

# Computation of Forces arising from the Polarizable Continuum Model within the Domain-Decomposition Paradigm

Paolo Gatto,<sup>1</sup> Filippo Lipparini,<sup>1</sup> and Benjamin Stamm<sup>1</sup>

Within implicit solvation models, the domain-decomposition strategy for the computation of the electrostatic energy due to the solvent based on the Polarizable Continuum Model (PCM) has recently been developed. The methodological development started with the so-called ddCOSMO method and has recently be generalized to the PCM equation resulting in the ddPCM-method [Stamm *et al.*, J. Chem. Phys. 144, 054101 (2016)] for which derive the forces within this article. We show the derivation of the forces and derive an efficient implementation followed by numerical tests.

## I. INTRODUCTION

The effects of solvation play a crucial role in many processes in chemistry and biochemistry. While the individual solvent molecules in a solute-solvent system can not always be taken explicitly into account, in particular when a quantum-mechanical description of the molecular system is considered, solvation effects can not be neglected altogether. Thus, a compromise is to resort to a limited description of the solvent, as the solvent *per se* is not of interest, but, rather, its effect on the solute molecule is.

Polarizable Continuum Solvation Models (PCSM's)<sup>1-18</sup> focus on the electrostatic interaction between the solute and the solvent by replacing the solvent with an infinite continuum with dielectric permittivity that matches the bulk permittivity of the solvent. Such approach turns out to be a good compromise between accuracy and the cost to compute electric, magnetic, vibrational and mixed properties<sup>1,5,19</sup>. As a result, PCSM's are nowadays available in most quantum Chemistry codes, and have become a standard tool which is successfully employed by an extended community.

Standard implementations of PCSM's usually employ the Boundary Element Method<sup>1,20-24</sup> (BEM) to numerically solve a discretized integral equation. In practical terms, this requires the solution of a linear system whose size scales linearly, although with a large proportionality constant (a rough estimate would be 50-100), with respect to the number of atoms. The system is usually solved through standard dense linear algebra techniques, such as the LU decomposition<sup>25</sup>, which require a computational effort of cubic complexity. Consequently, the solution step can rapidly become demanding when dealing with systems as large as those treated via QM/MM methods, and one should seek other strategies. An alternative is provided by iterative techniques, which reduce the computational cost to that of several matrix-vector multiplications<sup>26</sup>, which is, in general, quadratic. When fast summation techniques, e.g., the Fast Multipole Method (FMM)<sup>27</sup>, can be employed, the computational complexity can be further reduced. Thus, it is possible to solve the PCSM linear equations in a number of floating point operations which is linear with respect to the number of atoms. Nevertheless, the solution of the PCSM equations can still represent a formidable bottleneck for large systems<sup>28</sup>, especially when repeated computations are required for statistical sampling purposes or time-dependent simulations.

In recent years, we have proposed an alternative approach for PCSM's which is based

on domain-decomposition. We introduced a completely different strategy<sup>29</sup>, referred to as ddCOSMO<sup>28-31</sup>, to solve the PCSM equation for the Conductor-like Screening Model<sup>8</sup> (COSMO) in combination with Van der Waals molecular cavities. As a first step, the COSMO equation in its differential form is rewritten as a system of coupled linear differential equations on each sphere, where the coupling occurs only between overlapping spheres. Secondly, each differential equation is recast as an integral equation, so that it can be efficiently solved by using a (truncated) expansion of spherical harmonics<sup>29</sup>. The discretization produces a block-sparse linear system<sup>30</sup>, where only the blocks corresponding to overlapping spheres are nonzero. This structure allows for a computational cost that scales linearly with respect to the number of atoms, and is overall very small as compared to previous implementations. Indeed, as shown in<sup>28</sup>, as many as two or three orders of magnitude are gained by the ddCOSMO approach.

More recently, the method has been generalized to the Polarized Continuum Model (PCM) equation, which assumes a finite permittivity of the solvent. This resulted in the ddPCM method<sup>32</sup>, which is based on the same domain-decomposition approach. While the method has been developed in<sup>32</sup> to compute the electrostatic contribution to the solvation energy, the aim of this article is to present the derivation of analytical forces of the ddPCM solvation energy.

This paper is organized as follows. Section II reviews the ddPCM and ddCOSMO methods that we have previously developed. In Section III we describe the derivation of the ddPCM forces and discuss their efficient implementation. Section IV is devoted to numerical experiments. Finally, in Section V we draw conclusions from the presented work and point to possible future directions of research.

## II. A BRIEF REVIEW OF THE DDPCM STRATEGY

The foundation of Polarizable Continuum Solvation Models (PCSM's) is the assumption that the solvent in a solute-solvent system can be treated as either a dielectric, or a conducting continuum medium on the outside of the molecular cavity  $\Omega$  of the solute. We follow the customary approach of taking the cavity to be the so-called Van der Waals cavity<sup>1</sup>, i.e., the union of spheres centered at each atom with radii coinciding with the van der Waals radii. Within this approach, the topologically similar Solvent Accessible Surface (SAS) cavity can

be treated as well. Models based on the Solvent Excluded Surface (SES) have also been proposed<sup>33–36</sup>, but will not be considered here.

The electrostatic part of the solute-solvent interaction is given by  $E_s = \frac{1}{2} f(\varepsilon) \int_{\Omega} \rho(x) W(x) dx$ , where  $f(\varepsilon)$  is an empirical scaling that depends on the dielectric constant of the solvent (and which is only applied in the case of the COSMO),  $\rho$  is the charge density of the solute, and  $W$  is the polarization potential of the solvent. The quantities  $W$  and  $E_s$  are usually referred to, respectively, as the reaction potential and the electrostatic contribution to the solvation energy.

The reaction potential is defined as  $W = \varphi - \Phi$ , where  $\varphi$  is the total electrostatic potential of the solute-solvent system and  $\Phi$  is the potential of the solute *in vacuo*. In the case of the PCM, the total potential  $\varphi$  satisfies a (generalized) Poisson equation with suitable interface conditions<sup>11,12</sup>. Indeed, if  $\varepsilon_s$  is the macroscopic, zero-frequency relative dielectric permittivity of the solvent, and define  $\varepsilon(x) = 1$  when  $x \in \Omega$  and  $\varepsilon(x) = \varepsilon_s$  otherwise, the reaction potential fulfills

$$\begin{cases} \Delta W = 0 & \text{in } \mathbb{R}^3 \setminus \Gamma \\ [W] = 0 & \text{on } \Gamma \\ [\varepsilon \partial_{\nu} W] = (\varepsilon_s - 1) \partial_{\nu} \Phi & \text{on } \Gamma \end{cases} \quad (1)$$

Here  $\Gamma = \partial\Omega$  is the boundary of the cavity,  $\partial_{\nu}$  is the normal derivative on  $\Gamma$ , and  $[\cdot]$  is the jump operator (inside minus outside) on  $\Gamma$ .

Recalling potential theory,  $W$  can be represented as  $W(x) = (\tilde{\mathcal{S}}\sigma)(x)$  when  $x \in \mathbb{R}^3 \setminus \Gamma$ , or  $W(s) = (\mathcal{S}\sigma)(s)$  when  $s \in \Gamma$ . The surface density  $\sigma$  defined on  $\Gamma$  is the so-called apparent surface charge,  $\tilde{\mathcal{S}}$  is the single layer potential and  $\mathcal{S}$  is the single layer operator, which is invertible<sup>37</sup>. Note that both  $\tilde{\mathcal{S}}$  and  $\mathcal{S}$  are based on the surface  $\Gamma$ . It can be shown that  $\sigma$  satisfies the equation  $\sigma = 1/4\pi [\partial_{\nu} W]$ , so that it is possible to recast the PCM problem (1) as a single integral equation for  $\sigma$ . In fact, if we define the operators

$$\mathcal{R}_{\varepsilon} = 2\pi \frac{\varepsilon + 1}{\varepsilon - 1} \mathcal{I} - \mathcal{D} \quad , \quad \mathcal{R}_{\infty} = 2\pi \mathcal{I} - \mathcal{D} \quad (2)$$

where  $\mathcal{I}$  is the identity and  $\mathcal{D}$  is the double layer boundary operator (also based on  $\Gamma$ ). It can be shown<sup>1</sup> that the apparent surface charge satisfies

$$\mathcal{R}_{\varepsilon} \mathcal{S} \sigma = -\mathcal{R}_{\infty} \Phi \quad \text{on } \Gamma \quad (3)$$

which is known as the IEF-PCM equation. It involves operators  $\mathcal{R}_\infty$  and  $\mathcal{R}_\varepsilon$ , which are both invertible. Furthermore, when the dielectric constant  $\varepsilon_s$  approaches infinity, the IEF-PCM equation simplifies to  $\mathcal{S}\sigma = -\Phi$  on  $\Gamma$ , which is the Integral Equation Formulation of the Conductor-like Screening Model (COSMO)<sup>16</sup>.

Let us recall how to solve equation (3) within the domain-decomposition paradigm. The first step is to write the IEF-PCM integral equation (3) as a succession of two integral equations, one of which is equivalent to the COSMO equation<sup>38</sup>. Indeed, if we define  $\Phi_\varepsilon = \mathcal{S}\sigma$ , equation (3) becomes

$$\mathcal{R}_\varepsilon \Phi_\varepsilon = \mathcal{R}_\infty \Phi \quad \text{on } \Gamma \quad (4)$$

$$\mathcal{S}\sigma = -\Phi_\varepsilon \quad \text{on } \Gamma \quad (5)$$

The ddPCM strategy is an extension of ddCOSMO in the following sense: first, equation (4) is solved in order to compute the right-hand side  $-\Phi_\varepsilon$  of equation (5); secondly, ddCOSMO is employed to solve equation (5) with the modified potential  $-\Phi_\varepsilon$ , and compute the solvation energy  $E_s$ .

In order to discuss the domain-decomposition approach employed for both steps, let us introduce some notation. As anticipated, we take the cavity  $\Omega$  be the union of  $M$  spheres  $\Omega_j = B(x_j, r_j)$  with boundaries  $\Gamma_j$ . Let  $U_j : \Gamma_j \rightarrow \mathbb{R}$  be the characteristic function of  $\Gamma_j^{\text{ext}} := \Gamma_j \cap \Gamma$ , and define extensions  $\Phi_j, \Phi_{\varepsilon,j} : \Gamma_j \rightarrow \mathbb{R}$  as  $\Phi_j(s) = U_j(s) \tilde{\Phi}(s)$  and  $\Phi_{\varepsilon,j}(s) = U_j(s) \tilde{\Phi}_\varepsilon(s)$  for  $s \in \Gamma_j$ , where  $(\tilde{\cdot})$  indicates the trivial extension to  $\bar{\Omega}$ .

**Step 1.** We enforce the integral equation (4) on  $\Gamma_j^{\text{ext}}$ , along with the boundary condition  $\Phi_{\varepsilon,j} = 0$  on  $\Gamma_j^{\text{int}}$ , by the single local integral equation

$$\alpha(1 - U_j) \Phi_{\varepsilon,j} + U_j \mathcal{R}_\varepsilon \Phi_\varepsilon = U_j \mathcal{R}_\infty \Phi \quad \text{on } \Gamma_j \quad (6)$$

where  $\alpha$  is an arbitrary nonzero scalar, and  $\mathcal{R}_\varepsilon \Phi_\varepsilon$  and  $\mathcal{R}_\infty \Phi$  should be understood as their trivial extensions to  $\Gamma_j$ . Employing the extensions  $\Phi_j, \Phi_{\varepsilon,j}$ , the double layer operator  $\mathcal{D}$  can be decomposed as

$$(\mathcal{D}\Phi)(s) = (\mathcal{D}_j \Phi_j)(s) + \sum_{k \neq j} (\tilde{\mathcal{D}}_k \Phi_k)(s) \quad ; \quad s \in \Gamma_j^{\text{ext}} \quad , \quad j = 1, \dots, M$$

where  $\mathcal{D}_j$  and  $\tilde{\mathcal{D}}_j$  are, respectively, the local double layer operator and the local double layer potential on  $\Gamma_j$ . We refer to<sup>32</sup> for concise details for these operators. Thus, we can also

decompose  $\mathcal{R}_\varepsilon$  in (6) as

$$(\mathcal{R}_\varepsilon \Phi_\varepsilon)(s) = (\mathcal{R}_{\varepsilon,j} \Phi_{\varepsilon,j})(s) + \sum_{k \neq j} (\tilde{\mathcal{R}}_{\varepsilon,k} \Phi_{\varepsilon,k})(s) \quad ; \quad s \in \Gamma_j^{\text{ext}} \quad , \quad j = 1, \dots, M \quad (7)$$

where the local operators  $\mathcal{R}_{\varepsilon,j}$  and  $\tilde{\mathcal{R}}_{\varepsilon,j}$  are defined as

$$\mathcal{R}_{\varepsilon,j} = 2\pi \frac{\varepsilon + 1}{\varepsilon - 1} \mathcal{I} - \mathcal{D}_j \quad , \quad \tilde{\mathcal{R}}_{\varepsilon,j} = -\tilde{\mathcal{D}}_j$$

with the obvious extension to the case  $\varepsilon = \infty$ . If we insert (7) into (6) and select  $\alpha = 2\pi(\varepsilon + 1)/(\varepsilon - 1)$ , we obtain a convenient form in terms of the local double layer potentials and double layer operators:

$$2\pi \frac{\varepsilon + 1}{\varepsilon - 1} \Phi_{\varepsilon,j} - U_j \left( \mathcal{D}_j \Phi_{\varepsilon,j} + \sum_{k \neq j} \tilde{\mathcal{D}}_k \Phi_{\varepsilon,k} \right) = 2\pi U_j \Phi_j - U_j \left( \mathcal{D}_j \Phi_j + \sum_{k \neq j} \tilde{\mathcal{D}}_k \Phi_k \right) \quad \text{on } \Gamma_j \quad (8)$$

This constitutes our domain-decomposition strategy for equation (4). It is important to remark that, because of the summation, every subdomain  $\Omega_j$  interacts with all other subdomains. We anticipate that this contrasts with the ddCOSMO strategy for equation (5).

**Step 2.** The restriction  $W_j := W|_{\Omega_j}$  is harmonic over the subdomain  $\Omega_j$ , thus it can be represented as

$$W_j(x) = (\tilde{\mathcal{S}}_j \sigma_j)(x) \quad , \quad x \in \Omega_j \quad ; \quad W_j(s) = (\mathcal{S}_j \sigma_j)(s) \quad , \quad s \in \Gamma_j \quad (9)$$

where  $\sigma_j$  is an unknown surface charge, and  $\mathcal{S}_j$  and  $\tilde{\mathcal{S}}_j$  are, respectively, the single layer potential and the single layer operator on  $\Gamma_j$ . The local problems (9), are coupled together by decomposing  $W_j$  as

$$W_j(s) = -U_j(s) \Phi_{\varepsilon,j}(s) + (1 - U_j(s)) n_j(s) \sum_{k \in N_j} W_k(s) \quad s \in \Gamma_j \quad (10)$$

**(replace by following, since characteristic functions are redundant)**

$$W_j(s) = -\Phi_{\varepsilon,j}(s) + n_j(s) \sum_{k \in N_j} W_k(s) \quad ; \quad s \in \Gamma_j \quad , \quad j = 1, \dots, M \quad (11)$$

where  $N_j$  is the set of all neighboring subdomains of  $\Omega_j$ ,  $W_k$  is understood as its trivial extension to  $\Omega$ , and  $n_j$  is a normalization factor. If  $s$  does not belong to any neighbor of  $\Omega_j$ ,

then  $n_j(s)$  vanishes. Otherwise,  $n_j(s)$  is the reciprocal of the number of neighbors. When we substitute the local problems (9) into the decomposition (10), and define  $(\tilde{\mathcal{S}}_{jk} \sigma_k)(s) = n_j(s) (\tilde{\mathcal{S}}_k \sigma_k)(s)$ , we obtain

$$\mathcal{S}_j \sigma_j = -U_j \Phi_{\varepsilon,j} + (1 - U_j(s)) \sum_{k \in N_j} \tilde{\mathcal{S}}_{jk} \sigma_k \quad \text{on } \Gamma_j \quad (12)$$

(replace by following, since characteristic functions are redundant)

$$\mathcal{S}_j \sigma_j = -\Phi_{\varepsilon,j} + \sum_{k \in N_j} \tilde{\mathcal{S}}_{jk} \sigma_k \quad \text{on } \Gamma_j \quad (13)$$

As opposed to the local problem (8) which features a global interaction of all subdomains, the ddCOSMO step (13) is characterized by the interaction of subdomain  $\Omega_j$  with only its neighbors. This results in a sparse, rather than dense, discrete operator.

We discretize equation (8) and (13) by expanding  $\Phi_j$ ,  $\Phi_{\varepsilon,j}$  and  $\sigma_j$  as truncated series of spherical harmonics. If  $Y_\ell^m$  indicates the spherical harmonic of degree  $\ell$  and order  $m$  on the unit sphere  $\mathbb{S}$ , we approximate the surface charge  $\sigma_j$  as

$$\sigma_j(s) = \sigma_j(x_j + r_j y) = \frac{1}{r_j} \sum_{\ell=0}^{L_{\max}} \sum_{m=-\ell}^{\ell} [X_j]_\ell^m Y_\ell^m(y)$$

for some unknown coefficients  $X = [X_j]_\ell^m$  and a prescribed integer parameter  $L_{\max}$ . Here  $y$  is the variable on  $\mathbb{S}$ . We approximate  $\Phi_{\varepsilon,j}$  and  $\Phi_j$  in the same fashion, namely

$$\Phi_{\varepsilon,j} = - \sum_{\ell=0}^{L_{\max}} \sum_{m=-\ell}^{\ell} [G_j]_\ell^m Y_\ell^m \quad , \quad \Phi_j = - \sum_{\ell=0}^{L_{\max}} \sum_{m=-\ell}^{\ell} [F_j]_\ell^m Y_\ell^m$$

where  $G = [G_j]_\ell^m$  and  $F = [F_j]_\ell^m$  are the coefficients of the expansions, and the minus signs have been introduced for convenience. In the following, we shall use the condensed notation  $\sum_{\ell,m}$  to indicate the double sum. We interpret each local problem (8) and (13) in a variational setting that uses spherical harmonics as test functions, see Appendix A. We employ orthogonality conditions of the spherical harmonics, along with Lebedev grids to perform numerical quadrature to derive discretizations of the global problems (4) and (5). Respectively, we obtain

$$A_\varepsilon G = A_\infty F \quad , \quad L X = G \quad (14)$$

and the expressions for the entries of the discrete operator  $A_\varepsilon$  are given in (A2) and (A3).

### III. COMPUTATION OF THE DDPCM-FORCES

The solvation energy can be written as a sum of subdomain contributions, which perfectly fits the domain-decomposition paradigm.

First, for a classical solute's charge distribution of the form of  $\rho = \sum_j q_j \delta_{x_j}$ , we can develop

$$E_s = \frac{1}{2} \int_{\Omega} \rho(x) W(x) dx = \frac{1}{2} \sum_j q_j W(x_j) = 2\pi \sum_j q_j [X_j]_0^0 Y_0^0 = \sqrt{\pi} \sum_j q_j [X_j]_0^0$$

This concept easily generalizes to point multipolar charge distributions. The evaluation of the energy for charge distributions is a bit more involved as it requires a three-dimensional integration and we refer to<sup>31</sup> for more details. In all cases however, the energy can be written as

$$E_s = \frac{1}{2} \sum_j \sum_{\ell, m} [\Psi_j]_{\ell}^m [X_j]_{\ell}^m =: \frac{1}{2} \langle \Psi, X \rangle$$

where the angular brackets indicate the double scalar product over  $j$  and  $\ell, m$ . For example, for the classical charge as illustrated above, we have

$$[\Psi_j]_{\ell}^m = \sqrt{\pi} q_j \delta_{\ell, 0} \delta_{m, 0}.$$

The force acting on the  $i$ -th particle can be computed as

$$\mathcal{F}_i = -\nabla_i E_s = -\frac{1}{2} \langle \Psi, \nabla_i X \rangle$$

where the gradient  $\nabla_i$  is understood with respect to  $x_i$  and  $X$  denotes the solution to the ddPCM-equations (14). Here we used the fact that the entries of the vector  $\Psi$  are independent of the atomic positions. On the other hand, since both the matrices  $A_{\varepsilon}$ ,  $A_{\infty}$  and the right-hand side  $F$  depend on the nuclear positions  $x_1, \dots, x_M$ , so does the unknown  $X$  of (14), which can compactly be written as  $A_{\varepsilon} L X = A_{\infty} F$ . The idea adjoint differentiation is to consider the adjoint problem  $(A_{\varepsilon} L)^* s = \Psi$  and compute the quantity  $\langle \Psi, \nabla_i X \rangle = \langle s, A_{\varepsilon} L \nabla_i X \rangle$  by using the definition of the adjoint matrix.

If combine equations (14), the fully discretized problem becomes  $A_{\varepsilon} L X = A_{\infty} F$ , and Leibnitz differentiation rule allows to move derivatives from  $X$  onto the other terms, namely

$$A_{\varepsilon} L \nabla_i X = \nabla_i A_{\infty} F + A_{\infty} \nabla_i F - \nabla_i A_{\varepsilon} L X - A_{\varepsilon} \nabla_i L X =: h_i$$



Thus, once the solution  $s$  of the adjoint problem and vector  $h$  have been determined, the forces can be computed as

$$\mathcal{F}_i = -\frac{1}{2} \langle s, h_i \rangle$$

The derivatives  $\nabla_i L$  of the ddCOSMO discretization were discussed in<sup>30</sup> and we now focus on the new parts due to the ddPCM-method. The quantity  $\nabla_i F$  is *a priori* nonzero since  $F_j$  is the discretization of  $\Phi_j = U_j \tilde{\Phi}$ . **Wouldn't it be simple to put the formula at the end of this section?**

In the remainder of this section we discuss the derivatives of the ddPCM matrix  $A_\varepsilon$ . Let  $\{s_n\}$  be the  $N_{\text{grid}}$  Lebedev integration points and define the following quantities

$$t_n^{jk} = \frac{|x_j + r_j s_n - x_k|}{r_k} \quad , \quad s_n^{jk} = \frac{x_j + r_j s_n - x_k}{|x_j + r_j s_n - x_k|} \quad , \quad U_j^n = U_j(x_j + r_j s_n)$$

The blocks  $A_{jk}^\varepsilon$  of the ddPCM matrix  $A_\varepsilon$ , see (A2) and (A3), have the form

$$\begin{aligned} [A_{jj}^\varepsilon]_{\ell\ell'}^{mm'} &= 2\pi \frac{\varepsilon + 1}{\varepsilon - 1} \delta_{\ell\ell'} \delta_{mm'} + \frac{2\pi}{2\ell' + 1} \sum_{n=1}^{N_{\text{grid}}} w_n U_j^n Y_\ell^m(s_n) Y_{\ell'}^{m'}(s_n) \\ [A_{jk}^\varepsilon]_{\ell\ell'}^{mm'} &= -\frac{4\pi\ell'}{2\ell' + 1} \sum_{n=1}^{N_{\text{grid}}} w_n U_j^n Y_\ell^m(s_n) (t_n^{jk})^{-(\ell'+1)} Y_{\ell'}^{m'}(s_n^{jk}) \end{aligned}$$

where  $\{w_n\}$  are the weights associated to the integration points. Since the derivatives are independent of  $\varepsilon$ , we drop the  $\varepsilon$ -dependency for ease of notation.

The case of the diagonal blocks yields

$$[\nabla_i A_{jj}]_{\ell\ell'}^{mm'} = \frac{2\pi}{2\ell' + 1} \sum_n w_n \nabla_i U_j^n Y_\ell^m(s_n) Y_{\ell'}^{m'}(s_n)$$

so that it only requires the derivatives of the characteristic function. The function  $U_j$  is, in practice, a smoothed version of the (discontinuous!) characteristic function, and is defined as

$$U_j(x_j + r_j y) = \begin{cases} 1 - f_j(y) & f_j(y) \leq 1 \\ 0 & \text{otherwise} \end{cases} \quad , \quad f_j(y) = \sum_{k \in N_j} \chi\left(\frac{|x_j + r_j y - x_k|}{r_k}\right)$$

where  $y$  varies on  $\mathbb{S}^2$  and  $\chi$  is a regularized characteristic function of  $[0, 1]$ . We conclude that  $\nabla_i U_j$  and, consequently,  $\nabla_i A_{jj}$  are *a priori* nonzero only when  $i \in N_j$  or  $i = j$ .

The case of the off-diagonal blocks, i.e.,  $j \neq k$ , is more involved since it includes the gradient of the product of three functions, namely

$$[\nabla_i A_{jk}]_{\ell\ell'}^{mm'} = -\frac{4\pi\ell'}{2\ell' + 1} \sum_n w_n Y_\ell^m(s_n) \nabla_i \left[ U_j^n (t_n^{jk})^{-(\ell'+1)} Y_{\ell'}^{m'}(s_n^{jk}) \right] \quad (15)$$

However, since  $t_n^{jk}$  and  $s_n^{jk}$  depend only upon  $x_j$  and  $x_k$ , if we assume  $i \neq j$  and  $i \neq k$ , we obtain

$$[\nabla_i A_{jk}]_{\ell\ell'}^{mm'} = -\frac{4\pi\ell'}{2\ell'+1} \sum_n w_n Y_\ell^m(s_n) \nabla U_j^n (t_n^{jk})^{-(\ell'+1)} Y_{\ell'}^{m'}(s_n^{jk}) \quad (16)$$

Thus, since  $U_j$  depends only upon  $x_i$  such that  $i \in N_j$ , we conclude that  $\nabla_i A_{jk}$  vanishes whenever  $i \neq j$  and  $i \neq k$  and  $i \notin N_j$ . In order to discuss the opposite case, i.e.,  $i = j$  or  $i = k$  or  $i \in N_j$ , notice that the events  $(i = j)$  and  $(i = k)$  are mutually exclusive, as are  $(i = j)$  and  $(i \in N_j)$ . We obtain the three subcases  $i = j$ , and  $i = k$ , and  $i \in N_j, i \neq k$ , which we address individually.

Standard differentiation implies that

$$\begin{aligned} \nabla_i \left[ U_j^n (t_n^{jk})^{-(\ell'+1)} Y_{\ell'}^{m'}(s_n^{jk}) \right] &= \nabla_i U_j^n (t_n^{jk})^{-(\ell'+1)} Y_{\ell'}^{m'}(s_n^{jk}) - \\ &+ U_j^n (\ell' + 1) (t_n^{jk})^{-(\ell'+2)} \nabla_i t_n^{jk} Y_{\ell'}^{m'}(s_n^{jk}) + U_j^n (t_n^{jk})^{-(\ell'+1)} (D_i s_n^{jk})^T \nabla_i Y_{\ell'}^{m'}(s_n^{jk}) \end{aligned} \quad (17)$$

where  $D_i$  emphasizes that the gradient of the vector quantity  $s_n^{jk}$  is indeed its Jacobian matrix and where the extra subscripts refer to the variables with respect to which differentiation is taken. We proceed to evaluate  $\nabla_i t_n^{jk}$  and  $D_i s_n^{jk}$ . When  $i = j$ , differentiation implies

$$\nabla_j t_n^{jk} = \frac{s_n^{jk}}{r_k}, \quad D_j s_n^{jk} = \frac{I - s_n^{jk} \otimes s_n^{jk}}{|x_j + r_j s_n - x_k|^3}$$

where  $I$  is the identity matrix and  $\otimes$  indicates the outer product. We remark that the Jacobian matrix  $D_j s_n^{jk}$  is symmetric, so that the transpose in (17) is redundant. Due to the particular relation between  $x_j$  and  $x_k$ , we obtain  $\nabla_j t_n^{jk} = -\nabla_k t_n^{jk}$  and  $D_j s_n^{jk} = -D_k s_n^{jk}$ . We can therefore analogously derive the case  $i = k$  and those relationships imply

$$\begin{aligned} \nabla_j \left[ U_j^n (t_n^{jk})^{-(\ell'+1)} Y_{\ell'}^{m'}(s_n^{jk}) \right] + \nabla_k \left[ U_j^n (t_n^{jk})^{-(\ell'+1)} Y_{\ell'}^{m'}(s_n^{jk}) \right] &= \\ \left[ \nabla_j U_j^n + \nabla_k U_j^n \right] (t_n^{jk})^{-(\ell'+1)} Y_{\ell'}^{m'}(s_n^{jk}) \end{aligned}$$

which provide a convenient way of evaluating  $[\nabla_k A_{jk}]_{\ell\ell'}^{mm'}$  from  $[\nabla_j A_{jk}]_{\ell\ell'}^{mm'}$ . In fact, we obtain the quasi-skew-symmetric relation

$$[\nabla_j A_{jk}]_{\ell\ell'}^{mm'} + [\nabla_k A_{jk}]_{\ell\ell'}^{mm'} = -\frac{4\pi\ell'}{2\ell'+1} \sum_n w_n Y_\ell^m(s_n) \left[ \nabla_j U_j^n + \nabla_k U_j^n \right] (t_n^{jk})^{-(\ell'+1)} Y_{\ell'}^{m'}(s_n^{jk})$$

for  $\nabla_i A_{jk}$ . Finally, the case  $i \in N_j, i \neq k$  reduces to (16).

## IV. NUMERICAL EXPERIMENTS

### A. Convergence tests

We first verify the implementation of the computation of the forces. Let us denote by  $\mathcal{F}_{i,\alpha}$  the  $\alpha$ -component of the force  $\mathcal{F}_i$  for  $\alpha = 1, 2, 3$ . Let  $e_\alpha$  be the canonical unit vectors of  $\mathbb{R}^3$ . Then, for any atomic position  $x_i$  we consider a sequence  $\delta_1, \delta_2, \dots$  and consider the approximate force obtained by finite difference

$$\mathcal{F}_{i,\alpha} \approx \mathcal{F}_{i,\alpha}(\delta_n) = \frac{E_s(x_1, \dots, x_i + \delta_n e_\alpha, \dots, x_M) - E_s(x_1, \dots, x_i, \dots, x_M)}{\delta_n},$$

where we have made explicit the dependency of the solvation energy on the nuclear positions.

In table ??, we illustrate the RMS of the approximate forces with  $\delta_n = x^{-n}$  for caffeine.

...

### B. Timings

Timings for different molecular structures depending on the number of atoms (i.e. alanine chains, hemoglobin, etc).

## V. CONCLUSIONS

## ACKNOWLEDGMENTS

P.G. and B.S. acknowledge funding from the German Academic Exchange Service (DAAD) from funds of the “Bundesministeriums für Bildung und Forschung” (BMBF) for the project Aa-Par-T (Project-ID 57317909).

### Appendix A: ddPCM discretization

The derivation ddPCM discrete operator  $A_{jk}^\varepsilon$  rests upon the fact that the spherical harmonics  $Y_\ell^m$  are eigenfunctions of the double layer operator on  $\mathbb{S}$ , i.e.,  $\mathcal{D} Y_\ell^m = -2\pi/(2\ell + 1) Y_\ell^m$ , along with the following jump relation for the double layer potential

$$\lim_{\delta \rightarrow +0} (\tilde{\mathcal{D}} Y_\ell^m)(y \pm \delta \nu) = \pm 2\pi Y_\ell^m(y) + (\mathcal{D} Y_\ell^m)(y) \quad (\text{A1})$$

where  $\nu$  denotes the outward normal at  $y \in \mathbb{S}$ . We shall employ the invariance by translation and scaling  $(\mathcal{D}_j \Phi_{\varepsilon,j})(s) = (\mathcal{D} \hat{\Phi}_{\varepsilon,j})(y)$ , where  $y = (s - x_j)/r_j$  and  $\hat{\Phi}_{\varepsilon,j}$  is defined on  $\mathbb{S}$  through the push-forward-like transformation  $\hat{\Phi}_{\varepsilon,j}(y) = \Phi_{\varepsilon,j}(s)$ . We begin by discussing the diagonal term  $2\pi f_\varepsilon \Phi_{\varepsilon,j} - U_j \mathcal{D}_j \Phi_{\varepsilon,j}$  of (8) where, for brevity, we set  $f_\varepsilon = (\varepsilon + 1)/(\varepsilon - 1)$ . As customary, in order to obtain a numerical discretization, we multiply by a test function  $\varphi$  and integrate over  $\Gamma_j$ . The change of variable  $y = (s - x_j)/r_j$ , yields an integral over  $\mathbb{S}$  which involves the hatted quantities, namely

$$\int_{\Gamma_j} (2\pi f_\varepsilon \Phi_{\varepsilon,j} - U_j \mathcal{D}_j \Phi_{\varepsilon,j}) \varphi = 2\pi f_\varepsilon r_j^2 \int_{\mathbb{S}} \hat{\Phi}_{\varepsilon,j} \hat{\varphi} - r_j^2 \int_{\mathbb{S}} \hat{U}_j \mathcal{D} \hat{\Phi}_{\varepsilon,j} \hat{\varphi}$$

We proceed to expand  $\hat{\Phi}_{\varepsilon,j}$  as a series of spherical harmonics with coefficients  $-[G_j]_{\ell'}^{m'}$ , and select as a test function  $\hat{\varphi}$  the rescaled spherical harmonic  $r_j^{-2} Y_\ell^m$ . The orthogonality of the spherical harmonics, together with the fact that they are eigenfunctions of the double layer potential, yield

$$\begin{aligned} \int_{\Gamma_j} (2\pi f_\varepsilon \Phi_{\varepsilon,j} - U_j \mathcal{D}_j \Phi_{\varepsilon,j}) \varphi &= \\ &= -2\pi f_\varepsilon \sum_{\ell', m'} [G_j]_{\ell'}^{m'} \delta_{\ell\ell'} \delta_{mm'} - \sum_{\ell', m'} [G_j]_{\ell'}^{m'} \frac{2\pi}{2\ell' + 1} \int_{\mathbb{S}} \hat{U}_j Y_{\ell'}^{m'} Y_\ell^m \end{aligned}$$

The last step to obtain the diagonal block  $A_{jj}^\varepsilon$  is to approximate the integral through a suitable quadrature formula with weights  $\{w_n\}$  and nodes  $\{s_n\}$ . Once the numerical quadrature is carried out and the spherical harmonics expansion is truncated, we derive the final expression

$$[A_{jj}^\varepsilon]_{\ell\ell'}^{mm'} = 2\pi \frac{\varepsilon + 1}{\varepsilon - 1} \delta_{\ell\ell'} \delta_{mm'} + \frac{2\pi}{2\ell' + 1} \sum_{n=1}^{N_{\text{grid}}} w_n \hat{U}_j(s_n) Y_{\ell'}^{m'}(s_n) Y_\ell^m(s_n) \quad (\text{A2})$$

The computation of the off-diagonal term  $-U_j \tilde{\mathcal{D}}_k \Phi_{\varepsilon,k}$  employs again the fact that the double layer operator is invariant under translation and scaling, i.e.  $(\tilde{\mathcal{D}}_k \Phi_{\varepsilon,k})(x) = (\tilde{\mathcal{D}} \hat{\Phi}_{\varepsilon,k})(u)$  where  $x \in \mathbb{R}^3 \setminus \bar{\Omega}_k$  and  $u = (x - x_k)/r_k$ . In particular, when  $x \in \Gamma_j$ , i.e.,  $x = s = x_j + r_j y$  for some  $y \in \mathbb{S}$ , then  $u = u(y) = (x_j + r_j y - x_k)/r_k$ . As the quantity  $U_j \tilde{\mathcal{D}}_k \Phi_{\varepsilon,k}$  is indeed well-defined on the whole  $\Gamma_j$ , we can proceed as before and obtain

$$\begin{aligned} \int_{\Gamma_j} U_j(s) (\tilde{\mathcal{D}}_k \Phi_{\varepsilon,k})(s) \varphi(s) ds &= \int_{\mathbb{S}^2} \hat{U}_j(y) (\tilde{\mathcal{D}} \hat{\Phi}_{\varepsilon,k})(u(y)) Y_\ell^m(y) dy = \\ &= - \sum_{\ell', m'} [G_k]_{\ell'}^{m'} \int_{\mathbb{S}^2} \hat{U}_j(y) (\tilde{\mathcal{D}} Y_{\ell'}^{m'})(u(y)) Y_\ell^m(y) dy \end{aligned}$$

where  $-[G_k]_{\ell'}^{m'}$  are the coefficients of the expansion of  $\hat{\Phi}_{\varepsilon,k}$  as a series of spherical harmonics. The function  $\tilde{\mathcal{D}} Y_{\ell'}^{m'}$  is harmonic on  $\mathbb{R}^3 \setminus \overline{B(0,1)}$ , so that it has to coincide with the unique harmonic extension of its boundary value. The jump relation (A1), along with the eigenfunction property, provide the boundary value

$$\lim_{\delta \rightarrow +0} (\tilde{\mathcal{D}} Y_{\ell'}^{m'})(y + \delta \nu) = 2\pi Y_{\ell'}^{m'}(y) + (\mathcal{D} Y_{\ell'}^{m'})(y) = \frac{4\pi\ell'}{2\ell' + 1} Y_{\ell'}^{m'}(y)$$

and, by elementary notions on harmonic functions, we conclude

$$(\tilde{\mathcal{D}} Y_{\ell'}^{m'})(u) = \frac{4\pi\ell'}{2\ell' + 1} \frac{1}{|u|^{\ell'+1}} Y_{\ell'}^{m'}(u/|u|)$$

After truncation the series expansion and performing numerical integration we obtain the final result

$$[A_{jk}]_{\ell\ell'}^{mm'} = -\frac{4\pi\ell'}{2\ell' + 1} \sum_{n=1}^{N_{\text{grid}}} w_n \hat{U}_j(s_n) \frac{1}{|u(s_n)|^{\ell'+1}} Y_{\ell'}^{m'}\left(\frac{u(s_n)}{|u(s_n)|}\right) Y_{\ell}^m(s_n) \quad (\text{A3})$$

This concludes the derivation of the ddPCM discretization.

## Appendix B: ddCOSMO discretization

Here, just for convenience, we repeat some ddCOSMO equations and derivations. This will not be included in the article.

**Definition of operators:** Single layer potential:

$$(\tilde{\mathcal{S}}_j \sigma^j)(x) = \int_{\Gamma_j} \frac{\sigma^j(s')}{|s' - x|} ds' \quad , \quad x \in \Omega_j$$

Single layer operator:

$$(\mathcal{S}_j \sigma^j)(s) = \int_{\Gamma_j} \frac{\sigma_j(s')}{|s' - s|} ds' \quad , \quad s \in \Gamma_j$$

**ddCOSMO-equations:** We defined

$$W_j(x) = (\tilde{\mathcal{S}}^j \sigma^j)(x) \quad , \quad x \in \Omega_j \quad ; \quad W_j(s) = (\mathcal{S}_j \sigma^j)(s) \quad , \quad s \in \Gamma_j \quad (\text{B1})$$

and the ddCOSMO-equations become

$$W_j(s) = -U_j(s) \Phi_{\varepsilon,j}(s) + (1 - U_j(s)) n_j(s) \sum_{k \in N_j} W_k(s) \quad s \in \Gamma_j \quad (\text{B2})$$

This is equivalent to

$$(\mathcal{S}_j \sigma^j)(s) = -U_j(s) \Phi_{\varepsilon,j}(s) + (1 - U_j(s)) n_j(s) \sum_{k \in N_j} (\tilde{\mathcal{S}}_k \sigma^k)(s) \quad s \in \Gamma_j \quad (\text{B3})$$

**Mapping on unit sphere:** Then, introducing  $s = x_j + r_j \hat{s}$  and  $s' = x_j + r_j \hat{s}'$ , we deduce

$$(\mathcal{S}_j \sigma^j)(s) = \int_{\Gamma_j} \frac{\sigma^j(s')}{|s' - s|} ds' = \frac{1}{r_j} \int_{\Gamma_j} \frac{\sigma^j(s')}{\left| \frac{s' - x_j}{r_j} - \frac{s - x_j}{r_j} \right|} ds' = r_j \int_{\mathbb{S}^2} \frac{\sigma^j(x_j + r_j \hat{s}')}{|\hat{s}' - \hat{s}|} d\hat{s}',$$

so that

$$(\mathcal{S}_j \sigma^j)(x_j + r_j \hat{s}) = r_j (\mathcal{S}_{\mathbb{S}^2} \sigma^j(x_j + r_j \cdot))(\hat{s}).$$

In a similar manner, for  $x = x_j + r_j \hat{x}$  (and thus  $\hat{x} \in B_1(0)$ ):

$$(\tilde{\mathcal{S}}_j \sigma^j)(x) = \int_{\Gamma_j} \frac{\sigma_j(s')}{|s' - x|} ds' = r_j \int_{\mathbb{S}^2} \frac{\sigma^j(x_j + r_j \hat{s}')}{|\hat{s}' - \hat{x}|} d\hat{s}' = r_j (\tilde{\mathcal{S}} \sigma^j(x_j + r_j \cdot))(\hat{x}) \quad , \quad x \in \Omega_j.$$

**Operators applied to spherical harmonics:** We have

$$\left( \mathcal{S}_j Y_{\ell,m} \left( \frac{-x_j}{r_j} \right) \right) (s) = r_j (\mathcal{S}_{\mathbb{S}^2} Y_{\ell,m})(\hat{s}) = \frac{4\pi r_j}{2\ell+1} Y_{\ell,m} \left( \frac{s-x_j}{r_j} \right).$$

and

$$\left( \tilde{\mathcal{S}}_j Y_{\ell,m} \left( \frac{\cdot - x_j}{r_j} \right) \right) (x) = r_j \int_{\mathbb{S}^2} \frac{Y_{\ell,m}}{|\hat{s}' - \hat{x}|} d\hat{s}' = \frac{4\pi r_j}{2\ell+1} |\hat{x}|^\ell Y_{\ell,m} \left( \frac{\hat{x}}{|\hat{x}|} \right) = \frac{4\pi r_j}{2\ell+1} \frac{|x-x_j|^\ell}{r_j^\ell} Y_{\ell,m} \left( \frac{x-x_j}{|x-x_j|} \right)$$

**Approximation:** We now approximate

$$\sigma^j(x_j + r_j \hat{s}) = \frac{1}{r_j} \sum_{\ell,m} [X_j]_\ell^m Y_{\ell,m}(\hat{s}).$$

Observe the scaling by  $1/r_j$ . That implies that  $W_j$  at any point  $x \in \overline{\Omega_j}$  is given by

$$W_j(x) = (\tilde{\mathcal{S}}^j \sigma^j)(x) = \frac{1}{r_j} \sum_{\ell,m} [X_j]_\ell^m \left( \tilde{\mathcal{S}}_j Y_{\ell,m} \left( \frac{\cdot - x_j}{r_j} \right) \right) (x) = \sum_{\ell,m} [X_j]_\ell^m \frac{4\pi}{2\ell+1} \frac{|x-x_j|^\ell}{r_j^\ell} Y_{\ell,m} \left( \frac{x-x_j}{|x-x_j|} \right)$$

We will now multiply (B3) by  $Y_{\ell',m'}$  and integrate over  $\Gamma_j$ . We start with the left hand side:

$$\begin{aligned} \int_{\Gamma_j} (\mathcal{S}_j \sigma^j)(s) Y_{\ell',m'} \left( \frac{s-x_j}{r_j} \right) ds &= \frac{1}{r_j} \sum_{\ell,m} [X_j]_\ell^m \int_{\Gamma_j} \left( \tilde{\mathcal{S}}_j Y_{\ell,m} \left( \frac{\cdot - x_j}{r_j} \right) \right) (s) Y_{\ell',m'} \left( \frac{s-x_j}{r_j} \right) ds \\ &= \sum_{\ell,m} \frac{4\pi}{2\ell+1} [X_j]_\ell^m \int_{\Gamma_j} Y_{\ell,m} \left( \frac{s-x_j}{r_j} \right) Y_{\ell',m'} \left( \frac{s-x_j}{r_j} \right) ds \\ &= \sum_{\ell,m} \frac{4\pi r_j^2}{2\ell+1} [X_j]_\ell^m \int_{\mathbb{S}^2} Y_{\ell,m}(\hat{s}) Y_{\ell',m'}(\hat{s}) d\hat{s} = \frac{4\pi r_j^2}{2\ell'+1} [X_j]_{\ell'}^{m'} \end{aligned}$$

The first term of the right hand side becomes

$$- \int_{\Gamma_j} U_j(s) \Phi_{\varepsilon,j}(s) Y_{\ell',m'} \left( \frac{s-x_j}{r_j} \right) ds = -r_j^2 \int_{\mathbb{S}^2} U_j(x_j + r_j \hat{s}) \Phi_{\varepsilon,j}(x_j + r_j \hat{s}) Y_{\ell',m'}(\hat{s}) d\hat{s}$$

The second term on the right hand side can be developed as follows

$$\begin{aligned} &\int_{\Gamma_j} (1 - U_j(s)) n_j(s) \sum_{k \in N_j} (\tilde{\mathcal{S}}_k \sigma^k)(s) Y_{\ell',m'} \left( \frac{s-x_j}{r_j} \right) ds \\ &= \sum_{k \in N_j} \frac{1}{r_k} \sum_{\ell,m} [X_k]_\ell^m \int_{\Gamma_j} (1 - U_j(s)) n_j(s) \left( \tilde{\mathcal{S}}_k Y_{\ell,m} \left( \frac{\cdot - x_k}{r_k} \right) \right) (s) Y_{\ell',m'} \left( \frac{s-x_j}{r_j} \right) ds \\ &= \sum_{k \in N_j} \sum_{\ell,m} \frac{4\pi}{2\ell+1} [X_k]_\ell^m \int_{\Gamma_j} (1 - U_j(s)) n_j(s) \frac{|s-x_k|^\ell}{r_k^\ell} Y_{\ell,m} \left( \frac{s-x_k}{|s-x_k|} \right) Y_{\ell',m'} \left( \frac{s-x_j}{r_j} \right) ds \\ &= r_j^2 \sum_{k \in N_j} \sum_{\ell,m} \frac{4\pi}{2\ell+1} [X_k]_\ell^m \int_{\mathbb{S}^2} (1 - U_j(x_j + r_j \hat{s})) n_j(x_j + r_j \hat{s}) \frac{|x_j + r_j \hat{s} - x_k|^\ell}{r_k^\ell} Y_{\ell,m} \left( \frac{x_j + r_j \hat{s} - x_k}{|x_j + r_j \hat{s} - x_k|} \right) Y_{\ell',m'}(\hat{s}) d\hat{s} \end{aligned}$$

Combining all three results yields then the discrete ddCOSMO-equations.

## REFERENCES

- <sup>1</sup>J. Tomasi, B. Mennucci, and R. Cammi, Chem. Rev. **105**, 2999 (2005).
- <sup>2</sup>C. Cramer and D. Truhlar, Chem Rev **99**, 2161 (1999).
- <sup>3</sup>M. Orozco and F. J. Luque, Chem. Rev. **100**, 4187 (2000).
- <sup>4</sup>A. Klamt, WIREs Comput. Mol. Sci. **1**, 699 (2011).
- <sup>5</sup>B. Mennucci, WIREs Comput. Mol. Sci. **2**, 386 (2012).
- <sup>6</sup>B. Honig and A. Nicholls, Science **268**, 1144 (1995).
- <sup>7</sup>B. Roux and T. Simonson, Biophys. Chem. **78**, 1 (1999).
- <sup>8</sup>A. Klamt and G. Schuurmann, J. Chem. Soc., Perkin Trans. 2 , 799 (1993).
- <sup>9</sup>D. Chipman, J. Chem. Phys. **110**, 8012 (1999).
- <sup>10</sup>D. M. Chipman, J. Chem. Phys. **124**, 224111 (2006).
- <sup>11</sup>E. Cancès, B. Mennucci, and J. Tomasi, J. Chem. Phys. **107**, 3032 (1997).
- <sup>12</sup>E. Cancès and B. Mennucci, J.Math. Chem. **23**, 309 (1998).
- <sup>13</sup>B. Mennucci, E. Cancès, and J. Tomasi, J. Phys. Chem. B **101**, 10506 (1997).
- <sup>14</sup>V. Barone and M. Cossi, J. Phys. Chem. A **102**, 1995 (1998).
- <sup>15</sup>M. Cossi, N. Rega, G. Scalmani, and V. Barone, J. Comput. Chem. **24**, 669 (2003).
- <sup>16</sup>F. Lipparini, G. Scalmani, B. Mennucci, E. Cancès, M. Caricato, and M. J. Frisch, J. Chem. Phys. **133**, 014106 (2010).
- <sup>17</sup>F. Lipparini, G. Scalmani, B. Mennucci, and M. J. Frisch, J. Chem. Theory Comput. **7**, 610 (2011).
- <sup>18</sup>A. V. Marenich, C. J. Cramer, and D. G. Truhlar, J. Phys. Chem. B **113**, 6378 (2009).
- <sup>19</sup>B. Mennucci, J. Phys. Chem. Lett. **1**, 1666 (2010), and references therein.
- <sup>20</sup>S. Miertus, E. Scrocco, and J. Tomasi, Chem. Phys. **55**, 117 (1981).
- <sup>21</sup>J. Tomasi and M. Persico, Chem. Rev. **94**, 2027 (1994).
- <sup>22</sup>G. Scalmani and M. J. Frisch, J. Chem. Phys. **132**, 114110 (2010).
- <sup>23</sup>D. York and M. Karplus, J. Phys. Chem. A **103**, 11060 (1999).
- <sup>24</sup>A. W. Lange and J. M. Herbert, J. Chem. Phys. **133**, 244111 (2010).
- <sup>25</sup>R. Cammi and J. Tomasi, J. Comput. Chem. **16**, 1449 (1995).
- <sup>26</sup>G. Scalmani, V. Barone, K. Kudin, C. Pomelli, G. Scuseria, and M. Frisch, Theor. Chem. Acc. **111**, 90 (2004).
- <sup>27</sup>L. Greengard and V. Rokhlin, J. Comput. Phys. **73**, 325 (1987).



- <sup>28</sup>F. Lipparini, L. Lagardère, G. Scalmani, B. Stamm, E. Cancès, Y. Maday, J.-P. Piquemal, M. J. Frisch, and B. Mennucci, *J. Phys. Chem. Lett.* **5**, 953 (2014).
- <sup>29</sup>E. Cancès, Y. Maday, and B. Stamm, *J. Chem. Phys.* **139**, 054111 (2013).
- <sup>30</sup>F. Lipparini, B. Stamm, E. Cancès, Y. Maday, and B. Mennucci, *J. Chem. Theory Comput.* **9**, 3637 (2013).
- <sup>31</sup>F. Lipparini, G. Scalmani, L. Lagardère, B. Stamm, E. Cancès, Y. Maday, J.-P. Piquemal, M. J. Frisch, and B. Mennucci, *J. Chem. Phys.* **141**, 184108 (2014).
- <sup>32</sup>B. Stamm, E. Cancès, and Y. Maday, *J. Chem. Phys.* **144**, 054101 (2016).
- <sup>33</sup>C. Quan, B. Stamm, and Y. Maday, “Polarizable Continuum Model based on the Solvent Excluded Surface,” (2017), working paper or preprint.
- <sup>34</sup>M. Bugeanu, R. Di Remigio, K. Mozgawa, S. S. Reine, H. Harbrecht, and L. Frediani, *Phys. Chem. Chem. Phys.* **17**, 31566 (2015).
- <sup>35</sup>H. Harbrecht and M. Randrianarivony, *Computing* **92**, 335 (2011).
- <sup>36</sup>V. Weijs, M. Randrianarivony, H. Harbrecht, and L. Frediani, *Journal of Computational Chemistry* **31**, 1469 (2010).
- <sup>37</sup>W. Hackbusch, *Integral Equations—Theory and Numerical Treatment* (Birkhäuser, Basel, 1995).
- <sup>38</sup>E. Cancès, in *Continuum Solvation Models in Chemical Physics*, edited by B. Mennucci and R. Cammi (Wiley, New York, 2007) Chap. 1.2, pp. 29–48.