 000 points on the surface are obtained by the \textit{ab initio} coupled-cluster and multi-reference configuration interaction methods. Analytic forms for the diabatic potential energy surfaces are derived as expansions in complete sets of orthogonal functions depending on the three intermolecular angles. The numeric computation of the angular expansion coefficients is discussed. The distance-dependence of the angular coefficients is represented by the reproducing kernel Hilbert space method. It is checked that both diabatic potentials converge for large intermolecular separations to the values computed directly from the electrostatic multipole expansion. The final DPESs are discussed and illustrated by some physically meaningful one- and two-dimensional cuts through them. © 2005 American Institute of Physics. [DOI: 10.1063/1.1949198]

I. INTRODUCTION

The abundant OH(\(^2\Sigma^+)\) radical is very reactive and known to take part in many atmospheric and combustion processes. Reactions of OH with the hydrogen halides efficiently convert halogens to their active forms. In particular, the reaction \(\text{OH} + \text{HCl} \rightarrow \text{H}_2\text{O} + \text{Cl}\) is an important source of active chlorine in the atmosphere. Several experimental and computational studies have been devoted to this reaction, see Ref. 1 for a survey of the literature. An important aspect of the theoretical studies \(^2\)–\(^6\) is the potential energy surface for the reaction which was investigated either through the use of semi-empirical models or by means of \textit{ab initio} electronic structure calculations. Most of these studies concentrate on the transition state for the reaction. In the article of Yu and Nyman \(^2\) it is pointed out that a stable OH–HCl van der Waals complex occurs in the entrance channel of the reaction. Since it is known from theoretical work \(^7\) on the analogous reaction \(\text{O} + \text{HCl} \rightarrow \text{OH} + \text{Cl}\) that the existence of such a complex may have a large influence on the energy dependent reaction probabilities, detailed spectroscopic studies on the OH–HCl complex will be undertaken. \(^8\) Also rotationally inelastic state-to-state scattering cross sections of \(\text{OH}(\^2\Pi_{\text{u}})\) in collision with HCl were measured recently. \(^1\) To support this experimental work, computations on the same cross sections are underway in our group. In addition, computations on the bound states of the van der Waals complex \(\text{OH}(\^2\Sigma^+)\)–HCl are planned. The comparison of experimental and theoretical cross sections gives a thorough check on computed results and contributes importantly to the understanding of the physical origin of propensity rules observed in Ref. 1. The scattering and bound state calculations are based on the expanded potential energy surfaces reported in the present article.

A problem that occurs for many open-shell van der Waals molecules, including OH(\(^2\Sigma^+)\)–HCl, is the breakdown of the Born–Oppenheimer (BO) approximation. The BO approximation is applicable when the eigenvalues of the electronic Hamiltonian are everywhere (as function of the intermolecular nuclear coordinates) well separated in energy. However, for linear geometries of OH–HCl as well as for large intermolecular distances, the adiabatic states that correlate with the \(^2\Sigma^+\) state of OH become degenerate, i.e., two potential energy surfaces (PESs) come close or even coincide. Terms arising from the nuclear kinetic energy operator, neglected in the BO approximation, are then no longer negligible. A solution for this problem was pointed out by Smith, \(^9\) who introduced the concept of diabatic states. Diabatic states are linear combinations of adiabatic states—which by definition are eigenstates of the electronic Hamiltonian—that minimize (or annihilate) the non-BO nuclear kinetic coupling terms concurring with approaching PESs. The introduction of diabatic states is at the expense of giving a non-diagonal PES matrix and, correspondingly, it leads to multiple nuclear motion equations that are coupled by potential energy terms. Since we have here two orthonormal adiabatic states from which two orthonormal diabatic states are constructed, the transformation between them is completely specified by one rotation angle: “the diabatic angle” \(\gamma\). As is explained in more detail below, we compute this angle by inspection of molecular orbitals and the lowest two multi-reference configuration interaction (MRCI) states.

It is known from much work on closed-shell van der Waals molecules that the coupled-cluster-singles-doubles plus non-iterative triples [CCSD(T)] approach is at present the most reliable supermolecule method for the computation of PESs. For open-shell complexes the partially spin-restricted CCSD(T) [RCCSD(T)] method is therefore an obvious choice. However, for most points on the PES a dimer
consisting of two diatoms has no point group symmetry and accordingly RCCSD(T) is strictly speaking only suitable to give the ground state energy. This is regrettable as the RCCSD(T) method is size-extensive, while most CI methods are not. As we just pointed out, the diabatic angle \( \gamma \) will be extracted from MRCI wave functions, so that MRCI calculations are required anyway. Therefore, to alleviate the size-extensivity problem, we computed ground state energies by the RCCSD method and used the MRCI method for the first excitation energy \( \Delta E_{\text{MRCI}} \).

The outline of the paper is as follows: First we introduce the concept of diabatic potential energy surfaces (DPESs). Then we introduce an analytic expansion of these surfaces, which is suggested by long range theory and is appropriate for coupled-channel scattering and bound-state calculations. This is followed by a discussion of the practical details of the calculations. Finally we give the results and present some illustrative cuts through the diabatic and adiabatic potentials.

II. THEORY
A. Diabatic potentials

Under influence of the interaction the rotational symmetry of \( \mathrm{OH} \) is broken and the degeneracy of the \( \Pi \) states is lifted, because they are contaminated by states of symmetry other than \( \Pi \). This is a second-order effect. In first-order in the interaction, \( \Pi_1 \) and \( \Pi_2 \) are mixed in a definite manner. To a good approximation we can therefore write the two broken-symmetry adiabatic states \( \Psi_1 \) and \( \Psi_2 \), with interaction energies \( V_1 \) and \( V_2 \), as a linear combination of pure \( \Pi_1 \) and \( \Pi_2 \) states,

\[
(\Pi_1, \Pi_2) = (|\Psi_1\rangle, |\Psi_2\rangle) \mathcal{R}(\gamma)
\]

with \( \mathcal{R}(\gamma) = \begin{pmatrix} \cos \gamma & -\sin \gamma \\ \sin \gamma & \cos \gamma \end{pmatrix} \).

Since the diabatic states \( \Pi_1 \) and \( \Pi_2 \) are monomer states, they give vanishing distance-dependent (radial) nuclear kinetic coupling terms. If the expression (1) for the adiabatic states were exact, the coupling terms would be transformed away by this equation. In any case, we assume the radial coupling terms to become so small by this transformation that they can be neglected in the nuclear motion problem. Transformation of the (diagonal) adiabatic energy matrix gives

\[
\mathcal{R}(\gamma)^T \begin{pmatrix} V_1 & 0 \\ 0 & V_2 \end{pmatrix} \mathcal{R}(\gamma) = \frac{V_1 + V_2}{2} \begin{pmatrix} 1 & 0 \\ 0 & 1 \end{pmatrix} + \frac{V_2 - V_1}{2} \begin{pmatrix} -\cos 2\gamma & \sin 2\gamma \\ \sin 2\gamma & \cos 2\gamma \end{pmatrix}.
\]

(2)

The diabatic angle \( \gamma \), which is a function of the intermolecular coordinates, can be determined from the \textit{ab initio} states \( \Psi_1 \) and \( \Psi_2 \) via Eq. (1). Use of the eigenstates \( |\Lambda\rangle \) with \( \Lambda = \pm 1 \) of \( \hat{L}_z \), the projection of the electronic angular momentum on the O–H axis, gives the required reduction of the azimuthal nuclear kinetic coupling terms.\(^{10}\) The states \( |\Lambda\rangle \) are simply related to the diabatic states \( |\Pi_1\rangle \) and \( |\Pi_2\rangle \)

\[
(|1\rangle, |-1\rangle) = \frac{1}{\sqrt{2}} (|\Pi_1\rangle, |\Pi_2\rangle)(-1 \quad 1 \\ -i \quad -i) \quad \text{(3)}
\]

and further transformation of the energy matrix gives

\[
\begin{pmatrix} V_{-1,-1} & V_{-1,1} \\ V_{1,-1} & V_{1,1} \end{pmatrix} = \frac{1}{2} \begin{pmatrix} V_1 + V_2 & (V_2 - V_1) \exp(-2i\gamma) \\ (V_2 - V_1) \exp(2i\gamma) & V_1 + V_2 \end{pmatrix}.
\]

(4)

We will refer to the matrix elements \( V_{\pm 1, \pm 1} \) as diabatic potential energy surfaces (DPES).

For planar geometries the adiabatic states \( |\Psi_1\rangle \) and \( |\Psi_2\rangle \) have \( A'/A'' \) (even/odd) symmetry with respect to reflection in the plane of the complex. Except when the OH radical lies along the intermolecular bond axis, they must coincide with the diabatic states \( |\Pi_1\rangle \) and \( |\Pi_2\rangle \) and the diabatic angle \( \gamma \) equals 0° or 90°. When the OH radical lies along the bond axis the angle \( \gamma \) is determined by the orientation of the plane through the HCI axis and the bond axis. This will be discussed in more detail below.

The use of the diabatic states \( |\Lambda\rangle = |\pm 1\rangle \) is very convenient also for the inclusion of spin-orbit coupling in bound state and scattering calculations on the OH–HCl complex. Spin-orbit coupling in the \( X^2\Pi \) ground state of the OH radical can be represented by the operator \( \hat{A}_L \hat{S}_z \), with a coupling constant \( A = 139.21 \text{ cm}^{-1} \). The operator \( \hat{L}_z \) is diagonal in the basis \( |\Lambda\rangle \) with eigenvalues \( \Lambda = \pm 1 \) and \( \hat{S}_z \) is diagonal in the spin basis \( |\pm \frac{1}{2}\rangle \), so that only a diagonal term \( \pm A/2 \) has to be added to the diabatic potential matrix in Eq. (4) for the spin-orbit coupled eigenstates \( |\pm \frac{1}{2}\rangle \). The \( \hat{S}_z \) of \( \hat{c}_z \) is diagonal in the \( xz \)-plane.

B. Expansion of DPESs for two interacting open-shell diatoms

We now present an expansion of DPESs for two open-shell diatoms. Although in this paper we are studying the case of one open- and one closed-shell diatomic molecule, the generalization to two open-shell diatoms is easy and natural. This expansion will be derived from the first-order (in the interaction \( V \)) long-range (no exchange) perturbation matrix. In the long range diabatic states are equal to products of monomer (zeroth-order) states. The general expansion, derived in this section for the matrix elements of the first-order perturbation matrix, is in principle exact, as it is an expansion in a complete basis. There are, however, an infinite number of complete expansion bases. The expansion basis chosen here has the important advantage that the matrix elements go continuously over into the correct large \( R \) expressions and the correct expressions for linear dimer geometries.

Let us first consider one rigid diatom with its molecular axis parallel to the \( z \)-axis of an arbitrarily oriented space-fixed (SF) frame. Its electronic wave functions \( |\Lambda\rangle \) are adapted to \( C_{\text{sym}} \),

\[
\hat{R}_z(\phi)|\Lambda\rangle = e^{-i\phi \hat{c}_z}|\Lambda\rangle = e^{-i\phi |\Lambda\rangle}.
\]

(5)

Reflection \( \hat{c}_{xz} \) in the \( xz \)-plane is represented in the basis \( |\pm |\Lambda\rangle \) by
where the phase is a matter of convention. Rotate now the vector \( \mathbf{r} \), pointing from one nucleus of the diatom to the other, to its desired orientation given by polar coordinates \( \theta \) and \( \phi \) with respect to the SF frame. This requires the active rotation \( \hat{R}(\theta, \phi) = \hat{R}_z(\theta) \hat{R}_z(\phi) \) of \( \mathbf{r} \), which initially was parallel to the SF z-axis. In order to keep electrons and nuclei in the same relative position we must rotate the electronic wave function in the very same way,

\[
\hat{R}_z(\theta) \hat{R}_z(\phi) |\Lambda\rangle = |\Lambda(\theta, \phi)\rangle.
\]  

Note that rotation over \( \psi \) around the rotated diatomic axis \( z' \) still gives the factor \( \exp(-i\psi\Lambda) \),

\[
e^{-i\psi L_{z'}} |\Lambda(\theta, \phi)\rangle = e^{-i\psi\Lambda} |\Lambda(\theta, \phi)\rangle.
\]

where \( \hat{L}_{z'} = \hat{R}(\theta, \phi) \hat{L}_{z'} \hat{R}(\theta, \phi) \),

because

\[
\hat{L}_{z'} |\Lambda(\theta, \phi)\rangle = \hat{R}(\theta, \phi) \hat{L}_{z'} |\Lambda(0, 0)\rangle = \Lambda |\Lambda(\theta, \phi)\rangle,
\]

with \( |\Lambda(0, 0)\rangle = |\Lambda\rangle \). Hence, the most general electronic wave function obtained by a three-angle Euler rotation \( \hat{R}(\psi, \theta, \phi) = \hat{R}_{z'}(\psi) \hat{R}_z(\theta) \hat{R}(\phi) \) is

\[
|\Lambda(\psi, \theta, \phi)\rangle = \hat{R}(\psi, \theta, \phi) |\Lambda\rangle = e^{-i\phi\Lambda} |\Lambda(\theta, \phi)\rangle.
\]  

We will keep using also the two-angle rotation of the electronic wave functions, which should be understood as \( |\Lambda(\theta, \phi)\rangle = |\Lambda(0, 0)\rangle \).

Let us next consider two diatoms A and B with Euler angles \( (\psi_A, \theta_A, \phi_A) \) and \( (\psi_B, \theta_B, \phi_B) \) with respect to the same SF frame. In first-order perturbation theory one constructs a matrix of the interaction operator \( \hat{V} = \hat{H} - \hat{H}_A - \hat{H}_B \) with zeroth-order bra and ket that are products of the degenerate monomer states. Rotating each of the diatoms and assuming that the vector \( \mathbf{R} \) pointing from A to B has the SF polar angles \( (\Phi, \Theta) \), we find that the matrix element

\[
\tilde{V}_{\Lambda_A', \Lambda_B'}^{\Lambda_A, \Lambda_B} = \sum_{L_A} \sum_{L_B} \sum_{M_A, M_B} \sum_{M} C_{L_A M A}^{\Lambda_A} C_{L_B M B}^{\Lambda_B} \langle \Lambda_A' | L_A M A | \Lambda_A \rangle \langle \Lambda_B' | L_B M B | \Lambda_B \rangle |\Lambda(\theta_A, \phi_A)\rangle \langle \Lambda(\theta_B, \phi_B)\rangle \times \langle \Lambda(\theta_A, \phi_A) | L_A M A | \Lambda_A \rangle \langle \Lambda(\theta_B, \phi_B) | L_B M B | \Lambda_B \rangle
\]

\[
e^{-i\phi_A(\Lambda_A' - \Lambda_A)} e^{-i\phi_B(\Lambda_B' - \Lambda_B)} \langle \Lambda(\theta_A, \phi_A) | L_A M A | \Lambda_A \rangle \langle \Lambda(\theta_B, \phi_B) | L_B M B | \Lambda_B \rangle
\]

\[
e^{-i\phi_A(\Lambda_A' - \Lambda_A)} e^{-i\phi_B(\Lambda_B' - \Lambda_B)} |\Lambda(\theta_A, \phi_A)\rangle \langle \Lambda(\theta_B, \phi_B)\rangle
\]

(11)

is a function of two sets of Euler angles, one for diatom A and one for diatom B, and of the polar angles \( \Phi \) and \( \Theta \). By the Peter–Weyl theorem\(^{11}\) the matrix element can be expanded in products of rigid rotor functions (complex conjugates of Wigner D-matrices) and spherical harmonic functions \( C_M^L(\Phi, \Theta) \),

\[
\tilde{V}_{\Lambda_A', \Lambda_B'}^{\Lambda_A, \Lambda_B} = \sum_{L_A} \sum_{L_B} \sum_{M_A, M_B} \sum_{M} C_{L_A M A}^{\Lambda_A} C_{L_B M B}^{\Lambda_B} \langle \Lambda_A' | L_A M A | \Lambda_A \rangle \langle \Lambda_B' | L_B M B | \Lambda_B \rangle |\Lambda(\theta_A, \phi_A)\rangle \langle \Lambda(\theta_B, \phi_B)\rangle \times D_{M A}^{L A}(\phi_A, \theta_A, \psi_A)^* \times D_{M B}^{L B}(\phi_B, \theta_B, \psi_B)^* C_M^L(\Phi, \Theta)^*.
\]  

Equating factors containing \( \psi_A \) and \( \psi_B \), while realizing that the complex conjugate D-matrices contain \( \exp(iK_A \psi_A) \) and \( \exp(iK_B \psi_B) \), shows that \( K_A = \Lambda_A' - \Lambda_A \) and \( K_B = \Lambda_B' - \Lambda_B \), so that the summations over \( K_A \) and \( K_B \) each shrink to one term.

Any matrix element of an energy operator is invariant under any rotation \( \hat{R} \) of all its integration variables (the electron coordinates), because a rotation is unitary, so that \( \langle \Phi|V|\Psi\rangle = \langle \Phi|\hat{R}^\dagger \hat{R} \hat{V} |\Psi\rangle = \langle \Phi|\hat{R}^\dagger \hat{R} |\Psi\rangle \). The fact that the matrix element is a rotational invariant implies that it is useful to Clebsch–Gordan couple the right-hand-side of Eq. (12). The elements \( (D^L_{M A} \) of column \( M \) transform as spherical harmonic functions \( C_M^L \) and accordingly it can be shown by an argument, similar to the one used in the proof of the Wigner–Eckart theorem, that the expansion in terms of invariants becomes

\[

V_{\Lambda_A', \Lambda_B'}^{\Lambda_A, \Lambda_B} = \sum_{L_A} \sum_{L_B} \sum_{M_A, M_B} \sum_{M} C_{L_A M A}^{\Lambda_A} C_{L_B M B}^{\Lambda_B} \langle \Lambda_A' | L_A M A | \Lambda_A \rangle \langle \Lambda_B' | L_B M B | \Lambda_B \rangle \times \langle L, M | L_A, M_A | L_B, M_B \rangle \times D_{M A}^{L A}(\phi_A, \theta_A, 0)^* \times D_{M B}^{L B}(\phi_B, \theta_B, 0)^* C_M^L(\Phi, \Theta)^*.
\]

(13)

Note that the angles \( \psi_A \) and \( \psi_B \) involved only in the rotation of the electron coordinates of the diatoms around the molecular axes have dropped out from this expression. The potential matrix elements \( V_{\Lambda_A', \Lambda_B'}^{\Lambda_A, \Lambda_B} \) defined in Eq. (11) are integrals over electronic wave functions \( |\Lambda(\theta_A, \phi_A)\rangle \) and \( |\Lambda(\theta_B, \phi_B)\rangle \) and the electron coordinates are integration variables.

C. Properties of expansion coefficients

From certain invariances follow a few important properties of the expansion coefficients \( V_{L_A L_B}^{\Lambda_A \Lambda_B M M} \) appearing in Eq. (13). In the first place the operator \( \hat{V} \) is invariant under inversion of all particles (electrons and nuclei) with respect to an arbitrary point in space. If we choose the origin of the SF frame as inversion point, this operation can be expressed as a reflection \( \hat{\sigma}_{z'} \) in the xy-plane of the SF frame times an overall rotation around the SF y-axis. Since we already saw that \( \hat{V} \) is invariant under any overall rotation, we consider only the effect of \( \hat{\sigma}_{z'} \). For the azimuthal angle of the three nuclear vectors \( \mathbf{r}_A, \mathbf{r}_B, \) and \( \mathbf{R} \) this reflection gives the map \( \phi \rightarrow -\phi \), hence for diatom X \((X=A, B)\) holds

\[
\hat{\sigma}_{z'} |\Lambda(\theta_X, \phi_X)\rangle = \hat{R}(\theta_X, -\phi_X) \hat{\sigma}_{z'} |\Lambda(0, 0)\rangle,
\]

(14)

so that Eq. (6) gives

\[
(-1)^\Lambda \begin{pmatrix} 0 & 1 \\ 1 & 0 \end{pmatrix},
\]

(6)
\[ \hat{\sigma}_{\text{sc}}(\lambda_A(\theta_A, \phi_A), \lambda_B(\theta_B, \phi_B)) = (-1)^{L_A + L_B - |\lambda_A|} \lambda_A(\theta_A, -\phi_A) - \lambda_B(\theta_B, -\phi_B). \] (15)

Hence,
\[ V_{\lambda_A^\prime \lambda_B^\prime}^{\lambda_A \lambda_B} = (-1)^{L_A + L_B - |\lambda_A|} \lambda_A(\theta_A, -\phi_A) - \lambda_B(\theta_B, -\phi_B) \]
\[ \times (-\lambda_A(\theta_A, -\phi_A) - \lambda_B(\theta_B, -\phi_B)) \]
\[ \times |L| \lambda_A(\theta_A, -\phi_A) - \lambda_B(\theta_B, -\phi_B). \] (16)

Expand the right-hand side according to Eq. (13) and substitute in this expansion
\[ D_{\mu, \lambda}(-\phi, \theta, 0)^0 = (-1)^{L_A + L_B + \mu} D_{\mu, -\lambda}(-\phi, \theta, 0)^0, \]
(17)

\[ C_{\mu, \lambda}(\Theta, -\Phi)^0 = (-1)^{M} C_{\mu, -\lambda}(\Theta, \Phi)^0, \]
(18)

and (for integer \( L, M \))
\[ \langle L, -M | L, -M \rangle M \lambda_A + \lambda_B |- M \rangle M \lambda_A + \lambda_B = (-1)^{L_A + L_B + |\lambda_A|} \lambda_A \lambda_B \]
(19)

followed by equating the terms in this expansion with those in the expansion of the noninverted matrix element. In this manner we get the symmetry property
\[ V_{L_A L_B}^{L_A^\prime L_B^\prime} = (-1)^{L_A + L_B + L_A^\prime + L_B^\prime + |\lambda_A| - |\lambda_B|} \lambda_A \lambda_B \]
(20)

In a very similar manner follows from \( V' = V \)
\[ V_{L_A L_B}^{L_A^\prime L_B^\prime} = (-1)^{L_A + L_B + L_A^\prime + L_B^\prime} \left( V_{L_A L_B}^{L_A^\prime L_B^\prime} \right)^*, \]
(21)

where we used that the symmetric property
\[ \langle L, -M | L, -M \rangle M \lambda_A + \lambda_B |- M \rangle M \lambda_A + \lambda_B = (-1)^{L_A + L_B + |\lambda_A| - |\lambda_B|} \lambda_A \lambda_B \]
(22)

Combining the last three equations gives
\[ V_{L_A L_B}^{L_A^\prime L_B^\prime} = \left( V_{L_A L_B}^{L_A^\prime L_B^\prime} \right)^* = \left( V_{L_A L_B}^{L_A^\prime L_B^\prime} \right)^*. \]
(23)

In particular, we see that all expansion coefficients are real.

The expansion coefficients of the diagonal potentials with \( \lambda_A^\prime = \lambda_A \) and \( \lambda_B^\prime = \lambda_B \) obey the relation
\[ V_{L_A L_B}^{L_A^\prime L_B^\prime}(R) = (-1)^{L_A + L_B + L_A^\prime + L_B^\prime} V_{L_A L_B}^{L_A^\prime L_B^\prime}(R) \]
(24)

and, hence, they must vanish for odd values of \( L_A + L_B + L \).

### D. Open-plus closed-shell diatom

We specialize the earlier results to the case that \( A \) is open-shell and \( B \) is closed-shell with the centers of mass of both diatoms lying on the z-axis of the SF frame. If we first rotate the total dimer such that \( R \) is along the z-axis of the SF system, \( \Theta = \Phi = 0 \), we have for the Racah normalized spherical harmonic
\[ C_{\mu, \lambda}^2(\Theta, \Phi) = \delta_{\mu, 0}, \]
(25)

so that the expansion (13) for two open-shell diatoms becomes

\[ V_{\lambda_A^\prime \lambda_B^\prime}^{\lambda_A \lambda_B}(R, \theta_A, \phi_A, \theta_B, \phi_B) = \sum_{L_A L_B} \sum_{L_A^\prime L_B^\prime} \sum_{M} \langle L_A, M | L_A^\prime, M^\prime \rangle \langle L_B, M | L_B^\prime, M^\prime \rangle \times D_{\mu, \lambda}^{L_A, M}(-\lambda_A, \phi_A, 0)^0 D_{\mu, -\lambda}^{L_B, M}(-\lambda_B, \phi_B, 0)^0. \] (26)

Take now diatom B (HCl) in a \( \Sigma \) state: \( \lambda_B^\prime = \lambda_B = 0 \). Use
\[ D_{\mu, \lambda}^{L_B, M}(-\lambda_B, \phi_B, 0)^0 = C_{\mu, \lambda}^2(\theta_B, \phi_B) \]
and we get, suppressing \( \lambda_B \) and \( \lambda_B^\prime \) in the notation and dropping the subscript \( A \) on the \( \lambda_A \),
\[ V_{\lambda_A^\prime \lambda_B}^{\lambda_A \lambda_B} = \sum_{L_A L_B} \sum_{L_A^\prime L_B^\prime} \sum_{M} \langle L_A, M | L_A^\prime, M^\prime \rangle \langle L_B, M | L_B^\prime, M^\prime \rangle \times D_{\mu, \lambda}^{L_A, M}(-\lambda_A, \phi_A, 0)^0 C_{\mu, -\lambda}^2(\theta_B, \phi_B, 0). \] (27)

We see that the diagonal elements \( V_{1,1} = V_{1,1} \) of Eq. (4), which are real, are expanded in terms of \( D_{\mu, \lambda}^{L_B, M}(-\lambda_B, \phi_B, 0)^0 = C_{\mu, \lambda}^2(\theta_B, \phi_B) \) (times spherical harmonics for \( B \)), while the expansion of the complex-valued off-diagonal elements \( V_{1,-1} \) contains the functions \( D_{\mu, 0, \lambda}^{L_B, M}(-\lambda_B, \phi_B, 0)^0 \).

### E. Alternative expansion

In some cases, for example, if the potential is available on a grid of points \( (\theta_A, \theta_B, \phi) \) and the number of \( \phi \) points is smaller than the number of \( \theta \) and \( \theta_B \) points, it is convenient to expand the potential in uncoupled functions. We restrict our attention to the case of an open-shell diatom \( A \) in interaction with a closed-shell diatom \( B \) and rewrite Eq. (27),
\[ V_{\lambda_A^\prime \lambda_B}^{\lambda_A \lambda_B} = \sum_{L_A L_B} \sum_{L_A^\prime L_B^\prime} \sum_{M} \langle L_A, M | L_A^\prime, M^\prime \rangle \langle L_B, M | L_B^\prime, M^\prime \rangle \times D_{\mu, \lambda}^{L_A, M}(-\lambda_A, \phi_A, 0)^0 C_{\mu, -\lambda}^2(\theta_B, \phi_B, 0). \] (28)

with
\[ V_{L_A L_B}^{L_A^\prime L_B^\prime}(R) = \sum_{L} V_{L_A L_B}^{L_A^\prime L_B^\prime}(R) \langle L, 0 | L_A, M | L_B, -M \rangle \] (29)

Given the coefficients \( V_{L_A L_B}^{L_A^\prime L_B^\prime}(R) \) that we refer to as “LLL coefficients”, the “LLL coefficients” \( V_{L_A L_B}^{L_A^\prime L_B^\prime}(R) \) follow from the unitarity of the Clebsch–Gordan coefficients,
\[ V_{L_A L_B}^{L_A^\prime L_B^\prime}(R) = \sum_{M} \langle L, 0 | L_A, M | L_B, -M \rangle V_{L_A L_B}^{L_A^\prime L_B^\prime} \] (30)

From Eqs. (20)–(22) it follows that the LLM expansion coefficients must obey
\[ V_{L_A L_B}^{L_A^\prime L_B^\prime}(R) = V_{L_A L_B}^{L_A^\prime, -L}(R) \] (31)

\[ V_{L_A L_B}^{L_A^\prime L_B^\prime}(R) = V_{L_A L_B}^{L_A^\prime, -M}(R)^* \] (32)

\[ V_{L_A L_B}^{L_A^\prime L_B^\prime}(R) = V_{L_A L_B}^{L_A^\prime, -M}(R)^*. \] (33)

Combining these equations one finds again that all expansion coefficients must be real-valued.
which is also obvious from Eq. (30). Furthermore, the diagonal potentials \( V_{\Lambda',\Lambda}(R, \theta_A, \theta_B, \phi) \) and \( V_{\Lambda,\Lambda}(R, \theta_A, \theta_B, \phi) \) have equal expansion coefficients

\[
v_{L_A}^{\Lambda,\Lambda}(R) = v_{L_A}^{\Lambda,\Lambda}(R)^*.
\]

The expansion coefficients with \( \Lambda' = \Lambda \) for \(-M\) and \(+M\) are equal as well

\[
v_{L_A}^{\Lambda,\Lambda}(R) = v_{L_A}^{\Lambda,\Lambda}(R) = v_{L_A}^{\Lambda,\Lambda}(R)^*.
\]

and by combining the terms with \(-M\) and \(+M\) in the expansion of Eq. (28) we obtain an expansion of the diagonal potentials in terms of real-valued angular functions

\[
V_{\Lambda,\Lambda}(R, \theta_A, \theta_B, \phi) = \sum_{L_A} \sum_{L_B} (2 - \delta_{M0})(-1)^M \times v_{L_A}^{\Lambda,\Lambda}(R) C_M^L(\theta_A, 0) C_{-M}^L(\theta_B, 0) \cos M \phi.
\]

Here, we used \( D_{M0}^L(\theta, 0) = d_{M0}^L(\theta) = C_M^L(\theta, 0) \) and \( C_{-M}^L(\theta, 0) = (-1)^M C_M^L(\theta, 0) \).

**F. Multipole expansion**

For large distances \( R \) between two molecules we can write the intermolecular interaction operator \( V \) in the form of the multipole expansion

\[
V = \sum_{L_A} \sum_{L_B} (2 - \delta_{M0})(-1)^M (2L_A + 2L_B) \frac{1}{2L_A} R^{-L_A-L_B-1} \times \sum_{M_A=-L_A}^{L_A} \sum_{M_B=-L_B}^{L_B} (-1)^{M_A+M_B} \times \langle L_A, M_A; L_B, M_B | L_A + L_B, M_A + M_B \rangle \times C_{L_A}^{L_B}(\theta, \phi) C_{-M_A}^{M_B}(\Theta, \Phi),
\]

where \( Q_{L_A}^{L_B}(\theta, \phi) \) and \( Q_{M_B}^{M_B}(\Theta, \Phi) \) are spherical multipole operators depending on the electronic and nuclear coordinates of the diatom \( A \) and \( B \), respectively. The vector \( \mathbf{R} = (R, \Theta, \Phi) \) points from \( A \) to \( B \). Since multipole operators are irreducible tensor operators it follows for both diatoms that

\[
Q_{L_A}^{L_B}(\theta, \phi) = \hat{R}(\theta, \phi) Q_{L_B}^{L_B}(\theta, \phi) = \sum_k Q_{k}^{L_A} D_{M_A}^{L_B}(\phi, \theta, 0)^*.
\]

For HCl in its \( \Sigma \) ground state only the \( K_B = 0 \) component \( \langle Q_{00}^{L_B} \rangle \) is nonvanishing, so

\[
\langle Q_{L_B}^{L_B}(\theta_B, \phi_B) \rangle = \langle Q_{00}^{L_B} \rangle D_{M_B}^{L_B}(\phi_B, \theta_B, 0)^* = \langle Q_{00}^{L_B} \rangle C_{M_B}^{L_B}(\theta_B, \phi_B).
\]

For OH (=A) we must evaluate matrix elements

\[
\langle \Lambda'(\theta_A, \phi_A) | Q_{L_A}^{L_A} | \Lambda(\theta_A, \phi_A) \rangle = \langle \Lambda'(0,0) | \hat{R}(\theta_A, \phi_A) | Q_{L_A}^{L_A} \hat{R}(\theta_A, \phi_A) | \Lambda(0,0) \rangle
\]

\[
= \langle \Lambda'(0,0) | Q_{L_A}^{L_A} \Lambda(0,0) | D_{M_A}^{L_A} \Lambda(\phi_A, \theta_A, 0) \rangle^*.
\]

Computing matrix elements of Eq. (38) and taking \( R \) along the \( z \)-axis, so that \( M_A = M_B = 0 \), we find precisely the expansion of Eq. (27) with expansion coefficients expressed in closed form

\[
v_{L_A}^{\Lambda,\Lambda}(R) = \delta_{L_A+L_B}(1) (2L_A + 2L_B) \frac{1}{2L_A} \langle \Lambda' | Q_{L_A}^{L_A} \Lambda \rangle \times \langle Q_{00}^{L_A} \rangle R^{-L_A-L_B-1}.
\]

Hence, in first-order long-range theory one can directly compute the DPESs \( V_{\Lambda,\Lambda}(R, \theta_A, \theta_B, \phi) \). The long range diabatic angle \( \gamma \) can be extracted from the off-diagonal matrix element \( V_{1,-1} \). This element is complex-valued and its phase angle equals \( 2\gamma \), see Eq. (4).

Note that the dominant long-range contribution to the diagonal potentials \( V_{1,1} = V_{1,-1} \) in the OH-HCl complex is the dipole-dipole interaction which depends on the distance as \( R^{-3} \). This term does not occur in the off-diagonal potentials \( V_{1,-1} = V_{1,-1} \); the dominant long-range contribution to \( V_{1,-1} \) is the quadrupole-dipole interaction proportional to \( R^{-4} \), which contains the off-diagonal quadrupole matrix element \( \langle \Pi_1 | \hat{Q}_0^2 | \Pi_{-1} \rangle \) of the OH radical in its \( ^3\Pi \) ground state. Also the higher multipole operators of OH provide such off-diagonal matrix elements.

**III. COMPUTATIONS**

As pointed out in the introduction, we computed the lowest energy by the RCCSD(T) method and the first excitation energy by the MRCI method. Accordingly, the adiabatic energies are defined thus,

\[
V_1 = E_{\text{RCCSD(T)}}(\text{OH} - \text{HCl}) - E_{\text{RCCSD(T)}}(\text{OH}) - E_{\text{RCCSD(T)}}(\text{HCl}),
\]

\[
V_2 = V_1 + \Delta E_{\text{MRCI}},
\]

the idea being that the unlinked cluster contributions in the lowest and second lowest MRCI state mutually cancel to a large extent. (Recall that the size-extensivity problem is caused by a contamination of the CI wave functions by unlinked clusters.) Davidson-corrected\(^{15}\) energies were used throughout this work. To correct for basis set superposition errors the RCCSD(T) calculations were all performed in the dimer orbital basis, while we assume that the BSSEs in \( \Delta E_{\text{MRCI}} \) mutually cancel.

In the RCCSD(T) calculations we used the augmented correlation-consistent polarized-valence double zeta (aug-cc-pVTZ)\(^{14}\) basis enlarged with bond functions placed halfway between the centers of mass of OH and HCl. Exponents of these uncontracted functions are \( sp(0.9,0.3,0.1) \) and \( d(0.6,0.2) \). In the MRCI calculations the bond functions caused problems in extracting diabatic angles, which is why we discarded them and used the larger aug-cc-pVQZ basis.
instead. The 27-electron MRCI calculations were based on natural orbitals obtained from complete active space self-consistent field (CASSCF) calculations with 6 orbitals doubly occupied and 10 orbitals in the active space. Since the contracted MRCI calculations were used only to compute excitation energies and diabatic angles, they were kept fairly modest with 12 doubly occupied orbitals, 4 orbitals in the active space and 126 external orbitals. All calculations reported on this work were performed by the aid of the program MOLPRO.15

The center of mass A of OH was placed in the origin, and the atomic masses are 1H:1.0078250321 u, 16O:15.9949146221 u. The O–H bond length of 1.950 Å was taken from earlier work16,17 on the Ar–OH complex. The usual spherical polar angles of the vector rOH, pointing from O to H, are denoted by θH and φH. The center of mass B of HCl was placed on the positive z-axis at R, the mass of 35Cl (75.8% natural abundance) is 34.96885271 u. The H–Cl bond length of 1.275 Å (≈2.4094 a0) corresponds to the equilibrium distance. The usual spherical polar angles of rCIH, pointing from Cl to H, are denoted by θH and φH. Figure 1, which shows the equilibrium geometries of the OH–HCl dimer (discussed below), illustrates the coordinates. The computations were performed for 24 distances: R = 25, 22, 20, 18, 14, 10, 8, 6, 4.5, 4.25, 4, 3.75, 3.5, 3.25, 3.15, 3.05, 3, 2.75, 2.65, 2.5, 2.35, 2.25, and 2 Å. The angles θH and θH were varied over 0°, 15°, 30°, 45°, 60°, 80°, 90°, 100°, 120°, 135°, 150°, 165°, and 180°, while the dihedral angle φ = φH−φH was varied from 0° to 180° in steps of 45°. In total we computed 20280 points on the PES. (Actually, in this number a few geometries are counted more than once. For instance, when HCl is along the z-axis the interaction does not depend on φ.)

In order to extract the diabatic angle γ from natural orbitals and CI vectors by Eq. (1), it is necessary to first let the diatomic OH axis coincide with the SF z-axis. We used φH = 0 (OH in xz-plane) and φ = φH so that it is required to rotate the vectors R, rCIH, and rOH actively around the y-axis over an angle −θH. This rotation was performed on the Cartesian input vectors to the MOLPRO program. Next we must realize that the two lowest CI vectors, Ψ1 and Ψ2, are linear combinations of two dominant configuration state functions (CSFs), which can be seen as states with a hole in a π1 and π2 natural orbital created out of a closed-shell reference state |Φ0⟩. Of course, CSFs of non-II symmetry contribute to Ψ1 and Ψ2 as well, but as already discussed, these are of minor importance. The two highest occupied natural orbitals are given to a very good approximation by

$$(|\pi_1\rangle, |\pi_2\rangle) = (|\pi_i\rangle, |\pi_i\rangle) R(\gamma_{\text{orb}}),$$  \hspace{1cm} (43)

where πi and πi are linear combinations of pi and p j AO's, respectively. Introducing corresponding hole creation operators a†i and a†i, we can write the CI vectors (adiabatic states) in terms of the hole states a†i[|Φ0⟩] and a†i[|Φ0⟩].

$$(|\Psi_1\rangle, |\Psi_2\rangle) = (a†_i, a†_i) R(\gamma_{\text{CI}})|\Phi_0⟩,$$  \hspace{1cm} (44)

According to Eq. (43) and realizing that for an orthogonal transformation particles and holes transform the same,

$$(a†_i, a†_i) = (a†_i, a†_i) R(\gamma_{\text{orb}}) R(\gamma_{\text{CI}})|\Phi_0⟩,$$ 

so that

$$(|\Psi_1\rangle, |\Psi_2\rangle) = (a†_i, a†_i) R(\gamma_{\text{orb}}) R(\gamma_{\text{CI}})|\Phi_0⟩,$$  \hspace{1cm} (46)

and γ = γ_{\text{orb}} + γ_{\text{CI}}.

By inspection of the dominant coefficients of the two highest natural orbitals and of the lowest two CI vectors the diabatic angles γ_{\text{orb}} and γ_{\text{CI}} can be obtained from Eqs. (43) and (44). In practice we put all non-relevant (small) coefficients in the two orbitals and in the two states equal to zero and re-orthonormalized the two orbitals and the two states, before extracting the diabatic angles. This re-orthonormalization was done by the Löwdin algorithm.19 It was checked that the angles obtained by an alternative orthonormalization algorithm, the Gram–Schmidt method, did not differ from our values by more than four degrees, even for the shortest distances where II symmetry breaking is largest. Substitution of γ and the adiabatic energies V1 and V2 [cf. Eq. (42)] into Eq. (4) gives finally the DPESs V_{1,1} ± V_{-1,-1} and V_{1,-1} ± V_{-1,1}.

We also computed the diabatic potentials V_{1,1} and V_{1,-1} directly from the multipole expansion, see Eq. (41) for the expansion coefficients. The multipole matrix elements of OH in its II ground state and the multipole moments of HCl were calculated from MRCI wave functions with the MOLPRO program. The monomer basis sets were the same as used in the dimer MRCI calculations and also the choice of monomer orbitals that were kept doubly occupied or put into the active space was the same as in the dimer calculations. For OH we computed matrix elements $\langle \Lambda' | Q^{\ell}_{\alpha \beta} | \Lambda \rangle$ with $\Lambda' = \pm 1$ for $\Lambda = 1, 2, 3$. For HCl we computed multipole moments $Q^{\ell}_{\alpha \beta}$ for $L = 2, 3, 4$. The values obtained are listed in Table I. Hence, the long range diabatic potentials that were obtained from the multipole expansion contain terms up to R^{-7}, inclusive, but only the terms up to R^{-5} inclusive are complete. We already observed that the leading term in the diagonal potential V_{1,1} is the dipole-dipole term $\propto R^{-3}$, whereas the off-diagonal potential V_{1,-1} has the quadrupole-dipole term $\propto R^{-4}$ as its leading long range contribution.
respectively, were determined. Most of the programs that "Wigner well-known and easily calculated by recursion; they are also

\[
V_{1,1} = \begin{cases} 
V_{1,1} & \text{for } V_{1,1} \leq V_0 \\
V_0 + \beta^{-1}\tanh(\beta(V_{1,1} - V_0)) & \text{for } V_{1,1} > V_0
\end{cases}
\]

(48)

where \( \beta = [V_{\text{max}} - V_0]^{-1} \). With this scheme, the damped potential \( V_{1,1} \) is continuous around \( V_0 \) up to the second derivative. Care was taken to use sufficiently high values of \( V_0 \) and \( V_{\text{max}} \), so that the potential was affected only in regions that are not of any practical importance in bound state and scattering calculations. The actual values used were \( V_0 = 10,000 \text{ cm}^{-1} \), and \( V_{\text{max}} = 2V_0 \). At those grid points where \( V_{1,1} \) was damped the off-diagonal potential \( V_{1,-1} \) was scaled by the same factor \( V_{1,1}/V_{1,1} \).

We experimented with two different fitting strategies: (i) radial fits of the energies \( V_{1,1} \) and \( V_{1,-1} \) followed by angular expansions of the radial fit parameters and (ii) angular fits of the energies followed by radial fits of the angular fit parameters. Overall the second strategy gave the more accurate fit and—after the fit parameters are determined—the fastest program for generation of the fitted DPESS. For both the real \( V_{1,1} \) and the complex \( V_{1,-1} \) the fit was first performed for the coefficients in the \( \text{LLM} \) expansion in Eq. (28).

Since the \( \text{LLM} \) expansion is in terms of orthogonal functions, each expansion (or Fourier) coefficient can be looked upon as an overlap integral between the function to be expanded and one of the expansion functions. This means that the coefficients can be obtained by numerical integration. Before actually performing this integration we generated values of the potentials \( V_{1,1} \) and \( V_{1,-1} \) on a 10-point Gauss–Legendre (GL) quadrature grid for both \( \theta_A \) and \( \theta_B \) by interpolating the \( 13 \times 13 \) \( \{\theta_A, \theta_B\} \) \textit{ab initio} points. A two-dimensional cubic spline interpolation method was used for this purpose. In order to avoid the oscillations that may occur in a cubic spline method especially near the boundaries of the interval to be interpolated we extended the range of \( \theta_A \) and \( \theta_B \) from \( 0 \leq \theta \leq \pi \) to \( -\pi \leq \theta \leq 2\pi \). The values of the potentials for \( -\pi \leq \theta \leq 0 \) and for \( \pi \leq \theta \leq 2\pi \) were obtained from the original values for \( 0 \leq \theta \leq \pi \) by means of the transformations

\[
V_{1,\pm 1}(R, \theta_A, \theta_B, \phi) = V_{1,\pm 1}(R, 2\pi - \theta_A, \theta_B, \pi + \phi)
\]

(49)

which are simply a consequence of the periodic boundary conditions in spherical polar coordinates, and the property that the addition of \( \pm \pi \) to any of these angles makes no difference. In addition, we used that \( V_{1,1}(R, \theta_A, \theta_B, \phi) = V_{1,1}(R, \theta_A, \theta_B, -\phi) \) and \( V_{1,-1}(R, \theta_A, \theta_B, \phi) = V_{1,-1}(R, \theta_A, \theta_B, -\phi) \), which follows from the reflection symmetry treated in Sec. II C. Next the integration was performed by means of a 10-point GL quadrature for \( \theta_A \) and \( \theta_B \) (in the range \( 0 \leq \theta_A, \theta_B \leq \pi \)) and a 5-point trapezoidal rule for \( \phi \). From Eq. (36) follows that in the \( \text{LLM} \) expansion of the diagonal \( (\Lambda' = \Lambda) \) potential \( V_{1,1} \) only non-negative \( M \) values appear. Expansion coefficients \( V_{1,1,LM}(R) \) with maximum \( L_A, L_B = 7 \) and \( M = 0, 1, 2, 3, 4 \) were computed. In the off-diagonal potential \( V_{1,-1} \) also negative \( M \) values appear. We computed expansion coefficients \( V_{1,-1,LM}(R) \) with maximum \( L_A, L_B = 7 \) and \( M = -4, -3, \cdots, 4 \). All of these coefficients were obtained on

### IV. FIT OF THE DPESS

In this section will be described how the expansion coefficients \( v_{LM}^{\Lambda,\Lambda} \) and \( v_{LM}^{\Lambda',\Lambda} \), appearing in Eqs. (28) and (27), respectively, were determined. Most of the programs that perform this task were coded in the commercial program package Matlab. Both expansions of the DPESSs require associated Legendre polynomials \( P^L_M(\cos \theta) \) and "Wigner d functions" \( D^j_{m'm}(\theta, 0, 0) \) (also known as "reduced rotation matrix elements").

The functions \( P^L_M \) are well-known and easily calculated by recursion; they are also part of Matlab. The \( d \) functions may be written as

\[
d^j_{m'm}(\theta) = (-1)^k \frac{(2j - k)!}{(k + \alpha)!} \frac{\alpha!}{\beta!} (\sin \frac{\theta}{2})^\alpha \\
\times \left( \cos \frac{\theta}{2} \right)^\beta p_k^{\alpha,\beta}(\cos \theta),
\]

(47)

where \( p_k^{\alpha,\beta}(\cos \theta) \) is a Jacobi polynomial. Here \( k = \min(j + m, j - m, j + m', j - m') \). The phase \( \lambda \) and the index \( \alpha \) are given by

\[
\begin{cases}
  j + m & : \alpha = m' - m; \lambda = m' - m \\
  j - m & : \alpha = m - m'; \lambda = 0 \\
  j + m' & : \alpha = m' - m; \lambda = 0 \\
  j - m' & : \alpha = m - m'; \lambda = m' - m
\end{cases}
\]

The index \( \beta = 2j - 2k - \alpha \). Fortran-90 and Matlab routines for a variety of polynomials are available on the internet, among which is a routine for Jacobi polynomials. Use of this routine together with Eq. (47) gives a stable and efficient algorithm for the generation of the \( d \) functions.

Since for the smallest values of \( R \) the potential becomes extremely repulsive for certain orientations of the OH and HCl monomers, one would need very many terms in the expansion. To avoid this, the potential \( V_{1,1} \) was damped in these strongly repulsive regions by means of a \( \tanh \) function up to a value \( V_{\text{max}} \).
the $R$ grid. Note that the expansion coefficients are real-valued, even for the complex-valued potential $V_{1,1}$; the overlap integrals to be evaluated in that case contain both the real and imaginary parts of $V_{1,1}$ and the real and imaginary parts of the expansion functions.

Once the coefficients $u_{LAM}^{A}(R)$ in the $LLM$ expansions were obtained, the coefficients $u_{LAM}^{A}(R)$ in the $LLL$ expansions were computed from Eq. (30). As just stated, in the case of the diagonal potential $V_{AA}$ the $LLM$ coefficients on the right-hand side of this equation are only determined for non-negative $M$. From the properties of the expansion coefficients derived in Sec. II C it follows that the diagonal coefficients $u_{LAM}^{AA}$ vanish unless $L_{A}+L_{B}+L$ is even. For given $L_{A},L_{B}$ the number of $L$ values with even $L_{A}+L_{B}+L$ allowed by the triangular relation is equal to the number of non-negative $M$ values. Separating the summation in Eq. (30) in one over positive $M$ and one over negative $M$, using Eq. (36), we find that for odd $L_{A}+L_{B}+L$ the expansion coefficients $u_{LAM}^{AA}$ indeed vanish and that for even $L_{A}+L_{B}+L$:

$$u_{LAM}^{AA}(R) = \sum_{M=0}^{\infty} (2 - \delta_{M0})(L,0 \mid L_{A},M;L_{B},-M)u_{LAM}^{AA}(R).$$

Finally, the angular expansion coefficients $u_{LAM}^{AA}(R)$ were fitted as functions of $R$ by means of the reproducing kernel Hilbert space (RKHS) method\textsuperscript{24} with the reproducing kernel for distance-like variables. The dominant proportional ($\propto R^{-2}$) contribution to the diagonal potential $V_{1,1}$ in the long range is given by the dipole-dipole interaction with $L_{A}=1$ and $L_{B}=1$, see Sec. II F. The dominant ($\propto R^{-2}$) long range contribution to the off-diagonal potential $V_{1,-1}$ is the quadrupole-dipole interaction with $L_{A}=2$ and $L_{B}=1$. All contributions with $L_{A}+L_{B}<5$, which decay slower than the leading induction and dispersion terms $\propto R^{-6}$, were fitted with RKHS parameter $m=L_{A}+L_{B}$, so that they decay as $R^{-L_{A}-L_{B}-1}$ (Ref. 24) beyond the outermost grid point. All expansion coefficients with $L_{A}+L_{B}+1\geq 6$ were fitted with RKHS parameter $m=5$ and decay as $R^{-6}$ for very large $R$. The smoothness parameter $n$ was always 2. The Fortran code with the fitted diabatic potentials is available from the authors upon request.

V. RESULTS AND DISCUSSION

Before we discuss the potentials obtained we mention some results regarding the accuracy of the analytic potentials with respect to the original $ab$ initio data. For a distance $R$ of 3.5 Å the mean absolute error in the diagonal diabatic potential $V_{1,1}$ is 30.8 cm\textsuperscript{-1}, which is about 8\% of the average magnitude of $V_{1,1}$ at this distance. The mean absolute error in the off-diagonal potential $V_{1,-1}$ is 11.5 cm\textsuperscript{-1}, while the average of $|V_{1,-1}|$ is 82 cm\textsuperscript{-1} at $R=3.5$ Å. At $R=10$ Å the mean absolute error in $V_{1,1}$ is 0.4 cm\textsuperscript{-1}, which is about 4\% of the magnitude of $V_{1,1}$. The mean absolute error in $V_{1,-1}$ at $R=10$ Å is 0.05 cm\textsuperscript{-1}, while the average of $|V_{1,-1}|$ is 0.7 cm\textsuperscript{-1}.

First we discuss the diabatic potentials at $R=3.5$ Å, which is close to the distance of the hydrogen-bonded equilibrium geometry. Figure 2 shows the diagonal potential $V_{1,1}$ for planar geometries with $\phi=0^\circ$ and $\phi=180^\circ$, while Fig. 3 shows the same potential at dihedral angle $\phi=90^\circ$. The minimum of $\approx 901$ cm\textsuperscript{-1} occurs at $\phi=180^\circ$, $\theta_{\text{CIH}}=159^\circ$, and $\theta_{\text{HCH}}=176^\circ$. When relaxing the distance $R$ we find a binding energy of 905 cm\textsuperscript{-1} at $R=3.45$ Å and nearly the same angles. This corresponds to a nearly linear hydrogen bonded structure with HCl being the proton donor in the hydrogen bond and OH being the proton acceptor. The donor HCl molecule lies nearly along the bond, the acceptor OH radical makes an angle of $150^\circ$ with the bond axis. This OH angle will change substantially when we include the off-diagonal potential and consider the minimum in the lowest adiabatic potential, see Fig. 1 and the discussion below. A second (meta)stable struct-
ture occurs for the hydrogen bonded geometry with $\phi = 180^\circ$, $\theta_{OH} = 0^\circ$ and $\theta_{CH} = 87^\circ$. This geometry corresponds to a hydrogen bond with OH as the donor and HCl as the acceptor. The binding energy of 653 cm$^{-1}$ at the optimized distance $R = 3.52$ Å is substantially smaller than for the most stable hydrogen bonded geometry. As in the HCl-HCl dimer, the accepting HCl molecule is nearly perpendicular to a hydrogen bond with OH as the donor and HCl as the acceptor. The binding energy of 653 cm$^{-1}$ at the optimized geometry with $\phi = 180^\circ$, i.e., for T-shaped geometries with HCl along the bond axis, but the plane of reflection has rotated over an angle $\phi$ with respect to the $xz$-plane. Hence, also the adiabatic states have rotated over $\phi$ with respect to the bond axis.

The largest absolute values occur near $\theta_{OH} = 90^\circ$ and $\theta_{CH} = 180^\circ$, i.e., for T-shaped geometries with HCl along the bond axis and its hydrogen atom pointing to the OH center of mass. For planar geometries with $\theta_{CH} = 0^\circ$ or $180^\circ$ the potential $V_{1,-1}$ must be real-valued for all values of $\theta_{OH}$, because of the $A'$ and $A''$ reflection symmetry of the real diabatic states $|\Pi_\parallel\rangle$ and $|\Pi_\perp\rangle$ and the fact that the plane of reflection is the $xz$-plane. For linear geometries, with both OH and HCl lying along the bond axis, the $|\Pi_\parallel\rangle$ and $|\Pi_\perp\rangle$ states are degenerate and the off-diagonal potential $V_{1,-1}$ vanishes altogether. For planar geometries with $\theta_{OH} = 0^\circ$ and HCl not lying along the bond axis, the diabatic mixing angle $\gamma$ must equal the dihedral angle $\phi$. Analogously, $\gamma = -\phi$ when $\theta_{OH} = 180^\circ$. Since the phase of the potential $V_{1,-1}$ is exp$(2i\gamma)$, see Eq. (4), this potential is real-valued at $\phi = 90^\circ$ and purely imaginary at $\phi = 45^\circ$ and $135^\circ$. This is a consequence of the $A'$ and $A''$ reflection symmetry of the adiabatic states when the complex is planar. For $\phi = 0^\circ$ the adiabatic states coincide with the real diabatic states $|\Pi_\parallel\rangle$ and $|\Pi_\perp\rangle$, whereas for other values of $\phi$ the complex is still planar when OH lies along the bond axis, but the plane of reflection has rotated over an angle $\phi$ with respect to the $xz$-plane. Hence, also the adiabatic states have rotated over $\phi$ with respect to the bond axis.

Figure 4 shows the off-diagonal potential $V_{1,-1}$ at $R = 3.5$ Å for planar geometries with $\phi = 0^\circ$ and $180^\circ$, where it is real-valued. Figure 4 shows the real and imaginary parts of this potential at dihedral angle $\phi = 90^\circ$. It is clear that the imaginary part is as important as the real part. The largest absolute values occur near $\theta_{OH} = 90^\circ$ and $\theta_{CH} = 180^\circ$, for T-shaped geometries with HCl along the bond axis and its hydrogen atom pointing to the OH center of mass. For planar geometries with $\theta_{CH} = 0^\circ$ or $180^\circ$ the potential $V_{1,-1}$ must be real-valued for all values of $\theta_{OH}$, because of the $A'$ and $A''$ reflection symmetry of the real diabatic states $|\Pi_\parallel\rangle$ and $|\Pi_\perp\rangle$ and the fact that the plane of reflection is the $xz$-plane. For linear geometries, with both OH and HCl lying along the bond axis, the $|\Pi_\parallel\rangle$ and $|\Pi_\perp\rangle$ states are degenerate and the off-diagonal potential $V_{1,-1}$ vanishes altogether. For planar geometries with $\theta_{OH} = 0^\circ$ and HCl not lying along the bond axis, the diabatic mixing angle $\gamma$ must equal the dihedral angle $\phi$. Analogously, $\gamma = -\phi$ when $\theta_{OH} = 180^\circ$. Since the phase of the potential $V_{1,-1}$ is exp$(2i\gamma)$,
= 180°. By comparison with Eq. (4) it follows that the diabatic angle $\gamma$ equals $\phi$ when $\theta_{\text{OH}} = 0°$ and $-\phi$ when $\theta_{\text{OH}} = 180°$.

From Eq. (4) it follows that the lowest adiabatic potential is simply obtained by subtracting the absolute value of $V_{1,1}$ from $V_{1,-1}$. Figure 6 gives a picture of the lowest adiabatic potential for planar geometries with $\phi = 0°$ and 180°. This figure shows the minimum value of the lowest adiabatic potential obtained by varying the distance $R$ for each set of angles $\theta_{\text{OH}}, \theta_{\text{CH}}, \phi$. The corresponding equilibrium values of $R$ are plotted in Fig. 7. The geometry of the most stable hydrogen bonded structure, see Fig. 1(a), has substantially changed with respect to the minimum in the diagonal potential in Fig. 2. The HCl molecule is still the donor and lies nearly parallel with the bond axis ($\theta_{\text{CH}} = 176.4°, \phi = 180°$), but the acceptor OH now makes an angle $\theta_{\text{OH}}$ of 110.7° with the bond axis. The binding energy is $D_e = 1123 \text{ cm}^{-1}$ and the equilibrium distance is $R_e = 3.366 \text{ Å}$. We still find a second, metastable, hydrogen bonded structures, see Fig. 1(b), with OH as the donor and HCl as the acceptor at $\theta_{\text{OH}} = 5.5°, \theta_{\text{CH}} = 86.9°, \phi = 180°, R_e = 3.517 \text{ Å}$, and $D_e = 655 \text{ cm}^{-1}$.

The earlier calculations of the potential energy surface of $\text{OH} + \text{HCl} \rightarrow \text{H}_2\text{O} + \text{Cl}$ (Refs. 2–6) concentrate on the transition state for the reaction in which the HCl bond is substantially elongated with respect to free HCl, so we cannot compare the results with those of the present paper. The paper by Yu and Nyman also contains $ab$ $initio$ results for the geometry and energetics of the OH–HCl entrance channel complex. These authors performed second-order Møller–Plesset (MP2) calculations on the planar complex in a fairly small basis. They correct for the omission of BSSE corrections and incomplete correlation by a scaling method. Thus, they find a minimum energy of the complex of $5.46$ kcal/mol = 1910 cm$^{-1}$ at (in our coordinates) $\theta_{\text{OH}} = 112.8°$ and $\theta_{\text{CH}} = 163.5°$ with a distance between the centers of mass of 4.334 $a_0 = 2.294 \text{ Å}$. These equilibrium angles are not very different from ours, but compared to our results their minimum is considerably too deep and falls at too short a distance. The most likely source of this discrepancy is that their scaling method does not correct sufficiently for the BSSE.

In Figs. 8–11 we display the diabatic potentials at $R = 10 \text{ Å}$ and compare them with the corresponding potentials directly calculated from the multipole expansion with the multipole moments of Table I and the formulas of Sec. II F. Both the diagonal potential $V_{1,1}$ and the off-diagonal potential $V_{1,-1}$ agree very well with their multipole-expanded counterparts. One observes this for planar geometries with $\phi = 0°$ and 180° in Fig. 8 for the diagonal potential $V_{1,1}$ and in Fig. 9 for the off-diagonal potential $V_{1,-1}$. It holds also for non-planar geometries where the off-diagonal potential $V_{1,-1}$ is complex-valued, see Fig. 10 for $V_{1,1}$ and Fig. 11 for the real and imaginary parts of $V_{1,-1}$. At this large distance of 10 Å the dipole-dipole interactions dominate the potential $V_{1,1}$ and the most stable structure corresponds to $\theta_{\text{OH}} = 180°$ and $\theta_{\text{CH}} = 180°$, i.e., to a linear geometry with the HCl and OH dipoles both aligned along the bond axis. The largest absolute (negative) values of $V_{1,1}$, the leading term of which is the OH quadrupole–HCl dipole interaction, occur for $\theta_{\text{OH}} = 90°$ and $\theta_{\text{CH}} = 180°$. 

FIG. 6. Lowest adiabatic potential (in cm$^{-1}$) for optimized $R$ at planar geometries with $\phi = 180°$ and $\phi = 0°$, put together as described in Fig. 2.

FIG. 7. Equilibrium values of $R$ for the lowest adiabatic potential in Fig. 6 at planar geometries with $\phi = 180°$ and $\phi = 0°$, put together as described in Fig. 2.
In Fig. 12 we show that the multipole-expanded potentials agree with the \textit{ab initio} values for large $R$. For the OH and HCl angles in these plots we chose the angles of the most stable hydrogen-bonded structure. With decreasing $R$ the attractive dispersion interactions, not included in the first-order multipole expansion, become important and at small $R$ the repulsive exchange interactions dominate.

VI. SUMMARY AND CONCLUSIONS

We have reported \textit{ab initio} computations of the diabatic potential energy surfaces (DPESs) $V_{1,1}(R, \theta_A, \theta_B, \phi)$ and $V_{1,-1}(R, \theta_A, \theta_B, \phi)$ for the dimer OH$(^2\Sigma)$--HCl. The DPESs are obtained from the lowest two adiabatic surfaces by combining results from RCCSD(T) and MRCI calculations. The
RCCSD(T) method gives the ground state energy, while the first excitation energy is computed by the MRCI method. Diabatic states are obtained from the lowest two adiabatic states by a $2 \times 2$ rotation with rotation angle $\gamma$, the so-called "diabatic angle". The diabatic angle is obtained from inspection of CASSCF natural orbitals and MRCI states.

The two DPESs thus obtained are expanded in terms of orthogonal polynomials depending on the angles $\theta_A, \theta_B$, and $\phi$. The form of this expansion is suggested by the electrostatic interaction between the monomers, one of which has a spatially degenerate II ground state. It is shown that the ab initio results converge for large $R$ to the electrostatic interactions, so that the expansion of the DPES coincides asymptotically with the multipole expansion of the electrostatic interaction.

Some experimentation with the numerical calculation of the angular expansion coefficients led to a scheme in which the ab initio points are first spline-interpolated to Gauss–Legendre quadrature points. In this interpolation it turns out to be useful to continue the colatitude angles $\theta_A$ and $\theta_B$ on the sphere from $-\pi$ to $2\pi$. The expansion coefficients were then computed by a Gauss–Legendre quadrature, which is possible by virtue of the fact that the expansion is in terms of orthogonal polynomials. Finally, the angular expansion coefficients are given as functions of $R$ by the reproducing kernel Hilbert space method. The diabatic and adiabatic PESs are discussed and some representative cuts are shown.

Bound state and scattering calculations are in progress; comparison of the results of these ongoing computations with experimental data will provide useful information on the OH–HCl entrance channel complex of the chemical reaction $\text{OH} + \text{HCl} \rightarrow \text{H}_2\text{O} + \text{Cl}$ and, at the same time, decide on the quality of the PESs reported on in this paper.

FIG. 10. Diabatic potential $V_{1,1}$ (in cm$^{-1}$) for $\phi=90^\circ$ and $R=10$ Å. From the full ab initio calculations (left panel) and from the multipole expansion (right panel).

FIG. 11. Real and imaginary parts of diabatic potential $V_{1,-1}$ (in cm$^{-1}$) for $\phi=90^\circ$ and $R=10$ Å. From the full ab initio calculations (left panels) and from the multipole expansion (right panels).
FIG. 12. Potentials $V_{1,1}$ (upper panel) and $V_{-1,-1}$ (lower panel) from the full \textit{ab initio} calculations (closed curves) and from the multipole expansion (dashed curves). The angles are the values of the most stable hydrogen-bonded structure: $\theta_{OH}=110.7^\circ$, $\gamma_{HI}=176.4^\circ$, and $\phi=180^\circ$.

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