

Multiresolution quantum chemistry: Basic theory and initial applications

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We describe a multiresolution solver for the all-electron local density approximation Kohn-Sham equations for general polyatomic molecules. The resulting solutions are obtained to a user-specified precision and the computational cost of applying all operators scales linearly with the number of parameters. The construction and use of separated forms for operators (here, the Green's functions for the Poisson and bound-state Helmholtz equations) enable practical computation in three and higher dimensions. Initial applications include the alkali-earth atoms down to strontium and the water and benzene molecules. © 2004 American Institute of Physics. [DOI: 10.1063/1.1791051]

I. INTRODUCTION

We describe a multiresolution solver for the all-electron local density approximation (LDA) Kohn-Sham¹ equations for molecules. The objective of this work is to provide a practical approach for computation on general polyatomic systems without basis set error and with the computational cost of applying all operators scaling linearly with the number of parameters. Besides using a multiresolution approach, a critical step in attaining this objective has been the development of separable representations for kernels of Green's functions.² We describe a prototype, orbital-based implementation with test application to closed-shell systems including the alkali-earth atoms down to strontium, and the water and benzene molecules. We consider that a demonstration of a practical approach for solution of one-electron methods is an essential precursor to direct numerical solution of two- and many-electron problems.

With only a few notable exceptions,^{3,4} mainstream molecular quantum chemistry is performed with the linear combination of atomic orbitals (LCAO) approximation most commonly using atom-centered Gaussian functions. These methods date back 50 years to the work of Roothaan,^{5,6} Hall,⁷ and Boys,⁸ and many of the successes of modern

where s_l is a scalar, $s_1 \geq \text{fl} \geq s_M > 0$, and $K_i^l(x_{j_i}, x_{j'_i})$ are kernels of normalized (norm one) operators T_i^l in dimension $d = 1$. We require the error to be less than ϵ , namely,

$$\| T - \sum_{l=1}^M s_l T_1^l \otimes T_2^l \otimes \text{fl} \otimes T_d^l \| < \epsilon. \quad (2)$$

We call the scalars s_l separation values and the integer M the

We note that for isotropic kernels, it is sufficient to compute the transition matrix elements for positive translations only. Values for negative translations are obtained using symmetries of the basis functions $\phi_i(x) = (-1)^i \phi_i(1-x)$. The one-dimensional (1D) transition matrix elements are computed and a singular value decomposition is performed for these matrices. Away from the singularity, the matrices have very low operator rank.

When applying the operator, a test is performed to see if it is more efficient to use a low-operator-rank form. In 1D, the break-even point is an operator rank less than half the matrix dimension. In 3D, if the transformations are done in order of increasing rank, the break-even ratios are about 3/4, 2/3, and 1/2 for the first, second, and third transformations, respectively.

Efficient application of the separated kernel requires screening based upon both the magnitude of the coefficients of the source function and the coefficients of the nonstandard form of the three-dimensional operator. We currently estimate the l^2 norm of each block of the operator using the power method, namely, the rapidly convergent iteration $\|A\|_2 = \lim_{n \rightarrow \infty} \sqrt{x_n/x_{n-1}}$, where $x_n = A^t A x_{n-1}$. Typically, two to four iterations provide more than one digit precision starting from a random initial guess.

We then use an algorithm for constructing a nearly optimal representation using Ref. 38, where the authors extend the approach in Ref. 39. As a result, we find the generalized Gaussian nodes and weights such that

$$\frac{1}{r} \sum_{m=1}^M w_m e^{-p_m r^2} \leq \frac{\epsilon}{r}, \quad (11)$$

where ϵ is the required maximum relative error within the range and where the number of terms, M , is nearly optimal. The optimization procedure in Ref. 38 also allows us to obtain an approximation

$$\frac{1}{r} \sum_{m=1}^M w_m e^{-p_m r^2} \leq \frac{\epsilon}{r^2}, \quad (12)$$

which is sufficient to obtain Eq. 9) and requires fewer terms.

The third step in obtaining Eq. 9) is to substitute the approximation in Eq. 11) or 12) into Eq. 8) and compute one-dimensional integrals. A detailed derivation of these representations will appear separately.⁴⁰

For instance, a 52-term representation was constructed to obtain a relative precision of 10^{-8} in the interval $[10^{-5}, 1]$ using Eq. 11). The representation obtained using the trapezoidal rule in Eq. 10) used 300+ terms, which may be partially reduced through elementary methods.³⁷ The optimal representation was then formed using Ref. 38, yielding 52 terms for the relative accuracy displayed in Fig. 1.

For $r > 0$, the kernel G in Eq. 3) is not homogeneous, so that the approximation should be constructed for each scale separately which will yield a compact representation of the matrix elements [see Eq. 5)] rather than of the kernel as a whole. Although this is clearly the most efficient approach, details of which we will present in a sequel, in the prototype code we used an expansion that spans all length scales. To generate the separated representation of G , $r > 0$, we use instead of Eq. 10) the integral

$$\frac{e^{-r}}{r} = \frac{2}{\sqrt{\pi}} \int_{-\infty}^{\infty} e^{-r^2 e^{2s} - 2e^{-2s} + s} ds. \quad (13)$$

Due to the superexponential decay at both ends of the range, the trapezoidal rule may be directly applied to this integral to obtain the desired separated representation.³⁷ In our computations with G , $r > 0$, we select the range of validity to include $[0.001/(LZ_{\max}), 1.0]$, where Z_{\max} is the maximum nuclear charge and L is the simulation cell size.

III. KOHN-SHAM EQUATIONS

The nonlinear Kohn-Sham equations [e.g., Eqs. 7.2.7-9

nonlocal exchange potential [Eq. 2.5.12) in Ref. 1]. Implementation of gradient-corrected functionals, Hartree–Fock exchange, and hybrid potentials will be discussed in future publications.

The eigenvalues ϵ_i are also referred to as the orbital energies and will be negative for the occupied orbitals. The asymptotic form of the LDA occupied orbitals is obtained by substituting the form $r^{-1} \exp(-\alpha r)$ into Eqs. 15) and collecting terms at long range. The result is, for a neutral system, $\alpha_i = \sqrt{-2\epsilon_i}$ and $\beta_i = 1/\alpha_i - 1$. The nonlocal exchange potential causes the asymptotic behavior of the Hartree–Fock orbitals⁴² to differ from that of the LDA orbitals, but the actual iterative solutions of the LDA and Hartree–Fock equations are very similar.

A. Integral equation formulation

In 1962, Kalos⁴³ used the following Lippmann–Schwinger integral formulation to determine via Monte Carlo sampling the ground state wave function ψ and corresponding eigenvalue E ,

$$\psi = -2G V \psi, \quad (19)$$

where G is an integral operator with the kernel being the Green’s function defined by

$$G(r, r') = \int \frac{e^{-\sqrt{-2E} |r-r'|}}{|r-r'|} \quad (20)$$

and $\psi = \sqrt{-2E}$. For one particle in three dimensions with free-space boundary conditions,

$$G(r, r') = \frac{1}{4} \int_0^\infty \frac{e^{-\sqrt{-2E} |r-r'|}}{|r-r'|} \quad (21)$$

To determine the wave function, the integral equation is iterated and the eigenvalue adjusted to conserve the norm of the wave function. In combination with deflation (Sec. VI), it may be used to extract the eigenvectors of the three-dimensional Hamiltonian. Away from the origin, the bound-state Helmholtz Green’s function is smooth and decays more rapidly than the Green’s function for the Poisson equation. It is therefore very efficiently represented in the multiwavelet basis. This integral formulation of the DFT equations is also commonly used in band structure calculations.⁴⁴ However, it is important to point out that the scattering-state (positive energy) Green’s function does not have a sparse representation in wavelet bases since the function is oscillatory at long range and its higher derivatives do not decay rapidly.

Beyond providing a simple and rapidly convergent iteration to compute the eigenfunctions, this integral equation is of interest because it does not require the use of derivative

IV. ANALYTIC ENERGY DERIVATIVES

Derivatives of the variational Hartree–Fock or DFT energy with respect to the nuclear coordinates (or any parameter in the external potential) are straightforwardly evaluated since the Hellman–Feynman theorem is obeyed up to the truncation threshold:

$$\frac{\partial E}{\partial q} = \frac{\partial V}{\partial q}$$

essary work is done processing matrices that are currently treated as dense and/or full rank. Odd orders are more efficient due to the superior error in derivatives.¹⁴

At each threshold)

The H₂ molecule (bond length 1.4 bohrs) was also used to examine the convergence behavior for iteration of the integral equation with and without use of the subspace information, and with and without the multiscale solution. All calculations started from the STO-3G orbital generated at a bond length of 2 bohrs. Simple iteration of the integral equation, updating the eigenvalue according to Eq. (30), using a multiwavelet

populations can be readily accomplished by projection into a minimum atomic basis.⁵⁸

The quantum chemical algorithms are written at a very high level in terms of operators and functions. This should be compared with the conventional approach in atomic bases which manipulates explicitly indexed, sparse, multidimensional arrays of matrix elements (one- and two-electron integrals). This has led to the very rapid and compact implementation of energies, analytic derivatives, and linear response

