A method to compute periodic sums

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Abstract

In a number of problems in computational physics, a finite sum of kernel functions centered at \( N \) particle locations located in a box in three dimensions must be extended by imposing periodic boundary conditions on box boundaries. Even though the finite sum can be efficiently computed via fast summation algorithms, such as the fast multipole method (FMM), the periodized extension is usually treated via a different algorithm, Ewald summation, accelerated via the fast Fourier transform (FFT). A different approach to compute this periodized sum just using a blackbox finite fast summation algorithm is presented in this paper. The method splits the periodized sum in to two parts. The first, comprising the contribution of all points outside a large sphere enclosing the box, and some of its neighbors, is approximated inside the box by a collection of kernel functions ("sources") placed on the surface of the sphere or using an expansion in terms of spectrally convergent local basis functions. The second part, comprising the part inside the sphere, and including the box and its immediate neighborhood, is treated via available summation algorithms. The coefficients of the sources are determined by least squares collocation of the periodicity condition of the total potential, imposed on a circumspherical surface for the box. While the method is presented in general, details are worked out for the case of evaluating electrostatic potentials and forces. Results show that when used with the FMM, the periodized sum can be computed to any specified accuracy, at a cost that is twice that of the free-space FMM with the same accuracy. Several technical details and efficient algorithms for auxiliary computations are provided, as are numerical comparisons.

Keywords

periodic sums; fast multipole method; Ewald summation; GPU computing; kernel independent methods; molecular dynamics; long-range interactions

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In a number of problems in computational physics, a finite sum of kernel functions centered at N particle locations located in a box in three dimensions must be extended by imposing periodic boundary conditions on box boundaries. Even though the finite sum can be efficiently computed via fast summation algorithms, such as the fast multipole method (FMM), the periodized extension is usually treated via a different algorithm, Ewald summation, accelerated via the fast Fourier transform (FFT). A different approach to compute this periodized sum just using a blackbox finite fast summation algorithm is presented in this paper. The method splits the periodized sum in to two parts. The first, comprising the contribution of all points outside a large sphere enclosing the box, and some of its neighbors, is approximated inside the box by a collection of kernel functions (?sources?) placed on the surface of the sphere or using an expansion in terms of spectrally convergent local basis functions. The second part comprising the part inside the sphere, and including the box and its immediate neighborhood, is treated via available summation algorithms. The coefficients of the sources are determined by least squares collocation of the periodicity condition of the total potential, imposed on a circumspherical surface for the box. While the method is presented in general, details are worked out for the case of evaluating electrostatic potentials and forces. Results show that when used with the FMM, the periodized sum can be computed to any specified accuracy, at a cost that is twice that of the free-space FMM with the same accuracy. Several technical details and efficient algorithms for auxiliary computations are provided, as are numerical comparisons.
1 Introduction

Many problems in physics, chemistry and materials science lead to a free-space finite “particle” sum of \( N \) functions, \( K \), centered at locations \( x_i \in \Omega_0 \subset \mathbb{R}^3 \), where \( \Omega_0 \) is a rectangular box \( d_1 \times d_2 \times d_3 \) centered at the origin of the reference frame

\[
\tilde{\phi}(y) = \sum_{i=1}^{N} q_i K(y - x_i).
\]  

(1)

For evaluation at \( N \) locations \( y \), this sum has a quadratic cost. There are efficient and arbitrarily accurate approximation algorithms for this summation (e.g., the fast multipole method, FMM [6]).

Often, an extension to this sum for \( \tilde{\phi} \) must be computed in which periodic boundary conditions are enforced on box boundaries, resulting in the potential \( \phi \). This can be evaluated by replacing the sum (1) with the infinite sum

\[
\phi(y) = \sum_{p \in P} \sum_{i=1}^{N} q_i K(y - x_i + p), \quad p \in P = \{(i_1d_1, i_2d_2, i_3d_3) : (i_1, i_2, i_3) \in \mathbb{Z}^3\}.
\]  

(2)

For some functions \( K \), such as those representing the field of an electrostatic charge, this infinite sum may be divergent or conditionally convergent. In this case certain side conditions may be needed to compute a physically relevant sum. Usually such infinite sums are performed using Fourier-transform based Ewald summation [1], which is accelerated via the FFT. This method is described briefly in Appendix C. Accounting for all pairwise interactions the method can achieve \( O(N \log N) \) complexity, for \( N \) particles in the box \( \Omega_0 \) which is periodically replicated over the full space [2, 3]. Because of the technique used for grid-to-particle interpolation these methods are usually low-order. A high-order accurate Gaussian interpolation based Ewald summation algorithm was recently presented in [4, 5].

A criticism of FMM algorithms has been that they are relatively harder to implement, combining the need for efficient data structures, careful analysis and computation of special functions, and mixed memory access patterns. Nevertheless, several open-source and commercial packages implementing the FMM for standard kernels in free space have become available. The FMM is not often used in practice to compute periodic sums, even though several methods to handle periodic boundary conditions using extensions to the basic FMM have been proposed, starting from the first publication of the algorithm [6, 7, 8, 9, 10, 11, 12, 13]. One issue with these extensions is that, compared to the basic FMM for the summation of free-space functions, they are more expensive and may require a more complicated algorithm (data structures and more complex functions, e.g., periodic Green’s functions). Analysis of the FMM and its plane-wave variant, which is better-suited for large \( N \) and parallel architectures than the smooth particle mesh Ewald algorithm, is presented in [14]. However, all these methods require constructing a new and different algorithm – a periodic variant, for which efficient implementations are not in general available.

Special purpose hardware such as graphics processors or heterogeneous CPU/GPU architectures also allow the fast computation of finite sums, either via brute force summation [15], or via the mapping of the FMM onto these architectures [16, 17, 18, 19]. Yokota et al. [19] favorably compare a large scale FMM-based vortex element computations with a direct numerical simulation via periodic pseudospectral methods. Their simulations could have been faster and more accurate – the FMM was executed on a finite system composed of \( 3^3 \) images, which while not being truly periodic also makes using the FMM significantly more expensive.

The problem this paper seeks to address is: Given a black-box fast summation algorithm (FSA) the user has for computing finite sums with a given kernel \( K(y - x_i) \), is it possible to compute the same sum with periodic boundary conditions without any modification of the FSA? We provide a positive answer to this
question. Our algorithm has the same cost as the FSA, and can be computed to any user specified accuracy \( \epsilon \), and does not use the FFT. The basic idea of the method is to divide the sum (2) into two parts. One part computes a finite sum of particles that lie within a sphere centered at the box. This is computed using the available FSA. The other part, is an approximation of the field within the box due to all particles outside the sphere. The field due to these sources can be represented within the box in terms of local expansions. Such local expansions have also been proposed in other attempts to extend the FMM to periodic systems, but are derived by explicit translation of multipole expansions from outside the box of interest into it. In our method, we propose to determine the coefficients directly from the periodicity conditions on the potential, which results in solution of a relatively small overdetermined function-fitting problem, easily solved via standard algorithms – e.g., rank-revealing QR decomposition. This step is in the spirit of the “kernel-independent” FMM methods [23, 24]. The computational overhead of our method is of the order of the cost of the finite FMM for the non-periodic (free space) problem with \( N \) particles in a box, and overall the periodic sum is computable for about twice the cost of the finite sum.

We present this “periodization” approach in a general setting, but focus computational examples on the evaluation of the electrostatic potential \( \phi \) and its gradient \( \nabla \phi \) at \( M \) evaluation points \( y_j \in \Omega_0 \subset \mathbb{R}^3 \) due to \( N \) charged particles of charge \( q_i \) placed in \( \Omega_0 \), and subject to periodic boundary conditions on \( \partial \Omega_0 \). This reduces to computing the infinite, conditionally convergent, sum (1), with kernel function \( K \):

\[
K (y - x) = \frac{1}{|y - x|}, \quad y \neq x; \quad K (y - x) = 0, \quad y = x, \quad \sum_{i=1}^{N} q_i = 0. \tag{3}
\]

and the net charge in each box being zero. Of course, in practical applications, such as in molecular dynamics, there will be other computations in addition to the one discussed here, to stabilize the overall computations. This paper does not consider these, focusing on the electrostatic sum at a single time step.

The periodization method could be easily applied to other kernels for which a “fast summation algorithm” (FSA) is available. As long as the periodic sum makes sense, and if the kernel \( K \) can be expanded over some local basis the method should work. Conditions similar to the charge neutrality in (3) may be necessary. The proposed method can also be applied for 2D problems, though we present it in 3D. Also, the periodic extension of the computational domain (box) may be performed only along one or two coordinates, for which Ewald summation may have problems.

2 Proposed method

2.1 Periodization

The box on which the periodic sum is to be computed is denoted \( \Omega_0 \). Let \( S_0 \) be a ball of radius \( R_0 \) centered at the center of the box \( \Omega_0 \) and containing it. Let \( S_b \) be another ball with the same center and radius \( R_b > R_0 \) (see Fig. 1). We denote as \( \Omega_b \) a finite region which includes the ball \( S_b \) (the minimal \( \Omega_b \) is the ball \( S_b \)). We decompose the infinite sum as

\[
\phi (y) = \phi_{\text{near}} (y) + \phi_{\text{far}} (y), \quad \phi_{\text{near}} (y) = \sum_{x_j \in \Omega_b} q_j K (y - x_j), \quad \phi_{\text{far}} (y) = \sum_{x_j \notin \Omega_b} q_j K (y - x_j), \tag{4}
\]

where \( \phi_{\text{near}} (y) \) is to be computed using the FSA (assumed to be the FMM in the sequel) to the specified accuracy \( \epsilon \) while \( \phi_{\text{far}} (y) \) is computed by some other method at least to the same accuracy. The sources in the infinite domain are indexed as \( x_j = x_i - p, \ q_j = q_i \) for appropriate vectors \( p \in \mathbb{P} \) (see Eq. (2)).

To apply the FMM to the kernel function \( K (y) \) it should be possible to approximate it via a convergent series over a set of local basis functions \( \{ R_t (y) \} \). This means that for any source point \( x_j \notin \Omega_b \) we have
Figure 1: The particle sum (2) in the box $\Omega_0$, subject to periodic boundary conditions on the box boundary, is to be computed to a specified accuracy $\epsilon$. The boundary condition can be enforced by the image method, which results in an infinite set of copies of the box $\Omega_0$. The proposed method divides the sum into a near-field component in $\Omega_b$ and a far-field sum.

The factorization

$$ K (y - x_j) = \sum_{t=1}^{P} B_t (x_j) R_t (y) + \epsilon_j^{(P)}, \quad |x_j| > R_b, \quad |y| < R_0, $$

where $P$ is the number of terms retained in the infinite series, $B_t (x_j)$ are the expansion coefficients, and $\epsilon_j^{(P)}$ is the truncation error depending on $|x_j|$ and $R_0$. The basis functions can be some standard choices, or, as in the kernel independent FMM [24], can be taken to be a collection of kernel functions, centered outside the region of approximation

$$ R_t (y) = K \left( y - x_i^{(s)} \right), \quad |x_i^{(s)}| > R_b, $$

where $x_i^{(s)}$ are a collection of sources located outside ball $S_b$. This method has the flavor of “equivalent-source” methods. Since the function $K$ depends only on the distance between its argument, it is a “radial basis function”, or RBF [21]. While any approximation scheme may be used, the use of kernel $K$ as RBF is dictated by the fact that this kernel satisfies the underlying equation (e.g. the Laplace equation), so the approximation to the field via the sum of such RBFs also satisfies the equation (if it is linear and space
A necessary condition for our method is convergence of both infinite sums $C_y^A$ and $C_y^L$ determined via least-squares collocation as follows. Consider a set of computation of invariant). Substituting Eq. (5) into the expression for $\phi_{far}(y)$ from Eq. (4), we obtain

$$
\phi_{far}(y) = \sum_{t=1}^{P} C_t R_t(y) + \epsilon^{(P)},
$$

(7)

$$
C_t = \sum_{x_j \in \Omega_b} q_j B_t(x_j), \quad \epsilon^{(P)} = \sum_{x_j \in \Omega_b} q_j \epsilon^{(P)}.
$$

A necessary condition for our method is convergence of both infinite sums $C_t$ and $\epsilon^{(P)}$. The problem of computation of $\phi_{far}(y)$ has been reduced to that of determination of $P$ fitting coefficients $C_t$. These can be determined via least-squares collocation as follows. Consider a set of $L > P$ check points, $Y^{(c)} \subset S_0 \setminus \Omega_0$. A point $y_i^{(c)} \in Y^{(c)}$ has two properties: first, $|y_i^{(c)}| < R_0$, and, second, there exists $p \in \mathbb{P}$ such that point $\tilde{y}_i^{(c)} = y_i^{(c)} + p \in S_0$ (see Fig. 1). This means that

$$
\phi(y_i^{(c)}) = \phi(\tilde{y}_i^{(c)}), \quad l = 1, \ldots, L.
$$

(8)

In terms of decomposition (4) and representation of the far field (7) this system can be rewritten as

$$
\sum_{t=1}^{P} A_{lt} C_t = f_l + \epsilon^{(P)}_l, \quad l = 1, \ldots, L,
$$

(9)

$$
A_{lt} = R_t(y_l^{(c)}) - R_t(\tilde{y}_l^{(c)}), \quad f_l = \phi_{near}(\tilde{y}_l^{(c)}) - \phi_{near}(y_l^{(c)})
$$

where $|\epsilon^{(P)}_l| = \max_{y \in S_0} |\epsilon^{(P)}(y)|$. We have $L$ linear equations in $P$ unknowns $C_1, \ldots, C_P$. As we are not constrained with the size of the set of the set points, $L > P$ can be selected to provide a substantial oversampling, so minimization of functional

$$
F(C_1, \ldots, C_P) = \sum_{l=1}^{L} \left( \sum_{t=1}^{P} A_{lt} C_t - f_l \right)^2,
$$

(10)

should take care about the “noise” introduced into the approximation due to $\epsilon^{(P)}$. The least square minimization procedure is well known and formally it results in solution

$$
C = A^\dagger \mathbf{f}, \quad A^\dagger = (A^T A)^{-1} A^T,
$$

(11)

where $C = \{C_t\}$ and $\mathbf{f} = \{f_l\}$ are organized as column vectors of size $P$ and $L$, respectively and $A^\dagger$ is the $P \times L$ matrix, which is the pseudoinverse of $A$, and superscript $T$ denotes transposition. Note that this notation is formal, and is not the way the least-squares problem is solved in practice. Rather a stable algorithm such as the rank-revealing QR decomposition [28] is used.

The known coefficients $C$ allow computation of $\phi_{far}(y)$ and can be added to the $\phi_{near}$ obtained via the FSA. However, some technical details need to be specified. In the next section we provide analysis and details for the important case of the Coulombic kernel (2).

A similar collocation of kernel based RBF expansions at a relatively small amount of the check points is also used in the “kernel independent” FMM [23] with basis functions (6). There, the collocation is at the level of the boxes in the FMM octree data structure, and fitting takes the place of expansions and translations. Here, we collocate the differences of the overall solution at a set of check points $y_l^{(c)}$ and at their periodic images $\tilde{y}_l^{(c)}$, at the level of the overall domain to determine the expansion coefficients for $\phi_{far}$. 

5
Figure 2: The check point distributions over the surface of a unit sphere for \( p = 16 \): random distribution \((L = 3p^2\) points), the Gauss spherical grid \((L = 2p^2 - p)\), and the Thomson points \((L = p^2)\).

2.2 Check point set

Selection of an optimal set of check points \( T^{(c)} \subset S_0 \setminus \Omega_0 \) is not trivial. A good check point set should yield a well conditioned solution, and sample the solution well spatially. A simple way which does not yield such a well conditioned set is to select the points \( y_l^{(c)} \subset T^{(c)} \) on the boundary of box \( \Omega_0 \). This is because the corresponding periodic point \( \tilde{y}_l^{(c)} \) will then also be located on the box boundary on the face opposite to \( y_l^{(c)} \). In this case the distances from the box center to \( y_l^{(c)} \) and to \( \tilde{y}_l^{(c)} \) will be the same, so both points will be located on a sphere of radius \( r_l^{(c)} \). If the basis is based on spherical functions, then several of these will take the same value for symmetrical points on the sphere, and the fitting equations may be rank deficient.

To avoid this degeneracy, the set \( T^{(c)} \) was chosen from points on the surface of ball \( S_0 \). Three distributions were tried (see Fig. 2): i) random uniformly distributed points, ii) Gaussian nodes (zeros of the Legendre polynomial along \( \theta \), \( P_p(\cos \theta) \), see [29]) and equispaced with respect to \( \varphi \), and, iii) the almost uniform distribution of points over the sphere obtained by solving the so-called Thomson problem of the equilibrium position of mutually repelling electrons constrained to be on the surface of the sphere [30], see [31]. Methods ii) and iii) show better results than the random distribution, as shown in Section 3. Further, we found that using the Thomson points for the interpolation in Eq. 6 provides good accuracy.

2.3 Periodization algorithm

The algorithm has two parts. In a first preliminary set-up step, denoted “set”, the check points are determined, and the matrix decompositions necessary to compute the least squares fit with the matrix \( A \) in Eq. (11) are precomputed. In simulations where the domain \( \Omega_0 \) is fixed and the particles move, as in molecular dynamics, this matrix does not change, and the cost of the “set” step is amortized over the entire simulation. The second part of the algorithm, denoted “get,” computes the right hand side and solution of the fitting equations via an inexpensive step such as backsubstitution. The accuracy depends on the choice of basis functions, and the parameters \( P, L, \) and \( R_b \). In the next section we provide both a theoretical and an empirical study for the case of the Coulombic kernel (Green’s function of Laplace’s equation).
2.3.1 Algorithm “set”

1. Set the circumsphere radius \( R_0 = \frac{1}{2} \sqrt{d_1^2 + d_2^2 + d_3^2} \). Based on the required accuracy determine \( R_0 \), \( P \), and \( L \geq P \).

2. Generate \( L \) check points distributed over the surface of the ball \( S_0 \), \( y_l^{(c)} \in \partial S_0 \), \( l = 1,...,L \). Denote this point set as \( Y_{cl} \).

3. For each point \( y_l^{(c)} = (y_{l1}^{(c)}, y_{l2}^{(c)}, y_{l3}^{(c)}) \), \( l = 1,...,L \), find a point \( \tilde{y}_l^{(c)} \in S_0 \), such that two Cartesian coordinates of \( \tilde{y}_l^{(c)} \) are the same as those of the respective coordinates of \( y_l^{(c)} \), while the other coordinate, \( y_{lk}^{(c)} \), is shifted by \( d_k \) with respect to \( y_{lk}^{(c)} \). Denote \( Y_{c2} = \{ \tilde{y}_l^{(c)} \} \).

4. Form the \( L \times P \) fitting matrix \( A = \{ A_{lt} \} \), \( A_{lt} = R_t \left( y_l^{(c)} \right) - R_t \left( \tilde{y}_l^{(c)} \right) \), \( l = 1,...,L \), \( t = 1,...,P \), where \( R_t (y) \) are the basis functions at the checkpoints.

5. Compute matrix decomposition of \( A \) necessary to solve the least squares problem.

6. (optional) Precompute other parameters which do not depend on the source distribution. If the summation needs certain auxiliary computations to ensure convergence, do those steps. For the Coulomb kernels this may involve computation of the integrals of the basis functions over the box \( \Omega_0 \).

2.3.2 Algorithm “get”

1. Periodically extend the source box \( \Omega_0 \) to cover the ball \( S_b \) of radius \( R_b \). The newly generated sources and charges will have coordinates \( X_p = \{ x_i + p \} \) for the values of the periodization vector \( p_1, ..., p_b \) from set \( \mathbb{P} \) (see Eq. (2)) and charges \( Q_p = \{ q_i \} \). Denote the set of all sources as \( X = X_0 \cup X_{p1} \cup ... \cup X_{pb} \). These constitute \( \Omega_b \).

2. Find the set of \( N_b \) sources \( X_{near} \) by removing sources from the set \( X \) that are outside the sphere \( S_b \), and satisfy \( |x_j| > R_b (X_{near} = \{ x \in X_b : |x| \leq R_b \}) \).

3. Using the FSA compute \( \phi_{near} \) for a given set of evaluation points \( Y_0 = \{ y_j \}, j = 1,...,M \) residing in \( \Omega_0 \) and belonging to sets \( Y_{c1} \) and \( Y_{c2} \), i.e. for points from the set \( Y = \{ y \in Y_0 \cup Y_{c1} \cup Y_{c2} \} \). If gradient computations are needed, compute \( \nabla \phi_{near} \) at \( y_j \in Y_0 \).

4. Form the right hand side of the periodization equation \( f = \{ f_l \}, f_l = \phi_{near} (\tilde{y}_l^{(c)}) - \phi_{near} (y_l^{(c)}) \), \( l = 1,...,L \), organized in a column vector.

5. Using the matrix decompositions in the “set” step solve the fitting equations for the \( P \) expansion coefficients \( C = (C_1, ..., C_P)^T \). This step can formally be written as \( C = A^T f \).

6. (optional) Compute the constant shift or other modification of the far field potential, if needed.

7. Evaluate \( \phi_{far} (y_j) \), and if gradient computation is needed, \( \nabla \phi_{far} (y_j) \), at \( y_j \in Y_0 \).

8. Get the periodized solution of the problem, \( \phi (y_j) = \phi_{near} (y_j) + \phi_{far} (y_j), y_j \in Y_0 \), and if gradient computation is needed, \( \nabla \phi (y_j) = \nabla \phi_{near} (y_j) + \nabla \phi_{far} (y_j), y_j \in Y_0 \).

Remark 1 In molecular dynamics and other N-body simulations the source and evaluation points are the same, so \( Y_0 = X_0 \).
Remark 2 In Step 5 for the Laplacian kernel and spherical basis functions $C = (C_2, ..., C_P)^T$. In this case optional Step 6 provides $C_1$ (see Eqs (35)-(38)). Otherwise set $C_1 = 0$.

Remark 3 Step 2 reduces $\Omega_b$ to the ball $S_b$, which is not necessary, but is efficient.

2.3.3 Complexity

Evaluation of the kernel at a single point requires $O(1)$ operations. So, $O(P)$ operations are needed to evaluate all basis functions. The complexity, $O(P)$, is also achieved when basis functions at a point can be computed recursively (as for the spherical basis functions). With $L = O(P)$ we can estimate the complexity of the “set” part of the algorithm as

$$C^{(set)} = O(1) + O(P) + O(P) + O(P) + O(P^3) + \left[ O(P^{5/2}) \right] = O(P^3),$$

(12)

where we assumed that computation of the matrix decomposition (e.g., via QR) $O(P^3)$ operations, as $P \sim L$. The term in the square brackets is the cost of the optional step for the Laplace kernel and spherical basis functions, using the method presented in Appendix A.

For the “get” part of the algorithm, assuming that $N_b = O(N)$ and $P \ll N$ we have

$$C^{(get)} = O(N) + O(N) + C^{(FSA)} + O(P) + O(P^2) + [O(N)] + O(PN) + O(N) = O(PN) + C^{(FSA)},$$

(13)

where $C^{(FSA)}$ is the cost of the finite summation algorithm and the cost of the optional step of the algorithm is put in the square brackets (see Appendix B for this step for the Laplacian kernel). This cost for the brute force summation is $O(N^2)$, while if the FMM is used as the FSA it can be estimated as follows.

In the FMM, generation of the data structure for $M \sim N$ points is $O(Nl_{max})$ which for deep trees, $l_{max} = O(\log N)$, results in formal $O(N\log N)$ complexity. However, in practice the depth of the trees in three dimensions is relatively small (e.g. $l_{max} < 10$) for sizes $N < 10^7$ and even at larger $l_{max}$ this cost is much smaller than the cost of the run part of the FMM. Moreover, the translation time in the FMM usually dominates over the time of generation and evaluation of expansions of complexity $O(P_{FMM}N)$, where $P_{FMM}$ is the size of the expansions in the FMM. For optimal $l_{max}$ the FMM using $O(P_{FMM}^2)$ methods for translations in three dimensions scale as $O(P_{FMM}^2)N$ or $O(P_{FMM}^3)$, where $\frac{1}{2} \leq \alpha \leq \frac{3}{4}$ for well studied kernels such as those for the Laplace and Helmholtz equations (in the latter case additional $\log^\beta N$, $\beta > 0$ factors appear in the algorithm complexity, which we drop in the present estimate). Optimized kernel independent FMM has complexity $O(P_{FMM}N)$. This can be summarized as

$$C^{(get)} = O(P_{FMM}^\alpha N) + O(PN), \quad \frac{1}{2} \leq \alpha \leq 1.$$  

(14)

Note that for the Laplacian kernel in three dimensions truncation numbers $p = P^{1/2}$ and $P_{FMM} = P_{FMM}^{1/2}$ should increase as $O(\log N)$ for a fixed absolute $L_\infty$-norm error (see error bounds below), while they are constant for the relative $L_2$-norm errors.

3 Laplacian kernel

A fundamental feature of the Laplace equation that any constant is a solution. Thus, one of the basis functions is a constant (say, $R_1(y) \equiv 1$). Moreover, any constant satisfies periodic boundary conditions, and so this part of the solution cannot be determined, as $R_1(y)$ belongs to the null-space of the Laplacian operator. System (9) also shows that $A_{l1} = R_1(y^{(e)}_l) - R_1(\tilde{y}^{(e)}_l) = 0$ for any $l$, so for any $L \geq P$ the
rank of matrix $A$ cannot exceed $L - 1$, in which case the coefficient $C_1$ can be arbitrary. To remove this rank deficiency of $A$ we can simply remove the constant basis function from consideration, formulate the problem as the problem of determination of coefficients $C_t$ for $t = 2, \ldots, L$ and then add an arbitrary constant $C_1$ to the solution. Indeed, in many cases the value of the potential is not important, as only differences and gradients determine physical quantities such as the electric field, velocities, forces, etc. For comparison with the FFT-based Poisson equation solutions however, it is desirable to obtain $C_1$, in which case an additional condition, zero period average, needs to be imposed. These will be discussed in a separate subsection. Here we just mention that many other equations have the same problem (e.g. the biharmonic equation) and the null-space of some equations, such as the Helmholtz equation, may have larger dimension, and an analysis similar to the one for the Laplace equation presented here would be needed.

### 3.1 Spherical basis functions

While the basis functions can be simply selected according to Eq. (6), we instead consider the closely related polynomial basis, in which case we can establish error bounds. In spherical coordinates $(r, \theta, \varphi)$ related to the Cartesian coordinates via

$$ x = r \sin \theta \cos \varphi, \quad y = r \sin \theta \sin \varphi, \quad z = r \cos \theta, $$

the local and multipole solutions of the Laplace equation in 3D can be represented as

$$ R_n^m (r) = \alpha_n^m r^n Y_n^m (\theta, \varphi), \quad S_n^m (r) = \beta_n^m r^{-n-1} Y_n^m (\theta, \varphi), \quad n = 0, 1, \ldots, \quad m = -n, \ldots, n. \quad (16) $$

Here $R_n^m (r)$ are the regular (local) spherical basis functions and $S_n^m (r)$ the singular (or multipole) spherical basis functions; $\alpha_n^m$ and $\beta_n^m$ are normalization constants which can be selected by convenience, and $Y_n^m (\theta, \varphi)$ are the orthonormal spherical harmonics:

$$ Y_n^m (\theta, \varphi) = N_n^m P_n^{|m|} (\mu) e^{im\varphi}, \quad \mu = \cos \theta, $$

$$ N_n^m = (-1)^m \sqrt{\frac{2n+1}{4\pi} \frac{(n-|m|)!}{(n+|m|)!}}, \quad n = 0, 1, 2, \ldots, \quad m = -n, \ldots, n, $$

where $P_n^{|m|} (\mu)$ are the associated Legendre functions [29]. We will use the definition of the associated Legendre function $P_n^{|m|} (\mu)$ that is consistent with the value on the cut $(-1, 1)$ of the hypergeometric function $P_n^m (z)$ (see Abramowitz & Stegun, [29]). These functions can be obtained from the Legendre polynomials $P_n (\mu)$ via the Rodrigues’ formula

$$ P_n (\mu) = (-1)^m (1 - \mu^2)^{m/2} \frac{d^m}{d\mu^m} P_n (\mu), \quad P_n (\mu) = \frac{1}{2^n n!} \frac{d^n}{d\mu^n} (\mu^2 - 1)^n. \quad (18) $$

Straightforward computation of these basis functions involves several relatively costly operations with special functions and use of the spherical coordinates. Further, as defined above, these functions are complex, which is an unnecessary expense for real valued computations. In [16] real basis functions were defined as

$$ \tilde{R}_n^m = \begin{cases} \Re \{ R_n^m \}, & m \geq 0 \\ -\Im \{ R_n^m \}, & m < 0 \end{cases}, \quad \tilde{S}_n^m = \begin{cases} \Re \{ S_n^m \}, & m \geq 0 \\ -\Im \{ S_n^m \}, & m < 0 \end{cases}, $$

with $R_n^m$ and $S_n^m$ defined via Eq. (16) are

$$ \alpha_n^m = (-1)^n \sqrt{\frac{4\pi}{(2n+1) (n-m)! (n+m)!}}, \quad \beta_n^m = \sqrt{\frac{4\pi (n-m)! (n+m)!}{2n+1}}, $$

$$ n = 0, 1, \ldots, \quad m = -n, \ldots, n. $$
Only local basis functions are needed here, and can be computed via an efficient recursive process, without spherical coordinates,

\[ \tilde{R}_0 = 1, \quad \tilde{R}_1 = -\frac{1}{2}x, \quad \tilde{R}_1 = \frac{1}{2}y, \]

\[ \tilde{R}_m = -z\tilde{R}_m, \quad m = 0, \pm 1, \ldots, \]

\[ \tilde{R}_m = -\frac{(2n-1)z\tilde{R}_m + r^2\tilde{R}_m}{(n - |m|)(n + |m|)}, \quad n = |m| + 2, \ldots, \quad m = -n, \ldots, n. \]

This basis can be obtained in the same manner as \( \{ R_n^m \} \) was from the complex basis \( R_n^m \), where \( \alpha_n^m = (-1)^n \sqrt{4\pi/(2n+1)} \). Complex basis functions \( R_n^m \) with a similar normalization \( \alpha_n^m \) (without the factor \( (-1)^n \)) were used in [33]. Note further that \( p \)-truncated expansion of a harmonic function \( \phi_{far} \) over basis (22) can be written as

\[ \phi_{far}(y) = \sum_{n=0}^{p-1} \sum_{m=-n}^{n} \hat{C}_n^m \tilde{R}_n^m(y) = \sum_{t=1}^{P} C_t R_t(y), \quad C_t = \hat{C}_n^m, \quad R_t(y) = \tilde{R}_n^m(y), \quad P = p^2, \]

\[ t = (n+1)^2 - (n-m), \quad n = 0, 1, \ldots, p-1, \quad m = -n, \ldots, n. \]

The latter form of the sum, where stacking of coefficients is used, is consistent with (7). The gradient of the potential is needed to compute the force. This can be computed as

\[ \nabla \phi_{far}(y) = \sum_{n=1}^{p-1} \sum_{m=-n}^{n} \hat{E}_n^m \tilde{R}_n^m(y) = \sum_{t=2}^{P} E_t R_t(y), \]

where \( \hat{E}_n^m \) are vectors in \( \mathbb{R}^3 \). In [35] one can find relations between the potential and gradient coefficients.

### 3.2 Error bounds

For the Laplacian kernel \( K \), Eq. (2), \( p \)-truncated expansions (5) over the basis \( \{ R_n^m \} \), Eq. (16), has a well-known error bound

\[ |\epsilon_j^{(P)}| < \frac{1}{|x_j| - R_0} \left( \frac{R_0}{|x_j|} \right)^p, \quad p = P^{1/2}. \]

The expansion error (7) due to all sources located in \( \mathbb{R}^3 \setminus \Omega_b \) then can be bounded as

\[ |\epsilon^{(P)}| < \frac{\max |q_j|}{R_b - R_0} \sum_{x_j \in \Omega_b} \left( \frac{R_0}{|x_j|} \right)^p \leq \frac{\max |q_j|}{R_b - R_0} \sum_{x_j \in \mathbb{R}^3 \setminus \Omega_b} \left( \frac{R_0}{|x_j|} \right)^p \]

\[ \leq \frac{\max |q_j|}{R_b - R_0} \int_{\mathbb{R}^3 \setminus \Omega_b} n(x) \left( \frac{R_0}{r} \right)^p \, dV, \]

\[ n(x) = \sum_{x_j \in \mathbb{R}^3 \setminus \Omega_b} \delta(x - x_j), \quad r = |x|. \]
Here we introduced the the number density $n(x)$, which for integral estimates can be replaced with a constant density $n_0 = N/V_0$, where $V_0$ is the volume of box $\Omega_0$, $V_0 = d_1d_2d_3$. In this case the integral can be evaluated as

$$
\int_{\mathbb{R}^3} n(x) \left( \frac{R_0}{r} \right)^p dV \sim 4\pi n_0 \int_{R_b}^{\infty} \left( \frac{R_0}{r} \right)^p r^2 dr = \frac{4\pi n_0 R_b^3}{p-3} \lambda^{-p}, \quad \lambda = \frac{R_b}{R_0}.
$$

Hence, we have an approximate error bound

$$
|\epsilon^{(P)}| \lesssim \frac{4\pi n_0 R_b^2}{R_b - R_0} \max |q_i| \lambda^{-p}.
$$

For a cubic domain we have $d_1 = d_2 = d_3 = d$, $R_0 = \frac{1}{2}d\sqrt{3}$, $R_b = \lambda R_0$, and we get

$$
|\epsilon^{(P)}| \lesssim \frac{3\pi N \max |q_i|}{d} \left( \frac{1}{\lambda - 1} \right) \left( p - 3 \right) \lambda^{p-3}.
$$

The actual error achieved in practice is expected to be much smaller than this estimate since it neglects cancellation effects due to the total charge neutrality. Also in the above equation one can set $d = 1$ to obtain a non-dimensional measure of the absolute error (since the Laplace equation is scale independent).

### 3.3 Optimization when using the FMM

There are three free parameters, $p$, $p_{FMM}$, and $\lambda$, which can be selected to optimize algorithm performance. Assuming that the number of charges, their intensities and distribution as well as the domain $\Omega_0$ are fixed, and computations performed with some prescribed tolerance, $\epsilon$, the optimization problem can be formulated as

$$
\epsilon_P(p, \lambda) = \epsilon, \quad \epsilon_{FMM}(p_{FMM}, \lambda) = \epsilon, \quad C^{(get)}(p, p_{FMM}, \lambda) \rightarrow \min,
$$

where $\epsilon_P$ and $\epsilon_{FMM}$ are the error bounds for the periodization and the FMM respectively, while $C^{(get)}$ is the cost of the “get” step. In practice these functions should be determined experimentally, using the qualitative theoretical estimates provided below for guidance.

We set the parameter $p_{FMM}$ by the prescribed accuracy $\epsilon$, and approximate $\epsilon_P(p, \lambda)$ as

$$
\epsilon_P(p, \lambda) = B_P \lambda^{-p}, \quad p = \frac{\ln (B_P/\epsilon)}{\ln \lambda},
$$

where $B_P$ is some constant. The number of evaluation points is $M_b = N + 2L$, while the number of sources to be summed are $N_b = O(\lambda^2 N)$. For a perfectly optimized FMM, in which the cost of translations and direct summations dominates the cost of other procedures, the complexity is proportional to the geometric mean of the number of sources and evaluation points. Assuming $L \sim 2p^2$ the cost can be estimated as

$$
C^{(FMM)} = O \left( \frac{2\alpha}{p_{FMM}^2 \sqrt{N_b M_b}} \right) = O \left( \frac{2\alpha}{p_{FMM}^2 N \lambda^{3/2}} \left( 1 + \frac{4p^2}{N} \right)^{1/2} \right) = A_{FMM} \left( \frac{2\alpha}{p_{FMM}^2 N \lambda^{3/2}} \left( 1 + \frac{4 \ln^2 (B_P/\epsilon)}{N \ln^2 \lambda} \right)^{1/2} \right)^{1/2}.
$$

where $A_{FMM}$ is some constant. Of course, for $N \gg p^2$ the second term in the parentheses can be dropped. However, for $\lambda \rightarrow 1$, we may have $p^2 \sim N$, and as soon as the complexity is a product of the decreasing and increasing functions, some minimum is expected. Our tests show that the cost of the FMM is the major
contributor to the overall cost of the algorithm. Taking the derivative of $C^{(FMM)}$ with respect to $\lambda$ and setting it to zero, we obtain $\lambda_{opt}$ from the solution of the cubic equation,

$$2x^3 - x^2 - A = 0, \quad x = \frac{1}{\ln \lambda_{opt}}, \quad A = \frac{N}{4 \ln^2 (B_P/\epsilon)}.$$  

(33)

At very large $N$ ($A \gg 1$) we have

$$\lambda_{opt} \sim \exp \left[ \left( \frac{A}{2} \right)^{-1/3} \right] = \exp \left[ \left( \frac{N}{8 \ln^2 (B_P/\epsilon)} \right)^{-1/3} \right],$$

(34)

which shows that at large $N$ optimal $\lambda$ should be shifted towards the limit $\lambda = 1$.

### 3.4 Constant shift in potential

Several methods can be proposed to determine coefficient $C_1$ if it is needed. In particular because we choose to compare our results with the Ewald summation method, we need it. Of course, the simplest case is that when the potential value, $\phi_0$, is prescribed or known at some point $y_0$, in which case

$$C_1 = \phi_0 - \phi_{near}(y_0) - \sum_{t=2}^{P} C_t R_t(y_0).$$

(35)

Note then that the Fourier based methods for periodization of Green’s function produce solution with some mean of the potential $\phi_{mean}$ (since the zero mode of the Fourier transform is zeroed). This particular solution corresponds to

$$\langle \phi \rangle_{\Omega_0} = \frac{1}{V_0} \int_{\Omega_0} \phi(y) \, dV(y) = \phi_{mean}. \quad (36)$$

In this case for consistency we should set

$$C_1 = \phi_{mean} - \frac{1}{V_0} \int_{\Omega_0} \phi_{near}(y) \, dV(y) - \sum_{t=2}^{P} C_t R_t^{(0)}, \quad R_t^{(0)} = \frac{1}{V_0} \int_{\Omega_0} R_t(y) \, dV(y). \quad (37)$$

As shown in Appendix C, the Ewald summation produces $\phi_{mean} = 0$, and we do the same for our method. Integrals $R_t^{(0)}$ can be computed relatively easy, since $R_t(y)$ are polynomials in the Cartesian coordinates of $y$ of degree which does not exceed $p - 1$, for which case exact quadratures exist. In fact, for a given box size ratio (e.g. for cube) these can be precomputed, scaled and used independently of particular source distribution (see Appendix A). The first integral can be represented as a sum

$$\frac{1}{V_0} \int_{\Omega_0} \phi_{near}(y) \, dV(y) = \sum_{x_j \in \Omega_b} q_j \Phi_0(x_j), \quad \Phi_0(x) = \frac{1}{V_0} \int_{\Omega_0} K(y - x) \, dV(y). \quad (38)$$

In Appendix B we provide analytical expressions for functions $\Phi_0(x)$. Despite their unwieldiness, their computation for a given $x$ is $O(1)$. Overall it is a $O(N)$ procedure to compute the sum and constant $C_1$, which is consistent with the overall complexity of the method. There also exist symmetries for periodic location of sources, which can be used to accelerate these computations, if this becomes an issue. Note also that results of Appendix B can be applied for computation of integrals representing the far field (37) in the case when the RBF, Eq. (6), is used.
4 Numerical tests

To check the accuracy and performance of the method we conducted several numerical tests. There are very few known analytical solutions, so for comparison we also implemented and tested a simple version of the Ewald summation method as an alternative method (see Appendix C).

4.1 Small size tests

As validation, we performed tests with different number of sources in the box. First, we conducted a small size test, with a cubic domain $\Omega_0$ and eight sources of charges $q_i = \pm 1$ located at the vertices, so that neighboring sources have opposite charges and the infinite domain forms a regular equispaced grid. Physically this corresponds to crystal structures, such as formed by molecules NaCl. As the reference for accuracy tests we computed the Madelung constant for this crystal [34],

$$M_{a_{Na}} = -M_{a_{Cl}} = \phi(y_{Na}) = R_{NaCl} \sum_i^{N} q_i K (y_{Na} - x_i + p)$$

where $y_{Na}$ is the location of Na atom and $R_{NaCl}$ is the distance between the closest neighbor atoms (in the tests we used $d_1 = d_2 = d_3 = 1$, in which case $R_{NaCl} = 0.5$). We also computed this constant using the Ewald summation and compared spatial distributions of the potential. As the measures of the relative errors we used

$$\epsilon_M = \left| \frac{M_{a^{(comp)}}}{M_{a^{(true)}}} - 1 \right|, \quad \epsilon_2 = \left\| \phi^{(Present)} - \phi^{(Ewald)} \right\|_2, \quad \| \phi \|_2 = \sqrt{\frac{1}{M} \sum_{i=1}^{M} \phi^2 (y_i)},$$

where $y_i \in \Omega_0$ are the receivers located on the grid used for the Ewald summation.

For the high accuracy test, we selected a $44 \times 44 \times 44$ grid, $\xi = 12$, and the sampling neighborhood for each source $N_x = 20$ for the Ewald method (see Appendix C). This setting provides $\epsilon_M \approx 10^{-14}$ (i.e. 14 digits of the Madelung constant). High accuracy test for the present method was performed with $p = 35$, $R_b = 1.5$ ($\lambda = R_b/R_0 = \sqrt{3}$), which results in errors $\epsilon_M \approx 6 \cdot 10^{-14}$ and $\epsilon_2 \approx 9 \cdot 10^{-13}$. For the middle accuracy test we used $24 \times 24 \times 24$ grid, $\xi = 10$, and the sampling neighborhood for each source $N_x = 10$, in which case the Ewald method results in $\epsilon_M \approx 6 \cdot 10^{-9}$. In our method we used $p = 16$, $R_b = 1.5$, which produced errors $\epsilon_M \approx 7 \cdot 10^{-8}$ and $\epsilon_2 \approx 10^{-6}$. These tests show that errors $\epsilon_M$ and $\epsilon_2$ are related and the former one approximately one order of magnitude smaller than the latter. So in the following accuracy tests we measured only $\epsilon_M$ for our method, which is independent of the Ewald summation routine. These computations were performed for the check points $y_i^{(c)}$ distributed on the Gauss spherical grid.

Figure 3 shows the dependence of $\epsilon_M$ computed for 651 values of parameters $\lambda = R_b / R_0$ and $p$ controlling the accuracy. The chart on the right shows that the computational errors are consistent with theoretical error bound $\epsilon_{th} = C_\epsilon (R_0 / R_b)^p$. For very small values of $\epsilon_{th}$ the computational errors are affected by the double precision roundoff errors. This shows that the parameters can be set to achieve the required accuracy.

Table 1 shows some results of the tests with different distributions of the check points. Here for the case of random distributions for any set size we performed 100 runs and the maximum error is reported. It is seen that the lowest errors were achieved using the Thomson point distributions. The number of such points
Figure 3: The relative error, $\epsilon_M$, in computations of the Madelung constant for NaCl crystal using the present method (colors). The chart on the graph compares the theoretical computational error, $\epsilon_{th} = C \epsilon (R_0/R_b)^p$ with the actual error for all data points used to plot the chart on the left. Ideally, the graph on the right should be a straight line (shown). Constant $C$ was set as $\epsilon_M/\epsilon_{th}$ for $p = 10$ and $\lambda = R_b/R_0 = 1.5$.

Table 1: Error $\epsilon_M$ for different check point distributions at $\lambda = R_b/R_0 = \sqrt{3}$.

| $p$ | Gauss sph. grid | T(256) | T(400) | Rand.$L = 2p^2$ | Rand.$L = 3p^2$ | Rand.$L = 5p^2$ |
|-----|-----------------|--------|--------|-----------------|-----------------|-----------------|
| 8   | 4.85(-6)        | 3.87(-6) | 3.90(-6) | 3.39(-5)        | 2.21(-5)        | 1.92(-5)        |
| 12  | 1.17(-7)        | 4.60(-8) | 6.28(-8) | 1.53(-6)        | 1.04(-6)        | 5.10(-7)        |
| 16  | 7.16(-8)        | 3.43(-7) | 2.74(-8) | 6.26(-7)        | 3.50(-7)        | 1.85(-7)        |

should be not less than $p^2 - 1$, which is necessary (but not sufficient) condition for the use of the present method. When the number of check points approaches $p^2 - 1$ the accuracy of the method deteriorates. Conclusion here is that if a database of the Thomson points or some analogous method of deterministic uniform distribution of the check points exist, then that method is recommended. In fact, the Gauss spherical grid also provides good results (the order of the error is the same). This grid is easy to generate for any $p$, and that is why this was used in the tests. The error for random distributions is about one order of magnitude larger than that for the Gauss spherical grid or for the Thomson points. It slowly decays with the growing oversampling. Perhaps, there is no reason to use random sets, which anyway show strong dependence of the error on $p$ and also can be used if needed.

We also performed the accuracy test for different basis functions (RBF, Eq. (6)), where 256 Thomson sampling sources $x_i^{(s)}$ were located on ball $S_b$, while the check points were the same as for the last line of Table 1 ($p = 16$). In this case we obtained $\epsilon_M \approx 1.32 \cdot 10^{-7}$, which is approximately two times larger than the error when using the spherical basis functions.

### 4.2 Large scale tests

The large scale tests were conducted for systems with $N$ up to $2^{23} \sim 10^7$, for which $O(N \log N)$ summation algorithms are needed. The main purpose of these tests is to check the performance and scaling of the present
algorithm. The reported wall clock times were measured on an Intel Xeon (2.8 GHz) 4 core PC with 8 GB RAM and averaged over ten runs of the same case.

We used a well-tested standard version of the FMM for the Laplacian kernel in three dimensions, where all translations are performed with $O(p_{FMM}^3)$ complexity using the rotation-coaxial translation-back rotation (RCR) algorithm. The code implements a standard multipole-to-local translation stencil with the maximum 189 neighborhood (see details in [36]). The code was parallelized for 4 core CPU machine using OpenMP with parallelization efficiency close to 100%. While faster versions of the FMM were available to the researchers (say utilizing graphics processors, [16]), the version used for the tests was selected to provide consistent scaling of different algorithm parts, as the periodization algorithm was implemented on multicore CPUs. For the tests we treated the FMM as a black box FSA and used it “as is” without any modifications.

First we ran accuracy tests, when the charges have random intensity, $\pm 1$, and zero total sum and are located inside a unit cube at regularly spaced grid points (subgrids of $60 \times 60 \times 60$ grid). The potential at charge locations (reference solution) was computed with high accuracy using the Ewald method (grid $60 \times 60 \times 60$, $\xi = 15$, $N_r = 24$). The present method with $p = 40$ and $p_{FMM} = 30$ for this case showed error $\epsilon_2 \approx 1.2 \cdot 10^{-10}$. Further performance tests were conducted with lower tolerance to ensure that the error of the reference solution does not affect error and optimization studies. In all cases the Gauss spherical grid ($L = 2p^2 - p$) was used for the check points.

Figure 4: Comparison of periodic solutions at the median plane $z = 0$, obtained by different methods for $30 \times 30 \times 30$ sources of random intensity $q_i = \pm 1$ in a unit box. The top row shows respectively the periodic solution obtain by the Ewald method, the non-periodic solution, and the difference between the former and latter potentials. The bottom row shows the periodic solution obtained by the method proposed in this paper, and its near and far field components for $\lambda = 1.1$. Computations performed with $p = 80$, $p_{FMM} = 16$.

Figure 4 illustrates distribution of potentials generated by 27,000 charges in a box computed by two different methods sampled on $60 \times 60 \times 60$ grid. It is seen that periodic solutions obtained using the Ewald
method and present method are almost the same ($\epsilon_2 < 5 \cdot 10^{-7}$), while substantially different from the non-periodic solution (the free field generated by the same sources). This also shows that accounting for the near field (sources in $\Omega_b$) substantially improves the non-periodic solution, but the far field component is still of magnitude comparable with the near field. This far field is smooth and its addition to the near field results in an accurate periodic solution.

Table 2: Performance for different $\lambda$ for tolerance $\epsilon_{tol,2} = 5 \cdot 10^{-7}$ at $N = 27,000$, $p_{FMM} = 16$.

| $\lambda$ | $p_{\text{min}}$ | $T^{(\text{get})}$,s | $T^{(\text{get})}_{(\text{no mean})}$,s | $T^{(FMM)}$,s | $N_b$ | $N + 2L$ | $T^{(\text{set})}$,s |
|-----------|-------------------|---------------------|---------------------|---------------------|-------------------|-------------------|---------------------|
| 1.1       | 80                | 2.81                | 2.64                | 2.03                | 97,780            | 52,440            | 150                 |
| 1.2       | 44                | 1.95                | 1.73                | 1.58                | 126,744           | 34,656            | 5.31                |
| 1.3       | 34                | 2.03                | 1.78                | 1.66                | 161,072           | 31,556            | 1.27                |
| 1.5       | 24                | 2.21                | 1.77                | 1.69                | 247,742           | 29,256            | 0.29                |
| 1.7       | 19                | 2.53                | 1.98                | 1.90                | 360,640           | 28,406            | 0.13                |

As the theory predicts existence of an optimal value of parameter $\lambda$ for a given tolerance we conducted tests to determine this value experimentally. Tables 2 and 3 display the results of these tests. Here we computed potential alone (no gradient computations). The difference between the cases shown in the tables is in the number of charges (27,000 and 216,000, respectively). In these tests $p_{FMM} = 16$, which provided the relative $L_2$-norm error of the FMM itself smaller than tolerance $\epsilon_{tol,2}$. Since the truncation number $p$ changes discretely, there is the minimal integer $p = p_{\text{min}}$ at which $\epsilon_2 < \epsilon_{2,\text{tol}}$, where $\epsilon_2$ is defined by Eq. (40). This $p$ is slightly depends on $N$ and is shown in the tables. The periodization algorithm was executed with this $p$. The tables also show the number of sources, $N_b$, and the total number of evaluation points, $N + 2L$. These numbers provide data of the size of the problem solved by the FMM. It is seen that the FMM execution time is a non-monotonic function of $\lambda$, as at the increasing $\lambda$ we have increasing $N_b = O(N\lambda^3)$ and decreasing $L \sim 2p_{\text{min}}^2$. As the theoretical complexity of the optimized FMM is proportional to $\sqrt{N_b(N + 2L)}$ there can be a minimum of this function, which, indeed, is realized in experiments shown in Table 2. This is not the case for data from Table 3, but the explanation can be that optimization of the FMM (the maximum depth of the octree) changes discretely, so the performance depends on the source and receiver distribution. It is also noticeable that despite of substantial change of $N_b$ the FMM time does not change significantly. Finally, the time of the “set” part of the algorithm increases dramatically at large $p$ (as $p^4$). However, for a given $p$ and computational domain the pseudoniverse matrix can be precomputed and stored independently on the number of charges and their distribution. So this should not be considered as a limiting factor. These tests bring us to conclusion that practical optimal $\lambda$ are in the range $\sim 1.2 \sim 1.5$. Smaller or larger $\lambda$ can be also used based on particular problems and other issues (e.g. memory complexity).

Effect of the basis functions on the algorithm performance is shown in Table 4. As in the small size tests 256 Thomson points were selected as the centers of the RBFs. Comparison of the performance obtained using the RBFs vs the spherical basis functions, show that the latter choice is beneficial in terms of the accuracy, while the speed is approximately the same. Nonetheless, the errors for both bases are of the same order, while the RBF implementation is slightly simpler.

Figure 5 illustrates scaling of the algorithm with the problem size. In this test $N$ sources were distributed randomly in the unit box and both $\phi$ (with zero mean) and $\nabla \phi$ were computed at the source locations. Here only the time for the “get” part is displayed. For comparison we also plotted the wall clock time for solution of the same non-periodic problem, where the FMM with the same $p_{FMM}$ was used. (The discussion of GPU times is below). While for the non-periodic case the FMM is scaled approximately as
Table 3: Performance for different $\lambda$ for tolerance $\epsilon_{tol,2} = 5 \cdot 10^{-7}$ at $N = 216,000$, $p_{FMM} = 16$.

| $\lambda$ | $p_{\min}$ | $T^{(get)}$, $s$ | $T^{(get)}_{(no,mean)}$, $s$ | $T^{(FMM)}$, $s$ | $N_b$ | $N + 2L$ | $T^{(set)}$, $s$ |
|-----------|------------|-----------------|----------------|-------------|-------|---------|----------------|
| 1.1       | 80         | 15.7            | 14.6           | 11.5        | 782,131| 241,440| 150            |
| 1.2       | 44         | 12.6            | 11.2           | 10.2        | 1,015,037| 223,656| 5.31           |
| 1.3       | 34         | 12.8            | 11.1           | 10.3        | 1,291,247| 220,556| 1.27           |
| 1.5       | 25         | 15.0            | 12.3           | 11.7        | 1,983,665| 218,450| 0.32           |
| 1.7       | 19         | 18.7            | 14.7           | 14.1        | 2,887,393| 217,406| 0.13           |

Table 4: Performance for different basis functions at $\lambda = 1.5$, $p = 16$, $p_{FMM} = 12$.

| $N$       | Basis     | $\epsilon_2$ | $T^{(get)}$, $s$ | $T^{(get)}_{(no,mean)}$, $s$ | $T^{(set)}$, $s$ |
|-----------|-----------|---------------|-----------------|----------------|-------------|
| 27,000    | Spherical | 1.3(-5)       | 1.67            | 1.30           | 0.08        |
|           | RBF       | 2.6(-5)       | 1.74            | 1.35           | 0.06        |
| 216,000   | Spherical | 1.6(-5)       | 12.3            | 9.57           | 0.08        |
|           | RBF       | 2.6(-5)       | 12.7            | 9.95           | 0.06        |

$O(N)$ (with some small $N \log N$ addition due to data structures), the present algorithm is scaled sublinearly at small $N$, while it approaches the same scaling as the regular FMM. Qualitatively this can be explained by $O\left(\sqrt{N_b (N + 2L)}\right)$ scaling of the FMM as it used in the present algorithm for periodization. So the ratio of the FMM time in the present algorithm and the FMM for non-periodic problem can be estimated as

$$\frac{T^{(per)}_{FMM}}{T^{(non)}_{FMM}} \sim \frac{\sqrt{N_b (N + 2L)}}{N} \sim \left(\frac{\pi \sqrt{3}}{2} \lambda^3\right)^{1/2} \left(1 + \frac{2L (\lambda)}{N}\right)^{1/2}.$$  \hspace{1cm} (41)

This ratio for different $\lambda$ along with the experimentally measured time ratio of the “get” part of the algorithm and the FMM for non-periodic problem is plotted in Fig. 6. It is seen that qualitatively this explains the observed results, and at relatively small $N$ the wall clock time of the present algorithm requires is several times larger than the time of non-periodic FMM. At larger times the ratio should come to an asymptotic limit depending on $\lambda$. While the theory predicts that $\lambda = 1.1$ limit should be smaller than $\lambda = 1.3$ limit, we found that in the range of experiments they are approximately the same (time ratio about 2). So one can say that the periodization overhead at large $N$ is approximately the same as the run time of the FMM for the non-periodic problem. The reason why the experimental data deviate from Eq. (41) is that, first, there are some $O(N)$, $O(PN)$, and $O(P^2)$ overheads in the periodization algorithm, and, second, that the optimization of the FMM can be done only discretely (changing the number of the levels of the octree). Indeed, we checked the profiling of the FMM for both cases, and found that while this algorithm is performing at its minimum time for a given distribution and $p_{FMM}$, the parts of the FMM related to translations and direct summation were not balanced perfectly (several times difference).

Finally we note that it is not an easy task to compare the absolute performance of the present methods with the smooth particle mesh Ewald (SPME) and other algorithms for periodic summation due to a difference in implementation, accuracy, what is actually computed (potential and gradient), hardware, etc. However, some comparison with that approaches can be done using published data [14] comparing performance of the SPME and FMM-type PWA implementation for clusters, for relatively small size problems.
Figure 5: Wall clock times of the present algorithm with the FMM as FSA for $\lambda = 1.3$, $p = 34$ (filled squares), and with brute-force 4 core CPU summation (filled triangles), and Tesla C2050 GPU summation (filled circles), for a periodic system replicating distribution of $N$ random sources in a unit cube (potential and gradient computations, Laplacian kernel in three dimensions). The corresponding empty squares, triangles and circles, connected by a solid line show performance of the particular FSA for the non-periodic problem with the same source distribution in the unit cube and the same truncation number $p_{FMM} = 16$. The dashed lines continue the trendlines to indicate performance, if memory resources had not been exceeded. The quadratic scaling of the brute-force summation, and the linear scaling for the FMM are seen. Brute force GPU sums for this case are faster till about $N \simeq 40,000$.

($N = 10^4$ and $N = 10^5$). The absolute figures indicate that the wall clock time of the present algorithm for these sizes is of the same order as the reported times for those methods, while we are able to perform computations with larger problem sizes on relatively modest hardware.

4.3 Using graphics processing units (GPUs) for summation

GPUs are often used to accelerate molecular dynamics simulations where periodization may be required. We implemented the “get” part of the algorithm completely on the GPU. Using “brute-force” summation on the GPU and on a multicore CPU as the FSA, we did the same performance tests as for the FMM, see Fig. 5. The “set” part can be executed on CPU and transferred to the GPU. All parts of the algorithm are highly parallelizable and so while the scaling of the algorithm is quadratic, it is up to two orders of magnitude faster than the multicore CPU version.

Our tests reveal that at large truncation numbers $p$ (typically $p \gtrsim 30$) single precision GPU computations cannot be used for spherical basis function evaluation due to loss of precision in least-squares solution. Double-precision computations provide accurate results with $L_2$-norm relative errors of the order
Figure 6: The ratio of the solution times of the periodic and non-periodic problems using the FMM for the cases shown in Fig. 5. The dotted and the dashed curves show theoretical estimate of the FMM time ratio for $\lambda = 1.1$ and $\lambda = 1.3$, respectively, according to Eq. (41).

$10^{-13} - 10^{-15}$ in potential and gradient computations compared to the double precision CPU computations. Accordingly Fig. 5, only shows double precision times. High performance implementation of the brute-force summation is described in [16] and was used in the present tests with single and double precision, and run on a single NVIDIA Tesla C2050 card. The CPU wall clock times used for comparisons were measured for algorithm parallelized for 4 core PC described before.

Fig. 5 shows the results of tests. It is seen that despite the asymptotic quadratic scaling of the algorithm the GPU implementation can be faster or comparable in speed with the FMM running on CPU at $p_{FMM} = 16$ for problems of size $N \lesssim 10^5$. The ratio of times for solution of periodic and non-periodic problems for brute force computations at large $N$ tends to theoretical limit $T_{brute}^{(per)}/T_{brute}^{(non)} = \pi \sqrt{\frac{2}{3}} \lambda^3$, which for $\lambda = 1.3$ used for illustrations is about 6 times. Of course, this ratio can be reduced by decreasing $\lambda$. However, in the range $N \lesssim 10^5$ reduction of $\lambda$ below 1.1 does not benefit the overall performance due to increase of the size of the expansion and reduction of the performance of evaluation of the far field on the GPU, where the local memory is substantially smaller than the CPU cache. One of the advantages of brute force double precision GPU computing for problems of relatively small size ($N \lesssim 10^5$) is that besides the roundoff errors the accuracy is controlled only by parameters $p$ and $\lambda$, while for the FSA=FMM the error is controlled also by $p_{FMM}$. For efficient FMM implementations on GPU this number is usually small [16] while the efficiency of the FMM on GPU for high accuracy simulations is a subject for a separate study.

5 Conclusion

We have presented a kernel-independent method for the periodization of finite sums. The technique was presented in a general setting, and then applied to the particular case of the Laplacian kernel using different expansion bases. Tests showed that the method can be tuned to compute periodic sums with arbitrary
prescribed accuracy. In the case of use of the FMM as the fast summation algorithm the complexity of the method at large \( N \) is the same as the FMM. The computational time for large \( N \) (in tests up to \( N = 2^{23} \)) is about twice that for the finite box sum using the FMM. Similar results are seen for GPU based FSA, though here the scaling is quadratic, and the largest problem size that can be reasonably treated is about \( 10^5 \).

The ease of implementation of the periodization method, its performance, and capability to “retrofit” any available blackbox FSA without any modification makes it practical. This method may also be valuable on distributed architectures on which communication costs of an algorithm are as important as computational complexity. FMM-based approaches are known to be much more communication efficient than FFT-based approaches for solution of the same large problems on distributed architectures. Additional speedups of the method can be achieved by specialization of the FMM – these were specifically avoided in this paper to demonstrate the ability to use a blackbox sum algorithm, and should be investigated if the method is to be used in a “production” environment. Application to other kernels should be straightforward, though details will have to be worked out for them.

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A  Box integrals of the basis functions

Basis functions $R_{m}^{n}(y)$ are homogeneous polynomials of degree $n$ (sums of monomials $x_{n_{1}}y_{n_{2}}z_{n_{3}}$, $n_{1} + n_{2} + n_{3} = 1$). We can compute the integrals as

\[
R_{t}^{(0)} = \frac{1}{V_{0}} \int_{\Omega_{0}} R_{t}(y) dV(y) = \frac{1}{d_{1}d_{2}d_{3}} \int_{-d_{1}/2}^{d_{1}/2} \int_{-d_{2}/2}^{d_{2}/2} \int_{-d_{3}/2}^{d_{3}/2} R_{t}(x, y, z) dx dy dz
\]

\[
= \frac{1}{8} \sum_{i=1}^{N_{q}} \sum_{j=1}^{N_{q}} \sum_{k=1}^{N_{q}} w_{i}w_{j}w_{k}R_{t}\left(\frac{1}{2}d_{1}x_{i}, \frac{1}{2}d_{2}x_{j}, \frac{1}{2}d_{3}x_{k}\right),
\]

where $w_{i}$ and $x_{i}$ are the standard weights and abscissas of the Gauss quadrature of order $N_{q}$ [29]. Since this integration is exact for polynomials of degree $n < 2N_{q}$ and the maximum degree of the polynomials in the sum is $p - 1$, the choice

\[
N_{q} = \left\lceil \frac{p - 1}{2} \right\rceil + 1, \quad p = P^{1/2},
\]

provides an exact result. For evaluation of all $P$ required coefficients $R_{m}^{n(0)}$, the computational cost of this procedure is $O\left(\frac{P^{5}}{2}\right)$. Note that faster methods may be proposed for computation of this step. However, this is not crucial, as this integration is performed in the “set” part of the algorithm, which overall cost is $O\left(\frac{P^{3}}{2}\right)$, and this cost is amortized over the rest of the algorithm.
B Mean computations

To compute the integral (38) for the kernel (2), we first apply the Gauss divergence theorem to reduce the volume integral to a surface integral:

$$\Phi_0(x) = \frac{1}{V_0} \int_{\Omega_0} dV(y) = \frac{1}{V_0} \int_{\Omega_0} \nabla_y \cdot \left( \frac{y - x}{|y - x|} \right) dV(y) = \frac{1}{2V_0} \int_{\partial\Omega_0} n \cdot (y - x) dS(y),$$

(44)

where \(n\) is the outward normal to \(\Omega_0\). This result is valid for an arbitrary point \(x\) including when \(x\) is located in \(\Omega_0\) or on its boundary \(\partial\Omega_0\). This can be checked by consideration of \(\epsilon\)-vicinities of singularities, which are integrable. The surface integral can be decomposed into integrals over the box faces, \(S_k, k = 1, \ldots, 6\).

$$\Phi_0(x) = \frac{1}{2V_0} \sum_{k=1}^{6} \int_{S_k} n_k \cdot (y - x) dS(y) = -\frac{1}{2V_0} \sum_{k=1}^{6} (n_k \cdot x_k) L_k(x_k),$$

(45)

$$L_k(x_k) = \int_{S_k} \frac{dS(y)}{|y - x_k|}, \quad y_k = y - y_{k0}, \quad x_k = x - y_{k0},$$

where \(n_k\) and \(y_{k0}\) are the normal and the center of the \(k\)th face, while \(y_k\) and \(x_k\) are coordinates in the reference frame with the origin at the \(k\)th face center. The surface integral then can be reduced to the contour integral using, e.g. the Gauss divergence theorem in the plane of a particular face. Indeed, consider function

$$F_k(r_k) = r_k f_k(r_k; h_k), \quad f_k(r_k; h_k) = \frac{\rho_k - h_k}{r_k^2}, \quad \rho_k = \sqrt{r_k^2 + h_k^2} (= |y_k - x_k|),$$

(46)

$$r_k = y_k - x_k', \quad x_k' = x_k - n_k h_k, \quad h_k = n_k \cdot x_k.$$ 

The 2D divergence of this function in the plane of the \(k\)th face is \(1/\rho_k\). So

$$L_k(x_k) = \int_{S_k} \nabla r_k \cdot F_k(r_k) dS(r_k) = \int_{C_k} n_k' \cdot F_k(r_k) dl(r_k),$$

(47)

where \(n_k'\) is the outer normal to the contour \(C_k = \partial\Omega_k\). This integral can be decomposed into four integrals over the face edges. So

$$L_k(x_k) = \sum_{j=1}^{4} I_{kj}(x_k), \quad I_{kj}(x_k) = \int_{C_{kj}} n_{k,j} \cdot r_k \frac{\rho_k - h_k}{r_k^2} dl.$$ 

(48)

The latter integrals can be found analytically. Indeed, consider for the \(j\)th edge a local right hand oriented reference frame centered at its endpoint \(y_{kj0}\) from which integration starts, and unit basis vectors \(i'_{k,jx}\) directed along the integration path, \(i'_{k,jy} = n_k\) and \(i'_{k,jz} = n_{k,jx}' \times i'_{k,jy}'\). Denoting coordinates of \(x\) in this reference frame as

$$x_{kj} = (x_k - y_{kj0}) \cdot i'_{k,jx}, \quad y_{kj} = (x_k - y_{kj0}) \cdot i'_{k,jy}, \quad z_{kj} = (x_k - y_{kj0}) \cdot i'_{k,jz},$$

(49)

we obtain

$$I_{kj}(x_k) = -z_{kj} \int_{-x_{kj}}^{l_{kj} - x_{kj}} f(r_{kj}; |y_{kj}|) dx = H(l_{kj} - x_{kj}, |y_{kj}|, z_{kj}) - H(-x_{kj}, |y_{kj}|, z_{kj}),$$

(50)

where \(r_{kj}^2 = x^2 + z_{kj}^2\), \(l_{kj}\) is the length of edge \(C_{kj}\), and \(H(x, y, z)\) is the primitive,

$$H(x, y, z) = - \int f(r; |y|) dx, \quad f(r; y) = \frac{\rho - |y|}{r^2}, \quad r^2 = x^2 + z^2, \quad \rho = \sqrt{r^2 + y^2},$$

(51)
which can be computed analytically as
\[
H(x, y, z) = |y| \left( \arctan \frac{x}{z} - \arctan \frac{|y|x}{z\rho} \right) - z \ln |\rho + x| + C(y, z). \quad (52)
\]

The integration constant \( C(y, z) \) can be selected arbitrarily to eliminate possible singularities. Particularly for \( y \neq 0, z = 0 \) one can set \( H = 0 \). The above formulae are sufficient for numerical implementation, which in the simplest form can program the primitive (52) and implement the above decompositions. There exist some box symmetries (e.g. all local coordinates are nothing but permuted and shifted original Cartesian coordinates), which can be exploited to achieve better performance.

C Ewald summation

The Ewald summation method is based on decomposition of kernel (2)
\[
K(y - x) = K_1(y - x; \xi) + K_2(y - x; \xi),
\]
(53)
\[
K_1(y - x; \xi) = \frac{\text{erfc}(\xi |y - x|)}{|y - x|}, \quad y \neq x; \quad K_1(0; \xi) = -\frac{2\xi}{\sqrt{\pi}},
\]
\[
K_2(y - x; \xi) = \frac{\text{erf}(\xi |y - x|)}{|y - x|}, \quad \forall y, x \in \mathbb{R}^3, \quad \left( K_2(0; \xi) = \frac{2\xi}{\sqrt{\pi}} \right),
\]

which is exact for any value of parameter \( \xi \), since by definition of the error function, \( \text{erf}(x) \), and the complimentary error function, \( \text{erfc}(x) \), we have \( \text{erf}(x) + \text{erfc}(x) = 1 \) and the value of \( K_1(0; \xi) \) is set due to by definition \( K(0) = 0 \). So for the total potential (2) we have
\[
\phi(y) = \phi_1(y) + \phi_2(y), \quad \phi_1(y) = \sum_{p} \sum_{i=1}^{N} q_i K_1(y - x_i + p; \xi),
\]
\[
\phi_2(y) = \sum_{p} \sum_{i=1}^{N} q_i K_2(y - x_i + p; \xi). \quad (54)
\]

Both functions \( \phi_1(y) \) and \( \phi_2(y) \) are periodic.

Due to fast decay of \( \text{erfc}(x) \) computation of \( \phi_1(y) \) for \( y \in \Omega_0 \) can be done only using the sources in some neighborhood of \( \Omega_0 \), namely in \( \Omega_1 \supset \Omega_0 \) such that the minimum distance, \( a \), between the points on the boundaries \( \partial \Omega_0 \) and \( \partial \Omega_1 \) is much larger than \( 1/\xi \). Hence, this can be computed directly by evaluation of a finite sum with a controllable error as
\[
\phi_1(y) = \sum_{x_i \in \Omega_1(\xi)} q_j K_1(y - x_j; \xi) + O\left( e^{-\xi a^2} \right). \quad (55)
\]

For computation of \( \phi_2(y) \) one can notice that \( K_2 \) is a solution of the Poisson equation
\[
\nabla^2 K_2(y - x; \xi) = -4\pi \delta_{\xi} (y - x), \quad \delta_{\xi}(y - x) = \frac{\xi^3}{\pi^{3/2}} e^{-\xi |y - x|^2}, \quad (56)
\]
where \( \delta_{\xi}(y - x) \) is a compactly supported function, which turns to the Dirac delta-function as \( \xi \to \infty \). Periodic solution of the Poisson equation can be obtained via the FFT. For this purpose, we grid the domain \( \Omega_0 \) and select \( \xi \) in a way that \( \xi \ll 1/h \), and \( \xi \gg 1/\max(d_1, d_2, d_3) \) (an optimal setting can be found from analysis of the error bounds), where \( h \) is the minimum spatial step of the grid. This enables sampling of \( \delta_{\xi}(y - x) \) for source \( x = x_i \) at several grid points around \( x_i \). The number of these grid points determines
the accuracy of the method (at optimal settings), so we introduce additional parameter \( N_r \), so \( \delta_\xi (y) \) is sampled in a box \((2N_r + 1) \times (2N_r + 1) \times (2N_r + 1)\). We also take care about the points \( x_i \) located near the boundary of \( \Omega_0 \) by periodization (so we construct a periodic function \( \delta_\xi^{(p)} (y - x) \)). Further, we apply the forward 3D FFT to

\[
f_2 (y) = \nabla^2 \phi_2 (y) = -4\pi \sum_{i=1}^{N} q_i \delta_\xi^{(p)} (y - x_i),
\]

and zero the harmonic of the Fourier image \( f_2^* (k) \) corresponding to the wavenumber \( k = 0 \). The inverse 3D FFT of \( \phi_2^* (k) = -k^{-2} f_2^* (k) \), produces the required solution \( \phi_2 (y) \) with zero mean at grid points. Note then that solution obtained in this way has the following mean

\[
\phi_{\text{mean}} (\xi) = \left< \phi (y) \right>_{\Omega_0} = \frac{1}{V_0} \sum_{x_j \in \Omega_1 (\xi)} q_j \int_{\Omega_0} K_1 (y - x_j; \xi) \, dV \approx 0.
\]

The zero mean here is due to the compact support of the kernel \( K_1 \) and charge neutrality. This mean can be computed using decomposition \( K_1 (y - x_j; \xi) = K (y - x_j; \xi) - K_2 (y - x_j; \xi) \), where the integral with the first kernel can be computed analytically (see Appendix B), while the integral with the second kernel is regular and can be computed using, say, the trapezoidal rule (in the FFT-based method the space is gridded). To avoid interpolation errors, in the numerical tests where we compared our method for accuracy with the Ewald summation method, we used only cases when the source and evaluation points are located at the grid nodes.