



**HAL**  
open science

# A simple formula of the magnetic potential and of the stray field energy induced by a given magnetization

Tahar Zamene Boulmezaoud

► **To cite this version:**

Tahar Zamene Boulmezaoud. A simple formula of the magnetic potential and of the stray field energy induced by a given magnetization. 2023. hal-03966340

**HAL Id: hal-03966340**

**<https://hal.science/hal-03966340>**

Preprint submitted on 31 Jan 2023

**HAL** is a multi-disciplinary open access archive for the deposit and dissemination of scientific research documents, whether they are published or not. The documents may come from teaching and research institutions in France or abroad, or from public or private research centers.

L'archive ouverte pluridisciplinaire **HAL**, est destinée au dépôt et à la diffusion de documents scientifiques de niveau recherche, publiés ou non, émanant des établissements d'enseignement et de recherche français ou étrangers, des laboratoires publics ou privés.

# A simple formula of the magnetic potential and of the stray field energy induced by a given magnetization

Tahar Zamene BOULMEZAOUD<sup>1,2</sup>

ABSTRACT. The primary aim of this paper is the derivation and the proof of a simple and tractable formula for the stray field energy in micromagnetic problems. The formula is based on an expansion in terms of Arar-Boulmezaoud functions. It remains valid even if the magnetization is not of constant magnitude or if the sample is not geometrically bounded. The paper continues with a direct and important application which consists in a fast summation technique of the stray field energy. The convergence of this technique is established and its efficiency is proved by various numerical experiences.

## CONTENTS

1.	Introduction	1
2.	Overview of Arar-Boulmezaoud functions	4
3.	The first main result: the formulas	6
4.	The second main result: a new method for calculating the stray-field energy	9
5.	Implementation and computational tests	13
6.	Conclusion and perspectives	18
	References	22

## 1. INTRODUCTION

The description and the understanding of magnetic microstructures are often based on the theory of Landau and Lifshitz [28] (see also [16]) which consists in minimizing of the total free energy (see, e. g., [23], [33], [26] and [30]):

$$(1) \quad E_{tot}(M) = \alpha \int_{\Omega} |\nabla M|^2 dx + \int_{\Omega} \phi(M) dx - \frac{1}{2} \int_{\Omega} H_d \cdot J dx - \int_{\Omega} H_{ex} \cdot J dx + E_s,$$

---

<sup>1</sup> Université Paris-Saclay, UVSQ, LMV, Versailles, France.

<sup>2</sup> Department of Mathematics and Statistics, University of Victoria, Victoria, British Columbia, Canada.

*E-mail address:* `tahar.boulmezaoud@uvsq.fr`.

2020 *Mathematics Subject Classification.* 35J47, 35J15, 35C10, 35C20.

*Key words and phrases.* Micromagnetics, Landau–Lifshitz equation, Stray field, Arar-Boulmezaoud functions.

where  $\Omega$  is the sample (or the magnetic body),  $\alpha$  is the exchange stiffness (positive) constant,  $\phi$  is a function describing structural anisotropies,  $H_{ex}$  is an external field,  $H_d$  is the stray (or demagnetizing) field generated by the magnetic body itself and  $E_s$  is the sum the remaining energies (like magnetostrictive self-energy and magneto-elastic interaction energy). The magnetic polarisation  $J$  is given by the formula  $J = \mu_0 M$ , while the stray field  $H_d$  is related to  $M$  by the equations:

$$(2) \quad \mathbf{curl} H_d = 0 \text{ in } \mathbb{R}^3, \quad \text{div}(\mu_0(H_d + M\chi_\Omega)) = 0 \text{ in } \mathbb{R}^3,$$

where  $\chi_\Omega$  denotes the characteristic function of the sample.

The magnetization  $M$  is often subject to the Heisenberg-Weiss constraint

$$(3) \quad |M| = M_s \text{ a. e. in } \Omega,$$

where  $M_s$  is the spontaneous saturation magnetization which is assumed to be constant (and generally depending on the temperature). Although the reader may assume that  $M$  complies with this constraint, we will see that it is not necessary for the validity of the main results stated here; a much weaker constraint on  $M$  suffices (see assumption  $(\mathcal{H}_2)$  below).

In the litterature, much attention is paid to the calculation of the stray field energy resulting from demagnetizing field  $H_d$ :

$$(4) \quad \mathcal{E}_{sf}(M) := -\frac{\mu_0}{2} \int_{\Omega} H_d \cdot M dx.$$

In view of equations (2),  $H_d$  is curl free and can be written into the form

$$(5) \quad H_d = -\nabla U$$

(see [22]), where  $U$  is the magnetic potential which is solution of the Poisson equation in *the whole space*:

$$(6) \quad \Delta U = \text{div}(M\chi_\Omega) \text{ in } \mathbb{R}^3.$$

The stray field energy can be expressed as

$$(7) \quad \mathcal{E}_{sf}(M) = \frac{\mu_0}{2} \int_{\mathbb{R}^3} |\nabla U|^2 dx = \frac{\mu_0}{2} \int_{\mathbb{R}^3} |H_d|^2 dx.$$

Computing the stray field energy (7) is one of the most challenging issues in micromagnetics (see, e. g., [23] and [33]). The difficulty is mainly due to its non local nature. There are several methods for the effective calculation of this energy. Some of these methods are based on solving the elliptic partial differential equation (6) using finite differences method (see, e. g., [7], [38], [2]), or finite elements method (see, e. g., [20, 25], [6], [17]), or inverted finite elements method ([10], [14], [12, 13], [8], [15], [24] and [13]). Other methods are based on the calculation of  $U$  from the integral formula (see, e. g., [9], [29], [32], [19], [37], [27]):

$$(8) \quad U(x) = \frac{1}{4\pi} \int_{\Omega} \frac{(y-x) \cdot M(y)}{|y-x|^3} dy.$$

The primary aim of this work is to establish the following formula

$$(9) \quad \mathcal{E}_{sf}(M) = \frac{\mu_0}{2} \sum_{k=0}^{\infty} \frac{4}{4(k+1)^2 - 1} \sum_{\alpha \in \Lambda_k} \left( \int_{\Omega} M \cdot \nabla \mathcal{W}_{\alpha} dx \right)^2,$$

where  $(\mathcal{W}_{\alpha})_{\alpha}$  designate the Arar-Boulmezaoud functions<sup>1</sup> introduced in [5] and in [11]. These functions will be presented along with their properties in Section 2. We also prove the following formula for the magnetic potential

$$(10) \quad U = \sum_{k=0}^{\infty} \sum_{\alpha \in \Lambda_k} \frac{4}{4(k+1)^2 - 1} \left( \int_{\Omega} M \cdot \nabla \mathcal{W}_{\alpha} dx \right) \mathcal{W}_{\alpha}.$$

Another by-product, as we shall see, concerns approximation of the stray-field energy (7). More precisely, truncating formula (9) gives the approximation

$$(11) \quad \mathcal{E}_{sf}^N(M) = \frac{\mu_0}{2} \sum_{k=0}^N \frac{4}{4(k+1)^2 - 1} \sum_{\alpha \in \Lambda_k} \left( \int_{\Omega} M \cdot \nabla \mathcal{W}_{\alpha} dx \right)^2,$$

When  $M \cdot n = 0$ , we establish the estimate

$$(12) \quad 0 \leq \mathcal{E}_{sf}(M) - \mathcal{E}_{sf}^N(M) \leq CN^{-2} \|\operatorname{div} M\|_{L^2(\Omega)}.$$

The reader interested in formulas above but not in details of the proof can admit that (9) is valid for any connected open set  $\Omega$ , not necessarily bounded, and any (measurable) vector field  $M$  satisfying

$$\int_{\Omega} |M|^2 dx < +\infty$$

Nevertheless, the latter condition is obviously fulfilled when the sample  $\Omega$  has a finite volume and  $M$  satisfying the Heisenberg-Weiss constraint (3).

The rest of the paper is organized as follows. In Section 2 we present Arar-Boulmezaoud functions which are the key ingredient of this paper. Their most useful properties are listed. These properties are essentially known and no originality is claimed in Section 2. The formulas that form the main output of this paper are presented and proved in Section 3. In Section 4, a new method for calculating the energy resulting from these formulas is proposed and analyzed. In particular, the convergence of the method is established. In section 5 focus in on computational tests through several examples. The last section is devoted to a conclusion.

---

<sup>1</sup>Although these functions were discovered by N. Arar and the author in [5], the choice of this appellation is not due to the authors, but to a reviewer of one of [11] who asked to choose this appellation.

## 2. OVERVIEW OF ARAR-BOULMEZAOUD FUNCTIONS

In [5], Arar and the author introduced a family of multi-dimensionnal rational and quasi-rational functions ( $\mathscr{W}_\alpha$ ) which turned out to be particularly appropriate for solving second order elliptic equations in unbounded regions of space (see [11]). This is primarily due to their completeness, their orthogonal properties and their behavior at large distances.

The definition of these functions in  $\mathbb{R}^3$  necessitates the use of spherical harmonics on the unit sphere of  $\mathbb{R}^4$  and the stereographic projection (four dimensional spherical harmonics are less encountered than those on  $\mathbb{S}^2$  the unit sphere of  $\mathbb{R}^3$ ). For each integer  $k \geq 0$ ,  $\mathbb{H}_k$  will be the space of spherical harmonics of degree  $k$  over the unit sphere (see, e. g., [34], [31], [35], [36], [21]):

$$\mathbb{S}^3 := \{x \in \mathbb{R}^4 \mid |x| = 1\}$$

(spherical harmonics of degree  $k$  on  $\mathbb{S}^3$  are restrictions to  $\mathbb{S}^3$  of harmonic homogeneous polynomials of degree  $k$  on  $\mathbb{R}^4$ ). We know that

$$(13) \quad \dim \mathbb{H}_k = (k+1)^2 \text{ for all } k \geq 0.$$

In order to construct an orthogonal basis of  $\mathbb{H}_k$ , we set

$$\Lambda = \{(i, \ell, m) \in \mathbb{N}^2 \times \mathbb{Z} \mid 0 \leq \ell \leq i \text{ and } -\ell \leq m \leq \ell\},$$

and for each integer  $k \geq 0$

$$\begin{aligned} \Lambda_k &= \{(i, \ell, m) \in \Lambda \mid i = k\}, \\ \Lambda_k^* &= \bigcup_{i=0}^k \Lambda_i. \end{aligned}$$

If  $\alpha = (k, \ell, m)$ ,  $\beta = (k', \ell', m') \in \Lambda$ , then  $\delta_{\alpha, \beta}$  denotes the usual Kronecker symbol of  $\alpha, \beta$ , that is  $\delta_{\alpha, \beta} = \delta_{k, k'} \delta_{\ell, \ell'} \delta_{m, m'}$ . Define the spherical coordinates for  $\mathbb{S}^3$  as the triplet  $(\phi, \theta, \chi)$  such that  $0 \leq \phi < 2\pi$ ,  $0 \leq \theta \leq \pi$ ,  $0 \leq \chi \leq \pi$  and

$$(14) \quad \xi = (\sin \theta \cos \phi \sin \chi, \sin \theta \sin \phi \sin \chi, \cos \theta \sin \chi, \cos \chi),$$

Spherical harmonics on  $\mathbb{S}^3$  are defined by:

$$(15) \quad \mathscr{Y}_\alpha(\xi) = \frac{1}{\sqrt{a_{k, \ell}}} (\sin \chi)^\ell T_{k+1}^{(\ell+1)}(\cos \chi) Y_{\ell, m}(\phi, \theta), \text{ for } \alpha = (k, \ell, m) \in \Lambda.$$

Here

- $(T_k)_{k \geq 0}$  designate Chebyshev polynomials of the first kind satisfying

$$\cos(k\theta) = T_k(\cos \theta) \text{ for } \theta \in \mathbb{R},$$

- $(Y_{\ell, m})_{\ell, m}$  are the usual real spherical harmonics on  $\mathbb{S}^2$ :

$$(16) \quad Y_{\ell, m}(\phi, \theta) = \eta_\ell K_\ell^{|m|}(\cos \theta) y_m(\phi)$$

where

$$(17) \quad y_m(\phi) = \begin{cases} \cos(m\phi) & \text{if } m \geq 1, \\ \frac{1}{\sqrt{2}} & \text{if } m = 0, \\ \sin(|m|\phi) & \text{if } m \leq -1, \end{cases}$$

$$(18) \quad \eta_\ell = \sqrt{\frac{2\ell + 1}{2\pi}},$$

and

$$(19) \quad K_\ell^m(x) = (-1)^m \sqrt{\frac{(\ell - m)!}{(\ell + m)!}} P_\ell^m(x), \quad \text{for } -\ell \leq m \leq \ell$$

(thus,  $K_\ell^{-m} = (-1)^m K_\ell^m$ ). Here  $(P_\ell^m)_{\ell,m}$  designate the associated Legendre functions defined as:

$$P_\ell^m(t) = \frac{(-1)^m}{2^\ell \ell!} (1 - t^2)^{m/2} \frac{d^{\ell+m}}{dt^{\ell+m}} (t^2 - 1)^\ell, \quad -\ell \leq m \leq \ell$$

(some authors omit the  $(-1)^m$  factor, commonly referred to as the Condon-Shortley phase, or append it in the definition of  $Y_{\ell,m}$ ). We also adopt the convention  $P_\ell^m = 0$  and  $K_\ell^m = 0$  when  $|m| > \ell$ .

–  $(a_{k,\ell})$  are normalization constants given by

$$(20) \quad a_{k,\ell} = \frac{(k+1)\pi (k+\ell+1)!}{2 (k-\ell)!}.$$

The following properties hold true

- For all  $k \geq 0$ ,  $(\mathcal{Y}_\alpha)_{\alpha \in \Lambda_k}$  is a basis of  $\mathbb{H}_k$ .
- For all  $\alpha, \beta \in \Lambda$

$$(21) \quad \int_{\mathbb{S}^3} \mathcal{Y}_\alpha(\xi) \mathcal{Y}_\beta(\xi) dS(\xi) = \delta_{\alpha,\beta}.$$

- For all  $k \geq 0$  and  $\alpha \in \Lambda_k$ ,

$$-\Delta_S \mathcal{Y}_\alpha = k(k+2) \mathcal{Y}_\alpha,$$

where  $\Delta_S$  is the Laplace-Beltrami operator over the unit sphere  $\mathbb{S}^3$ . In terms of spherical coordinates, this operator is given by

$$\frac{1}{\sin^2 \chi} \left\{ \frac{\partial}{\partial \chi} \left( \sin^2 \chi \frac{\partial}{\partial \chi} \right) + \frac{1}{\sin \theta} \frac{\partial}{\partial \theta} \left( \sin \theta \frac{\partial}{\partial \theta} \right) + \frac{1}{\sin^2 \theta} \frac{\partial^2}{\partial \phi^2} \right\}.$$

In the three-dimensional situation (the only one that interests us here), Arar-Boulmezaoud functions are defined as follows (see [5] and [11]): for any  $\alpha \in \Lambda$

$$(22) \quad \mathcal{W}_\alpha(x) = \left( \frac{2}{|x|^2 + 1} \right)^{\frac{1}{2}} \mathcal{Y}_\alpha(\pi^{-1}(x)).$$

Here  $\pi$  denotes the stereographic projection defined on  $\mathbb{S}_*^3 = \mathbb{S}^3 - \{(0, \dots, 0, 1)\}$  by

$$\begin{aligned} \pi : \mathbb{S}_*^3 &\longrightarrow \mathbb{R}^3 \\ \xi &\longmapsto \left( \frac{\xi_1}{1 - \xi_4}, \frac{\xi_2}{1 - \xi_4}, \frac{\xi_3}{1 - \xi_4} \right). \end{aligned}$$

Its inverse is given by

$$\begin{aligned} \pi^{-1} : \mathbb{R}^3 &\longrightarrow \mathbb{S}_*^3 \\ x &\longmapsto \left( \frac{2x_1}{|x|^2 + 1}, \frac{2x_2}{|x|^2 + 1}, \frac{2x_3}{|x|^2 + 1}, \frac{|x|^2 - 1}{|x|^2 + 1} \right). \end{aligned}$$

Functions  $(\mathcal{W}_\alpha)_{\alpha \in \Lambda}$  were discovered by Arar and the author in [5] in studying spectrum of weighted Laplacians in  $\mathbb{R}^n$ . In Table 6 of Appendix B, the expressions of the first functions  $(\mathcal{W}_\alpha)_{\alpha \in \Lambda}$  are given explicitly. We can then see that these functions have a rational nature. This is a general property as will be announced later. One can also consult [5] and [11] for higher dimensions and for  $n = 1$  or  $n = 2$ .

In the following Proposition we summarize some useful properties of the functions  $(\mathcal{W}_\alpha)_{\alpha \in \Lambda}$ . We refer to [5] and [11] for their proofs.

**Proposition 2.1.** *Let  $k \geq 0$  be an integer and  $\alpha \in \Lambda_k$ . Then,*

– we have

$$(23) \quad -\Delta \mathcal{W}_\alpha = \frac{(2k+1)(2k+3)}{(|x|^2+1)^2} \mathcal{W}_\alpha,$$

– there exists  $k+1$  polynomial functions  $p_0, \dots, p_k$  such that:

$$(24) \quad \mathcal{W}_\alpha(x) = \sum_{i=0}^k \frac{p_i(x)}{(|x|^2+1)^{i+1/2}},$$

where for each  $i \leq k$ ,  $p_i$  is of degree less than or equal to  $i$ ,

– for all  $\beta \in \Lambda$

$$(25) \quad \int_{\mathbb{R}^3} \frac{\mathcal{W}_\alpha(x) \mathcal{W}_\beta(x)}{(|x|^2+1)^2} dx = \frac{1}{4} \delta_{\alpha,\beta},$$

$$(26) \quad \int_{\mathbb{R}^3} \nabla \mathcal{W}_\alpha(x) \cdot \nabla \mathcal{W}_\beta(x) dx = \frac{(2k+1)(2k+3)}{4} \delta_{\alpha,\beta}.$$

The orthogonality identities (25) and (26) are among the most important properties of Arar-Boulmezaoud functions. as we will see later. More particularly, these relations will be the cornerstone of the formula given in this paper and of the resulting numerical approximation.

Here ends this first enumeration of the properties of functions  $(\mathcal{W}_\alpha)$ . We will need other properties later on, in particular for the calculation of gradients (see paragraph 5.1).

### 3. THE FIRST MAIN RESULT: THE FORMULAS

The objective here is to prove formulas (9) and (10) announced in the introduction. These formulas will be used in the next section to propose a new method for computing stray-field energy. However, before stating the first main result, it is appropriate to give some basics concerning the underlying functional framework we use here. In particular, we show the well-posed nature of the

equation (6).

Here and subsequently, we assume that

( $\mathcal{H}_1$ ) the material fills a connected open set  $\Omega \subset \mathbb{R}^3$  having a lipschitzian boundary,

( $\mathcal{H}_2$ ) the magnetization field  $M$  is defined and measurable over  $\Omega$  and satisfies

$$(27) \quad \int_{\Omega} |M|^2 dx < \infty,$$

that is  $M \in L^2(\Omega)^3$ .

Assumption ( $\mathcal{H}_2$ ) is obviously fulfilled when  $|M|$  satisfies the Heisenberg-Weiss constraint (3) and  $|\Omega| < \infty$  since

$$\|M\|_{L^2(\Omega)^3}^2 = \int_{\Omega} |M|^2 dx = |M|^2 |\Omega| < +\infty.$$

Despite this, we assume neither that  $\Omega$  is bounded nor that  $|M|$  is satisfying the Heisenberg-Weiss constraint (3). Only assumptions ( $\mathcal{H}_1$ ) and ( $\mathcal{H}_2$ ) are needed here.

We now introduce some weighted function spaces. For all integers  $\ell \in \mathbb{Z}$  and  $m \geq 0$ ,  $W_{\ell}^m(\mathbb{R}^3)$  stands for the space of functions  $v$  satisfying

$$\forall |\lambda| \leq m, (1 + |x|^2)^{(\ell + |\lambda| - m)/2} D^{\lambda} v \in L^2(\mathbb{R}^3).$$

This space is equipped with the norm

$$(28) \quad \|v\|_{W_{\ell}^m(\mathbb{R}^3)} = \left( \sum_{|\lambda| \leq m} \int_{\mathbb{R}^3} (|x|^2 + 1)^{|\lambda| + \ell - m} |D^{\lambda} v|^2 dx \right)^{1/2}.$$

When  $m \geq 1$ , the following inclusions hold:

$$W_{\ell}^m(\mathbb{R}^3) \hookrightarrow W_{\ell-1}^{m-1}(\mathbb{R}^3) \hookrightarrow \dots \hookrightarrow W_{\ell-m+1}^1(\mathbb{R}^3) \hookrightarrow W_{\ell-m}^0(\mathbb{R}^3).$$

The following asymptotic property holds true for any function  $v \in W_{\ell}^m(\mathbb{R}^3)$  (see, e. g., [3])

$$(29) \quad \lim_{|x| \rightarrow +\infty} |x|^{\ell - m + 3/2} \|v(|x|, \cdot)\|_{L^2(\mathbb{S}^2)} = 0.$$

where  $\mathbb{S}^2$  is the unit sphere of  $\mathbb{R}^3$  and

$$(30) \quad \|v(|x|, \cdot)\|_{L^2(\mathbb{S}^2)}^2 = \int_{\mathbb{S}^2} |v(|x|, \sigma)|^2 d\sigma.$$

Let us mention the following Hardy's type inequality in  $W_0^1(\mathbb{R}^3)$  (see [13]):

$$(31) \quad \forall v \in W_0^1(\mathbb{R}^3), \int_{\mathbb{R}^3} \frac{|v|^2}{|x|^2 + 1} dx \leq 4 \int_{\mathbb{R}^3} |\nabla v|^2 dx.$$

Thus, from now on, we shall consider that the Hilbert space  $W_0^1(\mathbb{R}^3)$  is endowed with the scalar product

$$((v, w))_{W_0^1(\mathbb{R}^3)} = \int_{\mathbb{R}^3} \nabla v \cdot \nabla w dx,$$



and with the corresponding norm

$$\|v\|_{W_0^1(\mathbb{R}^3)} = \|\nabla v\|_{L^2(\mathbb{R}^3)},$$

which is equivalent to the norm  $\|\cdot\|_{W_0^1(\mathbb{R}^3)}$ .

Here, we look for a solution  $U$  of (6) satisfying

$$(32) \quad \int_{\mathbb{R}^3} |\nabla U|^2 dx < \infty,$$

The first main result of this paper is summarized as follows:

**Theorem 3.1.** *Assume that assumptions  $(\mathcal{H}_1)$  and  $(\mathcal{H}_2)$  hold true. Then (6) has a unique solution  $U \in W_0^1(\mathbb{R}^3)$  which is given by*

$$(33) \quad U = \sum_{k=0}^{\infty} \sum_{\alpha \in \Lambda_k} \frac{4}{(2k+1)(2k+3)} \left( \int_{\Omega} M \cdot \nabla \mathcal{W}_{\alpha} dx \right) \mathcal{W}_{\alpha},$$

where the serie in the right-hand side converges in  $W_0^1(\mathbb{R}^3)$ . The corresponding stray field energy is given by

$$(34) \quad \mathcal{E}_{sf}(U) = \sum_{k=0}^{\infty} \frac{2\mu_0}{(2k+1)(2k+3)} \sum_{\alpha \in \Lambda_k} \left( \int_{\Omega} M \cdot \nabla \mathcal{W}_{\alpha} dx \right)^2.$$

Moreover,

(1) we have

$$(35) \quad \|(1 + |x|^2)^{-1/2} U\|_{L^2(\mathbb{R}^3)} \leq 4 \|M\|_{L^2(\Omega)},$$

$$(36) \quad \mathcal{E}_{sf}(U) \leq \frac{\mu_0}{2} \|M\|_{L^2(\Omega)}^2.$$

(2)  $U \in L^2(\mathbb{R}^3)$ ,  $(1 + |x|^2)^{1/2} \nabla U \in L^2(\mathbb{R}^3)^3$  and

$$(37) \quad \|U\|_{L^2(\mathbb{R}^3)} + \|(1 + |x|^2)^{1/2} \nabla U\|_{L^2(\mathbb{R}^3)^3} \leq C_0(\Omega) \|M\|_{L^2(\Omega)},$$

for some constant  $C_0(\Omega) > 0$  depending only on  $\Omega$ .

(3) we have

$$(38) \quad \lim_{|x| \rightarrow +\infty} |x|^{3/2} \|U(|x|, \cdot)\|_{L^2(\mathbb{S}^2)} = 0,$$

Issues concerning the regularity of the solution  $U$  are postponed to next section (see Theorem 4.4).

*Proof.* We can reformulate equation (6) as follows: find  $U \in W_0^1(\mathbb{R}^3)$  such that

$$(39) \quad \forall v \in W_0^1(\mathbb{R}^3), \quad \int_{\mathbb{R}^3} \nabla U \cdot \nabla v dx = \int_{\Omega} M \cdot \nabla v dx.$$

The existence and uniqueness of solutions is a direct consequence of the Lax-Milgram theorem. Estimate (36) results from the use of Cauchy-Schwarz inequality on the right when  $v = u$  in (39). Combining with Hardy inequality (31) gives (35).

We also have the following lemma (see [5] and [11]):

**Lemma 3.2.** *The family  $(\mathcal{W}_\alpha)_{\alpha \in \Lambda}$  is a Hilbert basis of  $W_0^1(\mathbb{R}^2)$  endowed with the norm  $|\cdot|_{W_0^1(\mathbb{R}^3)}$ .*

Thus,

$$\begin{aligned} U &= \sum_{\alpha \in \Lambda} \frac{((U, \mathcal{W}_\alpha))_{W_0^1(\mathbb{R}^3)}}{((\mathcal{W}_\alpha, \mathcal{W}_\alpha))_{W_0^1(\mathbb{R}^3)}} \mathcal{W}_\alpha, \\ &= \sum_{k=0}^{+\infty} \frac{4}{(2k+1)(2k+3)} \sum_{\alpha \in \Lambda_k} (\nabla U, \nabla \mathcal{W}_\alpha)_{L^2(\mathbb{R}^3)^3} \mathcal{W}_\alpha. \end{aligned}$$

Combining with (39) and (26) gives (33). Since convergence of the right-hand side holds in  $W_0^1(\mathbb{R}^3)$  we also get (34). The reader can refer to [13] for estimate (37). Hence,  $U \in W_1^1(\mathbb{R}^3)$  and (29) holds true with  $\ell = m = 1$ . This gives (38).  $\square$

#### 4. THE SECOND MAIN RESULT: A NEW METHOD FOR CALCULATING THE STRAY-FIELD ENERGY

The main purpose of this section is to show that from the two formulas (33) and (34) results a very efficient and easy to implement numerical method for calculating the stray field energy. This numerical method could be seen as a spectral method in an unbounded domain. However, unlike the usual spectral methods in a bounded domain and which use polynomial functions or trigonometric functions, here we use (quasi)-rational functions guaranteeing a decay of the solution at large distances. Indeed, in view of Proposition 2.1, the functions  $(\mathcal{W}_\alpha)_\alpha$  are rationals up to a multiplicative factor.

**4.1. The method.** In view of (34), the energy  $\mathcal{E}_{sf}(U)$  can be reasonably approximated by truncating the sum. For this end, we set for each  $N \geq 1$

$$(40) \quad \mathcal{E}_{sf}^N(U) = \sum_{k=0}^N \sum_{\alpha \in \Lambda_k} \frac{2\mu_0}{(2k+1)(2k+3)} \left( \int_{\Omega} M \cdot \nabla \mathcal{W}_\alpha dx \right)^2.$$

We observe that

$$(41) \quad \mathcal{E}_{sf}^N(U) = \frac{\mu_0}{2} \int_{\mathbb{R}^3} |\nabla U_N|^2 dx,$$

where

$$(42) \quad U_N = \sum_{\ell=0}^N \sum_{\alpha \in \Lambda_k} \frac{4}{(2k+1)(2k+3)} \left( \int_{\Omega} M \cdot \nabla \mathcal{W}_\alpha dx \right) \mathcal{W}_\alpha.$$

Let us give another interpretation of  $U_N$ . Define the family of finite dimensional spaces  $(H_N)_{N \geq 0}$  as follows: for  $N \geq 0$ ,  $H_N$  is the space of functions of the form

$$(43) \quad v(x) = \sum_{k=0}^N \frac{p_k(x)}{(|x|^2 + 1)^{k+1/2}}, \quad x \in \mathbb{R}^3,$$

where, for each  $k \leq N$ ,  $p_k$  is a polynomial of degree less than or equal to  $k$ . Obviously,

$$(44) \quad H_0 \subset H_1 \subset H_2 \subset \cdots \subset H_N \subset \cdots$$

The following inclusion holds for  $N \geq 0$ :

$$(45) \quad H_N \hookrightarrow W_0^1(\mathbb{R}^3).$$

It can be easily proved that (see, e. g., [5])

$$(46) \quad \dim H_N = \binom{3+N}{3} + \binom{N+2}{3} = \frac{(N+1)(N+2)(2N+3)}{6}.$$

Since

$$|\Lambda_k| = (k+1)^2 \text{ for } k \geq 0$$

( $|\Lambda_k|$  designates the cardinal of the set  $\Lambda_k$ ), we deduce the identity

$$(47) \quad \dim H_N = \sum_{k=0}^N |\Lambda_k| = |\Lambda_N^*|.$$

On the other hand, in view of Proposition 2.1, we have for all  $N \geq 0$ ,

$$(\alpha \in \Lambda_\ell \text{ for some } \ell \leq N) \implies \mathscr{W}_\alpha \in H_N.$$

In other words,

$$\{\mathscr{W}_\alpha; \alpha \in \Lambda_N^*\} \subset H_N.$$

Combining the latter with (47) and with orthogonality properties (25) and (26) gives

**Lemma 4.1.** *For all  $N \geq 1$ , the family  $(\mathscr{W}_\alpha)_{\alpha \in \Lambda_N^*}$  is a basis of  $H_N$ .*

Now, we state this

**Proposition 4.2.** *The function  $U_N$  given by formula (42) is also the unique solution of the well-posed discrete problem*

$$(48) \quad \forall v_N \in H_N, \int_{\mathbb{R}^3} \nabla U_N \cdot \nabla v_N dx = \int_{\Omega} M \cdot \nabla v_N dx.$$

*In addition,  $U_N$  is the projection of  $U$  on  $H_N$  with respect to the inner product  $((\cdot, \cdot))_{W_0^1(\mathbb{R}^3)}$ .*

One could therefore consider that the approximation (42) is none other than the solution of the discrete problem (48) which consists to approximate the original problem (6) by a spectral method using the functions of  $H_N$ . The use of the family  $(\mathscr{W}_\alpha)_{\alpha \in \Lambda_N^*}$  as a basis of  $H_N$  reduces the discrete problem (48) to a simple *diagonal linear system*

$$(49) \quad DX = B$$

with  $D$  the diagonal matrix

$$D = \text{diag} \left( \frac{4}{3}, \underbrace{\frac{4}{15}, \dots, \frac{4}{15}}_{4 \text{ coefficients}}, \dots, \underbrace{\frac{4}{4(N+1)^2-1}, \dots, \frac{4}{4(N+1)^2-1}}_{(N+1)^2 \text{ coefficients}} \right).$$

Here  $X$  contains the components of  $U_N$  with respect to the basis  $(\mathcal{W}_\alpha)_{\alpha \in \Lambda_N^*}$  and  $B$  covers the integrals  $\int_{\Omega} M \cdot \nabla \mathcal{W}_\alpha dx$ . Thus, solution of (48) is obviously given by formula (33). This is a significant observation which demonstrates the benefits of using functions  $(\mathcal{W}_\alpha)_\alpha$ .

#### 4.2. Convergence of the method and error estimate.

Focus now is on convergence when  $N \rightarrow +\infty$ . We have:

**Lemma 4.3.** *Assume that  $(\mathcal{H}_1)$  and  $(\mathcal{H}_2)$  are fulfilled. Then,*

$$(50) \quad \mathcal{E}_{sf}(U) - \mathcal{E}_{sf}(U_N) = \frac{\mu_0}{2} |U - U_N|_{W_0^1(\mathbb{R}^3)}^2.$$

and

$$(51) \quad \lim_{N \rightarrow +\infty} \mathcal{E}_{sf}^N(U_N) = \mathcal{E}_{sf}(U).$$

*Proof.* We first observe that (51) is a direct consequence of (33). Indeed, in view of (39) and (48) we get

$$\int_{\mathbb{R}^3} (\nabla U - \nabla U_N) \cdot \nabla U_N dx = 0,$$

and (50) follows immediately.  $\square$

**Theorem 4.4.** *Assume that  $(\mathcal{H}_1)$  and  $(\mathcal{H}_2)$  are fulfilled. Assume also that  $\Omega$  is bounded,  $\operatorname{div} M \in L^2(\Omega)$  and  $M \cdot n = 0$  on  $\partial\Omega$ . Then,  $U \in W_2^2(\mathbb{R}^3)$  and there exists a constant  $C_1$  depending only on  $\Omega$  such that*

$$(52) \quad \|U - U_N\|_{W_0^1(\mathbb{R}^3)} \leq \frac{C_1}{N} \|\operatorname{div} M\|_{L^2(\Omega)},$$

$$(53) \quad 0 \leq \mathcal{E}_{sf}(U) - \mathcal{E}_{sf}(U_N) \leq \frac{C_1^2}{N^2} \|\operatorname{div} M\|_{L^2(\Omega)}^2.$$

If in addition  $\operatorname{div} M \in H_0^{k-1}(\Omega)$  for some integer  $k \geq 2$  and if

$$(54) \quad \int_{\Omega} M \cdot \nabla q dx = 0 \text{ for all } q \in \mathbb{P}_{k-1}^\Delta,$$

then  $U \in W_{2k}^{k+1}(\mathbb{R}^3)$  and there exists a constant  $C_k$  depending only on  $k$  and  $\Omega$  such that

$$(55) \quad \|U - U_N\|_{W_0^1(\mathbb{R}^3)} \leq C_k N^{-k} \|\operatorname{div} M\|_{H^{k-1}(\Omega)}^2,$$

$$(56) \quad 0 \leq \mathcal{E}_{sf}(U) - \mathcal{E}_{sf}(U_N) \leq C_k^2 N^{-2k} \|\operatorname{div} M\|_{H^{k-1}(\Omega)}^2$$

Here, the usual Sobolev space  $H_0^{k-1}(\Omega)$  designates the closure of  $\mathcal{C}_0^\infty(\Omega)$  in the usual Sobolev space  $H^{k-1}(\Omega)$ .

*Proof.* Firstly, we adopt the following notation: given a function  $f$  defined over  $\Omega$ , we denote by  $\widetilde{f}$  its extension to  $\mathbb{R}^3$  defined as

$$\widetilde{f} = \begin{cases} f & \text{in } \Omega, \\ 0 & \text{in } \mathbb{R}^3 \setminus \Omega. \end{cases}$$

The following lemma is due to [4] (Theorem 6.6):

**Lemma 4.5.** *Let  $m \geq 1$  and  $\ell \geq 1$  be two integers. Then, the Laplace operator  $\Delta$  defined by*

$$\Delta : W_{\ell+m}^{1+m}(\mathbb{R}^3) \rightarrow W_{\ell+m}^{-1+m}(\mathbb{R}^3) \perp \mathbb{P}_{\ell-1}^\Delta,$$

*is an isomorphism. Here  $\mathbb{P}_{\ell-1}^\Delta = \{p \in \mathbb{P}_{\ell-1} \mid \Delta p = 0\}$  and*

$$W_{\ell+m}^{-1+m}(\mathbb{R}^3) \perp \mathbb{P}_{\ell-1}^\Delta = \left\{ f \in W_{\ell+m}^{-1+m}(\mathbb{R}^3) \mid \int_{\mathbb{R}^3} f q dx = 0 \text{ for all } q \in \mathbb{P}_{\ell-1}^\Delta \right\}.$$

Assume now that  $\operatorname{div} M \in L^2(\Omega)$  and  $M.n = 0$  on  $\partial\Omega$ . Then,  $U$  is solution of the problem

$$(57) \quad \Delta U = \widetilde{\operatorname{div} M} \text{ in } \mathbb{R}^3.$$

Obviously  $\widetilde{\operatorname{div} M} \in W_2^0(\mathbb{R}^3)$  and

$$\int_{\mathbb{R}^3} \widetilde{\operatorname{div} M} dx = \int_{\Omega} \operatorname{div} M dx = 0.$$

In view of condition (54) and Lemma 4.5, we deduce that  $U \in W_2^2(\mathbb{R}^3)$ . If in addition  $\operatorname{div} M \in H_0^{k-1}(\Omega)$  for some  $k \geq 1$  and if  $M$  satisfies condition (54) when  $k \geq 2$ , then  $\widetilde{\operatorname{div} M} \in W_s^{k-1}(\mathbb{R}^3)$  for any real number  $s$  (since  $\widetilde{\operatorname{div} M}$  vanishes outside  $\Omega$ ). In particular  $\widetilde{\operatorname{div} M} \in W_{2k}^{k-1}(\mathbb{R}^3)$ . By Green's formula we also have

$$\forall q \in \mathbb{P}_{k-1}^\Delta \int_{\mathbb{R}^3} \widetilde{\operatorname{div} M} q dx = \int_{\Omega} \operatorname{div} M q dx = - \int_{\Omega} M . \nabla q dx = 0.$$

Hence,  $U \in W_{2k}^{k+1}(\mathbb{R}^3)$ , thanks to Lemma 4.5. Moreover, there exists a constant  $C_k$  depending only on  $k$  such that

$$(58) \quad \|U\|_{W_{2k}^{k+1}(\mathbb{R}^3)} \leq C_k \|\widetilde{\operatorname{div} M}\|_{W_{2k}^{k-1}(\mathbb{R}^3)} \leq \widetilde{C}_k \|\operatorname{div} M\|_{H^{k-1}(\Omega)}.$$

Let  $\pi_N$  be the orthogonal projector on  $H_N$  with respect to the scalar product associated to the norm  $|\cdot|_{W_0^1(\mathbb{R}^3)}$ . The following result is due to [11]:

**Lemma 4.6.** *Assume that  $v \in W_{2k}^{k+1}(\mathbb{R}^3)$  for some integer  $k \geq 0$ . Then,*

$$(59) \quad \|\nabla v - \nabla(\pi_N v)\|_{L^2(\mathbb{R}^3)^3} \leq C_k^* N^{-k} \|v\|_{W_{2k}^{k+1}(\mathbb{R}^3)},$$

where  $C_k^*$  is a constant which depends neither on  $N$  nor on  $v$ .

We know that  $U_N = \pi_N U$ . The inequalities (52) and (53) result from (59) and (58) with  $k = 1$  and from (31). The inequalities (55) and (56) are deduced in a similar way. □

REMARK – Assumption  $M.n = 0$  on  $\partial\Omega$  means that the effective magnetic charges are zero. One can easily see that if  $M.n \neq 0$  on  $\partial\Omega$  then  $U$  does not belong to  $W_2^2(\mathbb{R}^3)$ . Indeed, equation (6) can be rewritten as

$$\left\{ \begin{array}{ll} \Delta u = \operatorname{div} M & \text{in } \Omega, \\ \Delta u = 0 & \text{in } \mathbb{R}^3 \setminus \bar{\Omega}, \\ [u] = 0 & \text{on } \partial\Omega, \\ \left[ \frac{\partial u}{\partial n} \right] = -M.n & \text{on } \partial\Omega, \end{array} \right.$$

where  $n$  is the exterior normal on  $\partial\Omega$ . Thus,  $\left[ \frac{\partial u}{\partial n} \right] \neq 0$  on  $\partial\Omega$  and  $U \notin W_2^2(\mathbb{R}^3)$ .

## 5. IMPLEMENTATION AND COMPUTATIONAL TESTS

The first purpose of this section is to examine the numerical results obtained after implementation of the method suggested in the previous section and to check whether the theoretical error estimates are confirmed numerically and whether they are optimal. Another goal is to give some additional details regarding the implementation of the method, including the calculation of integrals. It is worth noting at this early stage that despite the three-dimensional nature of the problem, and despite the fact that it is posed in an open domain, the implementation of the method remains rather easy and fast.

### 5.1. Additional details about gradients of Arar-Boulmezaoud functions.

Formulas in Theorem 3.1 as well as the approximation method proposed in Section 4.1 involve functions  $(\mathcal{W}_\alpha)_\alpha$  by their gradients, particularly in the integral coefficients

$$(60) \quad \int_{\Omega} M \cdot \nabla \mathcal{W}_\alpha dx.$$

In practice, during the implementation of the method, the precise calculation of these gradients could be of great importance. It is consequently preferable to compute them by exact analytical expressions and not by discretization of the differentiation operators. Of course, one can use a Green's formula in (60) to make these gradients disappear:

$$\int_{\Omega} M \cdot \nabla \mathcal{W}_\alpha dx = - \int_{\Omega} (\operatorname{div} M) \mathcal{W}_\alpha dx + \langle M.n, \mathcal{W}_\alpha \rangle_{\partial\Omega},$$

( $\langle \cdot, \cdot \rangle_{\partial\Omega}$  designates the duality pairing between  $H^{1/2}(\partial\Omega)$  and  $H^{-1/2}(\partial\Omega)$ ). However, this requires a little more regularity on the magnetization vector field  $M$  (for example that  $\operatorname{div} M \in L^2(\Omega)$ ) and, moreover, it makes surface integrals appear. It will therefore not be useless to spell out the gradients  $(\nabla \mathcal{W}_\alpha)_\alpha$ . Actually, in view of (15) and (22), these gradients are not quite easy to calculate, especially because of the special functions appear in their formulas (that is, Chebyshev polynomials and associated Legendre functions of Legendre).

In this paragraph, we deduce simpler and exact expressions to the gradients of

the functions  $(\mathcal{W}_\alpha)_\alpha$ , in order to facilitate the computation of magnetic potential and the stray-field energy by formulas (42) and (40).

The starting point is the following proposition

**Proposition 5.1.** *For  $\alpha \in \Lambda$  and  $x \in \mathbb{R}^3$ :*

$$(61) \quad \nabla \mathcal{W}_\alpha(x) = (1 - \xi_4)^{1/2} (\mathcal{V}_\alpha(\xi) - \frac{1}{2} \mathcal{Y}_\alpha(\xi) \hat{\xi}),$$

where  $\xi = (\xi_1, \xi_2, \xi_3, \xi_4) = \pi^{-1}(x) \in \mathbb{S}^3$ ,  $\hat{\xi} = (\xi_1, \xi_2, \xi_3)$ ,  $(\phi, \theta, \chi)$  are the spherical coordinates of  $\xi$  (see (14)) and

$$(62) \quad \mathcal{V}_\alpha(\xi) = (1 - \cos \chi) \begin{pmatrix} -\sin \phi & \cos \phi \cos \theta & -\cos \phi \sin \theta \\ \cos \phi & \sin \phi \cos \theta & -\sin \phi \sin \theta \\ 0 & -\sin \theta & -\cos \theta \end{pmatrix} \begin{pmatrix} \frac{1}{\sin \theta \sin \chi} \frac{\partial \mathcal{Y}_\alpha}{\partial \phi}(\xi) \\ \frac{1}{\sin \chi} \frac{\partial \mathcal{Y}_\alpha}{\partial \theta}(\xi) \\ \frac{\partial \mathcal{Y}_\alpha}{\partial \chi}(\xi) \end{pmatrix}.$$

By the sake of simplicity, proof of Proposition 5.1 is postponed to Appendix A.

REMARK – In Proposition 5.1,  $\frac{\partial \mathcal{Y}_\alpha}{\partial \phi}$ ,  $\frac{\partial \mathcal{Y}_\alpha}{\partial \theta}$  and  $\frac{\partial \mathcal{Y}_\alpha}{\partial \chi}$  designate (abusively) the derivatives of  $\mathcal{Y}_\alpha$  considered as a function of  $\theta$ ,  $\phi$  and  $\chi$ .

At this stage, all that remains is the calculation of the partial derivatives

$$\frac{\partial \mathcal{Y}_\alpha}{\partial \phi}, \quad \frac{\partial \mathcal{Y}_\alpha}{\partial \theta} \quad \text{and} \quad \frac{\partial \mathcal{Y}_\alpha}{\partial \chi}.$$

In view of formula (15), the first two ones can be easily expressed in terms of derivatives of spherical harmonics on  $\mathbb{S}^2$ . For example, if  $\alpha = (k, \ell, m)$  then

$$(63) \quad \begin{aligned} \frac{\partial \mathcal{Y}_\alpha}{\partial \phi}(\xi) &= \frac{1}{\sqrt{a_{k,\ell}}} (\sin \chi)^\ell T_{k+1}^{(\ell+1)}(\cos \chi) \frac{\partial Y_{\ell,m}}{\partial \phi}(\phi, \theta) \\ &= -\frac{m}{\sqrt{a_{k,\ell}}} (\sin \chi)^\ell T_{k+1}^{(\ell+1)}(\cos \chi) Y_{\ell,-m}(\phi, \theta) \end{aligned}$$

In order to avoid division by zero in (62) (when  $\sin \theta = 0$ ), which is useless, one can employ in the definition (16) of  $Y_{\ell,-m}$  the recurrence property on associated Legendre functions:

$$(64) \quad 2mK_\ell^m(\cos \theta) = \sin \theta (\tau_{\ell,m} K_{\ell+1}^{m+1}(\cos \theta) + \tau_{\ell,-m} K_{\ell+1}^{m-1}(\cos \theta)),$$

with

$$\tau_{\ell,m} = \sqrt{(\ell + m + 2)(\ell + m + 1)} \quad \text{for} \quad -\ell \leq m \leq \ell.$$

Thus, for  $m \neq 0$  we have

$$(65) \quad \frac{2|m|}{\sin \theta} \frac{\partial \mathcal{Y}_\alpha}{\partial \phi}(\xi) = -m \frac{\eta_\ell}{\sqrt{a_{k,\ell}}} (\sin \chi)^\ell T_{k+1}^{(\ell+1)}(\cos \chi) y_{-m}(\phi) \left( \tau_{\ell,|m|} K_{\ell+1}^{|m|+1}(\cos \theta) + \tau_{\ell,-|m|} K_{\ell+1}^{|m|-1}(\cos \theta) \right)$$

Similarly, we have

$$(66) \quad \frac{\partial \mathcal{Y}_\alpha}{\partial \phi}(\xi) = \frac{\eta_\ell}{\sqrt{a_{k,\ell}}} (\sin \chi)^\ell T_{k+1}^{(\ell+1)}(\cos \chi) \left( c_{\ell,-|m|} K_\ell^{|m|-1}(\cos \theta) - c_{\ell,|m|} K_\ell^{|m|+1}(\cos \theta) \right) y_m(\phi),$$

where

$$(67) \quad c_{\ell,m} = \frac{1}{2} \sqrt{(\ell-m)(\ell+m+1)} \text{ for } \ell \geq 0 \text{ and } -\ell \leq m \leq \ell.$$

Note that we used the following recurrence formula:

$$(68) \quad (\sin \theta)(K_\ell^m)'(\cos \theta) = c_{\ell,m} K_\ell^{m+1}(\cos \theta) - c_{\ell,-m} K_\ell^{m-1}(\cos \theta).$$

(and with the convention  $K_\ell^j = 0$  when  $|j| > \ell$ ). Hence

$$(69) \quad \frac{\partial \mathcal{Y}_\alpha}{\partial \theta}(\xi) = \frac{\eta_\ell}{\sqrt{a_{k,\ell}}} (\sin \chi)^\ell T_{k+1}^{(\ell+1)}(\cos \chi) \left( c_{\ell,-|m|} K_\ell^{|m|-1}(\cos \theta) - c_{\ell,|m|} K_\ell^{|m|+1}(\cos \theta) \right) y_m(\phi).$$

Finally, we also have

$$(70) \quad \frac{\partial \mathcal{Y}_\alpha}{\partial \chi}(\xi) = \frac{(\sin \chi)^{\ell-1}}{\sqrt{a_{k,\ell}}} Y_{\ell,m}(\phi, \theta), \left( \ell \cos(\chi) T_{k+1}^{(\ell+1)}(\cos \chi) - (\sin \chi)^2 T_{k+1}^{(\ell+2)}(\cos \chi) \right),$$

for all  $\alpha = (k, \ell, m) \in \Lambda$ .

By using these expressions of partial derivative of functions ( $\mathcal{Y}_\alpha$ ) in (61), we obtain a complete formula which is readily available for practical use and for implementation.

**5.2. Computational tests and numerical validation.** In this section, focus is on some numerical results that allow to assess the practical usability of the formula (33) and (34) and the performances of the resulting numerical method outlined in section 3. Three different examples are investigated in the following. In the first example we deal with non homogeneously magnetized spherical domain for which we have an error estimate by Theorem 4.4. In the two last examples, the domain is homogeneously magnetized. In all these three cases, we derive expressions of the exact stray field, to which the numerical solution is compared. In all these computational tests we set  $\mu_0 = 1$ .

**Example 1: a non homogeneously magnetized sphere with  $M.n = 0$ .**

We prefer starting numerical experiences with the case of a non homogeneously magnetized spherical sample, that is

$$\Omega = \{x \in \mathbb{R}^3 \mid |x| < r_0\}$$

and

$$(71) \quad M = (\cos \theta)e_\varphi + (\sin \theta)e_\theta \text{ in } \Omega.$$



N	$\mathcal{E}_{sf}(u)$	$\mathcal{E}_{sf}(u_N)$	$\frac{ \mathcal{E}_{sf}(u) - \mathcal{E}_{sf}(u_N) }{\mathcal{E}_{sf}(u)}$	$e_0(H_d)$
10	0.07757018	0.07696625	7.78E-3	8.82E-2
20	-	0.07750001	9.03E-4	3.00E-2
30	-	0.07754315	3.48E-4	1.86E-2
40	-	0.07756016	1.29E-4	1.13E-2
50	-	0.07756414	7.79E-5	8.82E-3
60	-	0.07756708	4.00E-5	6.32E-3
The log. slope			-2.90	-1.45

TABLE 1. the exact and the approximate stray-field energy due to a non homogeneously magnetized sphere (example 1).

It may be noted that  $M$  is complying with Heisenberg-Weiss constraint (3) since  $|M| = 1$  in  $\Omega$ . Besides,  $M$  is tangential on the boundary of  $\Omega$  since  $M.n = 0$  on  $\partial\Omega$  (here  $n[x] = x/|x|$ ).

We are able to give an analytical expression of the exact solution (see [24, 13]). More precisely,

$$(72) \quad U(x) = \begin{cases} -\frac{2z}{9} + \frac{2z}{3} \ln\left(\frac{|x|}{r_0}\right) & \text{if } |x| \leq r_0, \\ -\frac{2r_0^3 z}{9|x|^3} & \text{if } |x| \geq r_0. \end{cases}$$

The exact stray-field energy is given by

$$(73) \quad \mathcal{E}_{sf}(U) = \frac{\mu_0}{2} \int_{\mathbb{R}^3} |\nabla U|^2 dx = \frac{16}{81} \pi r_0^3.$$

Here we choose  $r_0 = 1/2$ . In Table 1 we outline the computed stray-field energy (40) for several values of  $N$  (considered as a discretization parameter). We also outline the relative  $L^2$  error on the stray field  $H_d = -\nabla U$  defined by

$$e_0(H_d) = \frac{|U_N - U|_{W_0^1(\mathbb{R}^3)}}{|U|_{W_0^1(\mathbb{R}^3)}}.$$

We can then observe that the error  $e_0(H_d)$  decreases in as  $N^{-1.45}$ . This is in accordance with Proposition 4.3 in which it is forecasted that

$$|U - U_N|_{W_0^1(\mathbb{R}^3)} \leq CN^{-1} \|\operatorname{div} M\|_{L^2(\Omega)}.$$

Actually, the solution  $u$  belongs  $W_2^2(\mathbb{R}^3)$  since  $\operatorname{div} M \in L^2(\Omega)$  and  $M.n = 0$  on  $\partial\Omega$ . There is even a superconvergence with respect to this estimate. Note also that the error on the stray field energy decreases as  $N^{-2.90}$  (in agreement with the identity  $|\mathcal{E}_{sf}(U) - \mathcal{E}_{sf}(U_N)| = |U - U_N|_{W_0^1(\mathbb{R}^3)}^2$ ).

N	$\mathcal{E}_{sf}(u)$	$\mathcal{E}_{sf}(u_N)$	$\frac{ \mathcal{E}_{sf}(u) - \mathcal{E}_{sf}(u_N) }{\mathcal{E}_{sf}(u)}$	$e_0(H_d)$
10	0.08726646	0.07845252	10.10E-2	0.3153
20		0.08252939	5.42E-2	0.2322
30	-	0.08402011	3.72E-2	0.1924
40	-	0.08479348	2.83E-2	0.1680
50	-	0.08526692	2.29E-2	0.1511
60	-	0.08558669	1.92E-2	0.1385
The log. slope			-0.93	-0.46

TABLE 2. The exact and the approximate stray-field energy due to an homogeneously magnetized sphere (example 2).

**Example 2: a homogeneously magnetized sphere.**

In this second benchmark test, we consider a spherical sample  $\Omega = \{x \in \mathbb{R}^3 \mid |x| < r_0\}$  with a constant magnetization  $M = M_0$ . It is easy to prove that the exact solution of (6) is given by the formula:

$$(74) \quad U(x) = \begin{cases} \frac{1}{3}M_0 \cdot x & \text{if } |x| < r_0, \\ \frac{r_0^3}{3} \frac{M_0 \cdot x}{|x|^3} & \text{if } |x| \geq r_0. \end{cases}$$

The exact energy is

$$(75) \quad \mathcal{E}_{sf}(U) = \frac{1}{2} \int_{\mathbb{R}^3} |\nabla U|^2 dx = -\frac{1}{2} \int_{\Omega} M_0 \cdot h dx = \frac{2\pi|M_0|^2}{9} r_0^3.$$

Here, we choose  $M_0 = (0, 0, 1)$  and  $r_0 = 0.5$ . Thus,

$$\mathcal{E}_{sf}(u) = \frac{\pi}{36} = 0.08726646$$

It may be observed that  $[\frac{\partial U}{\partial n}] = -M_0 \cdot n \neq 0$  on  $\partial\Omega$ . Thus,  $U \notin W_2^2(\mathbb{R}^3)$  although  $U|_{\Omega} \in H^2(\Omega)$  and  $U|_{\mathbb{R}^3 \setminus \bar{\Omega}} \in W_2^2(\mathbb{R}^3 \setminus \bar{\Omega})$  (here  $U|_{\Omega}$  and  $U|_{\mathbb{R}^3 \setminus \bar{\Omega}}$  designate the restrictions of  $U$  to  $\Omega$  and to  $\mathbb{R}^3 \setminus \bar{\Omega}$  respectively). We are therefore not within the validity assumptions of Theorem 4.4 and the error estimates (52) and (53) are no longer necessarily true.

In Table 2, the approximate energy  $\mathcal{E}_{sf}(U_N)$  is given for several values of the discretization parameter  $N$ . We also compute the relative  $L^2$  error on the stray field  $H_d = -\nabla U$ . One can observe that this error decreases as  $N^{-0.46}$ . The error on the energy decreases as  $N^{-0.93}$ .

Here again, convergence of the approximate solution to the exact one holds although the normal component of  $h = -\nabla U$  is not continuous across the boundary of the sample.

N	$\mathcal{E}_{sf}(u)$	$\mathcal{E}_{sf}(u_N)$	$\frac{ \mathcal{E}_{sf}(u) - \mathcal{E}_{sf}(u_N) }{\mathcal{E}_{sf}(u)}$	$e_0(H_d)$
10	0.16666666	0.14711046	0.1173	0.3397
20	-	0.15617466	6.3E-2	0.2499
30	-	0.15951131	4.2E-2	0.2066
40	-	0.16123614	3.2E-2	0.180
50	-	0.16229007	2.62E-2	0.1618
60	-	0.16300181	2.19E-2	0.1481
The log. slope			-0.94	-0.47

TABLE 3. The exact and the approximate stray-field energy due to an homogeneously magnetized cube (example 3).

**Example 3: homogeneously magnetized cube.**

In this last test, we change the geometry of the sample and we consider a homogeneously magnetized cubic rod  $\Omega = ]-\gamma, \gamma[^3$ , with  $\gamma = 1/2$ , and  $M = (0, 1, 0)$ . The stray-field energy in this case is (see, e. g., [1])

$$(76) \quad \mathcal{E}_{sf}(U) = \frac{1}{6}.$$

The exact analytical expression of the demagnetizing field is (see [18]):

$$\begin{aligned} H_d(x) &= \frac{1}{4\pi} \left( \sum_{k,\ell,m=1}^2 (-1)^{k+\ell+m} \ln(z + (-1)^m \gamma + \varrho) \right) e_x, \\ &- \frac{1}{4\pi} \left( \sum_{k,\ell,m=1}^2 (-1)^{k+\ell+m} \arctan\left(\frac{(x + (-1)^k \gamma)(z + (-1)^m \gamma)}{(y + (-1)^\ell \gamma) \varrho}\right) \right) e_y, \\ &+ \frac{1}{4\pi} \left( \sum_{k,\ell,m=1}^2 (-1)^{k+\ell+m} \ln(x + (-1)^k \gamma + \varrho) \right) e_z, \end{aligned}$$

where  $\varrho = \sqrt{(x + (-1)^k \gamma)^2 + (y + (-1)^\ell \gamma)^2 + (z + (-1)^m \gamma)^2}$ .

It may be observed that  $M.n \neq 0$  on  $\partial\Omega$ . Thus,  $U \notin W_2^2(\mathbb{R}^3)$  (see Remark 4.2). The numerical results summarized in Table 3 confirm the convergence of the method and show that here too the  $L^2$  error on the stray field  $H_d$  decreases as  $N^{-0.47}$ . The error on the energy decreases as  $N^{-0.93}$ , while the error on the energy decreases like  $N^{-0.97}$ .

## 6. CONCLUSION AND PERSPECTIVES

The formula (33), in addition to being original, has several advantages both theoretically and numerically. From a computational point of view, it has been established that the formula inspires a particularly efficient and easy to implement numerical method to calculate the demagnetizing field and the associated energy. Indeed, the numerical results show a rapid convergence of the method especially when  $M.n = 0$  on  $\partial\Omega$ . In the latter case, the observed convergence

is even faster than that predicted by the error estimate in Theorem 4.4 since the convergence in energy is of order close to  $O(\frac{1}{N^3})$ . This suggests that these estimates are not optimal and could possibly be improved theoretically. In the case  $M.n \neq 0$ , the method also converges in accordance with the lemma, but one notes that convergence of the energy is of order close to  $O(\frac{1}{N})$ . This fact remains to be proven theoretically.

From a theoretical point of view, one could exploit formula (33) to give a new expression to the functional to be minimized. Actually, the total free energy can be expressed as:

$$(77) \quad E_{tot}(M) = \alpha \int_{\Omega} |\nabla M|^2 dx + \int_{\Omega} \phi(M) dx - \mu_0 \int_{\Omega} H_{ex} \cdot M dx + \sum_{k=0}^{\infty} \frac{2\mu_0}{4(k+1)^2 - 1} \sum_{\alpha \in \Lambda_k} \left( \int_{\Omega} M \cdot \nabla \mathcal{W}_{\alpha} dx \right)^2 + E_s.$$

It is well known that the minimization of the functional  $E_{tot}$  with respect to the variable  $M$  under Heisenberg-Weiss constraint (3) leads to the following partial differential equation (see, e. g., [23] and references therein):

$$(78) \quad -2\alpha \Delta M + \nabla_M \phi(M) - \mu_0 (H_d + H_{ext}) = \lambda M \text{ in } \Omega,$$

where  $\lambda$  is a lagrangian multiplier. By sake of simplicity we assumed here that  $E_s = 0$  (the reader can refer to, e. g., [23] for the general equations taking into account this term).

Formula (33) simplifies the system (78) and reduces it to only one equation

$$(79) \quad -2\alpha \Delta M + \nabla_M \phi(M) - \mu_0 H_{ext} + \mu_0 \sum_{k=0}^{\infty} \sum_{\alpha \in \Lambda_k} \frac{2}{4(k+1)^2 - 1} \left( \int_{\Omega} M \cdot \nabla \mathcal{W}_{\alpha} dx \right) \nabla \mathcal{W}_{\alpha} = \lambda M \text{ in } \Omega.$$

The study of this non-local PDE could provide new information about the best configuration minimizing the functional  $E_{tot}$ . If we truncate the serie on the left-hand side, keeping only the first term, we obtain the simplified approximate non local equation

$$(80) \quad -2\alpha \Delta M + \nabla_M \phi(M) - \mu_0 H_{ext} + \frac{2\mu_0}{3\pi^2(|x|^2 + 1)^{3/2}} \left( \int_{\Omega} \frac{M \cdot x}{(|x|^2 + 1)^{3/2}} dx \right) x = \lambda M \text{ in } \Omega.$$

The study of this kind of equations is beyond the scope of this paper; it will be the subject of a forthcoming paper.

**A. Proof of Proposition 5.1.** The objective here is to prove formula (61). Let  $Y$  be an arbitrary smooth function defined on  $\mathbb{S}^3$  and set

$$W(x) = \left( \frac{2}{|x|^2 + 1} \right)^{1/2} Y(\pi^{-1}(x)),$$

(thus, if  $Y = \mathcal{Y}_\alpha$ ,  $\alpha \in \Lambda$ , then  $W = \mathcal{W}_\alpha$ ). In [5] and [11] (formula A.9), the authors prove the following identity (linking the gradient of  $W$  to  $Y$  and its tangential derivatives on the unit sphere):

$$(81) \quad \nabla W(x) = (1 - \xi_4)^{1/2} \left( S(\xi) \nabla_\xi Y(\xi) - \frac{1}{2} Y(\xi) \hat{\xi} \right), \text{ for } x \in \mathbb{R}^3,$$

where  $\xi = \pi^{-1}(x) \in \mathbb{S}^3$ ,  $\hat{\xi} = (\xi_1, \xi_2, \xi_3)$  is the orthogonal projection of  $\xi = (\xi_1, \xi_2, \xi_3, \xi_4)$  on  $\mathbb{R}^3$ ,  $\nabla_\xi Y$  is the tangential gradient of  $Y$  on  $\mathbb{S}^3$  and  $S(\xi)$  is the  $3 \times 4$  rectangular matrix

$$(82) \quad \begin{aligned} S(\xi) &= \begin{pmatrix} 1 - \xi_4 & 0 & 0 & \xi_1 \\ 0 & 1 - \xi_4 & 0 & \xi_2 \\ 0 & 0 & 1 - \xi_4 & \xi_3 \end{pmatrix} \\ &= \begin{pmatrix} 1 - \cos \chi & 0 & 0 & \cos \phi \sin \theta \sin \chi \\ 0 & 1 - \cos \chi & 0 & \sin \phi \sin \theta \sin \chi \\ 0 & 0 & 1 - \cos \chi & \cos \theta \sin \chi \end{pmatrix}. \end{aligned}$$

It remains to spell out the expression of the tangential gradient  $\nabla_\xi Y(\xi)$  in terms of partial derivatives of  $Y$  with respect to  $\phi$ ,  $\theta$  and  $\chi$ , the spherical coordinates of  $\xi$  (see section 2). We state this

**Lemma .1.** *If  $\sin \chi \neq 0$ , then*

$$(83) \quad \nabla_\xi Y(\xi) = \frac{1}{\sin \chi} \begin{pmatrix} -\sin \phi & \cos \phi \cos \theta & \cos \phi \sin \theta \cos \chi \\ \cos \phi & \sin \phi \cos \theta & \sin \phi \sin \theta \cos \chi \\ 0 & -\sin \theta & \cos \theta \cos \chi \\ 0 & 0 & -\sin \chi \end{pmatrix} \begin{pmatrix} \frac{1}{\sin \theta} \frac{\partial \tilde{Y}}{\partial \phi}(\phi, \theta, \chi) \\ \frac{\partial \tilde{Y}}{\partial \theta}(\phi, \theta, \chi) \\ \sin \chi \frac{\partial \tilde{Y}}{\partial \chi}(\phi, \theta, \chi) \end{pmatrix}.$$

where  $\tilde{Y}(\phi, \theta, \chi) = Y(\cos \phi \sin \theta \sin \chi, \sin \phi \sin \theta \sin \chi, \cos \theta \sin \chi, \cos \chi)$ .

*Proof.* Consider the 0-homogeneous function  $F$  defined over  $\mathbb{R}^4 \setminus \{0\}$  by

$$F(y) = Y\left(\frac{y}{|y|}\right), \quad y \in \mathbb{R}^4 \setminus \{0\}.$$

It follows that

$$(84) \quad \nabla_\xi Y(\xi) = \nabla F(\xi) \text{ for } \xi \in \mathbb{S}^3.$$

In view of Euler's homogeneous function lemma, we have

$$(85) \quad \sum_{i=1}^4 y_i \frac{\partial F}{\partial y_i}(y) = 0 \text{ for } y \in \mathbb{R}^4 \setminus \{0\}.$$

Since

$$\tilde{Y}(\phi, \theta, \chi) = F(\cos \phi \sin \theta \sin \chi, \sin \phi \sin \theta \sin \chi, \cos \theta \sin \chi, \cos \chi),$$

we deduce that

$$\begin{aligned}\frac{\partial \tilde{Y}}{\partial \phi}(\phi, \theta, \chi) &= \sin \theta \sin \chi \left( -\sin \phi \frac{\partial F}{\partial y_1}(y) + \cos \phi \frac{\partial F}{\partial y_2}(y) \right), \\ \frac{\partial \tilde{Y}}{\partial \theta}(\phi, \theta, \chi) &= \cos \theta \sin \chi \left( \cos \phi \frac{\partial F}{\partial y_1}(y) + \sin \phi \frac{\partial F}{\partial y_2}(y) \right) \\ &\quad - \sin \theta \sin \chi \frac{\partial F}{\partial y_3}(y), \\ \frac{\partial \tilde{Y}}{\partial \chi}(\phi, \theta, \chi) &= \sin \theta \cos \chi \left( \cos \phi \frac{\partial F}{\partial y_1}(y) + \sin \phi \frac{\partial F}{\partial y_2}(y) \right) \\ &\quad + \cos \theta \cos \chi \frac{\partial F}{\partial y_3}(y) - \sin \chi \frac{\partial F}{\partial y_4}(y),\end{aligned}$$

where  $y = (\cos \phi \sin \theta \sin \chi, \sin \phi \sin \theta \sin \chi, \cos \theta \sin \chi, \cos \chi)$ .

Completing these identities with equation (85) gives a square linear system in terms of the derivatives  $\frac{\partial F}{\partial y_i}(y)$ ,  $1 \leq i \leq 4$ . Inverting this system gives

$$\nabla F(y) = R(\xi) \begin{pmatrix} \frac{\partial \tilde{Y}}{\partial \phi} \\ \frac{\partial \tilde{Y}}{\partial \theta} \\ \frac{\partial \tilde{Y}}{\partial \chi} \\ 0, \end{pmatrix}$$

where

$$R(\xi) = \begin{pmatrix} -\frac{\sin \phi}{\sin \theta \sin \chi} & \frac{\cos \phi \cos \theta}{\sin \chi} & \cos \phi \sin \theta \cos \chi & \cos \phi \sin \theta \sin \chi \\ \frac{\cos \phi}{\sin \theta \sin \chi} & \frac{\sin \phi \cos \theta}{\sin \chi} & \sin \phi \sin \theta \cos \chi & \sin \phi \sin \theta \sin \chi \\ 0 & -\frac{\sin \theta}{\sin \chi} & \cos \theta \cos \chi & \cos \theta \sin \chi \\ 0 & 0 & -\sin \chi & \cos \chi \end{pmatrix}$$

This ends the proof of (83). Formula (61) is a direct consequence of (81) and (83).  $\square$

### B. The first few three-dimensional Arar-Boulmezaoud functions.

In this appendix we an give explicit formulas of the first few three-dimensional Arar-Boulmezaoud functions defined by (22). These functions are illustrated in Table 6.

#### Declarations.

*Conflict of interest:* The author declares no competing interests.

$k$	$\ell$	$m$	$\mathcal{W}_{(k,\ell,m)}(x)$
0	0	0	$\frac{1}{\pi( x ^2 + 1)^{1/2}}$
1	0	0	$\frac{2}{\pi} \frac{ x ^2 - 1}{( x ^2 + 1)^{3/2}}$
	1	0	$\frac{4}{\pi} \frac{x_3}{( x ^2 + 1)^{3/2}}$
		1	$\frac{4}{\pi} \frac{x_1}{( x ^2 + 1)^{3/2}}$
		-1	$\frac{4}{\pi} \frac{x_2}{( x ^2 + 1)^{3/2}}$
2	0	0	$\frac{1}{\pi} \frac{3 x ^4 - 10 x ^2 + 3}{( x ^2 + 1)^{5/2}}$
	1	0	$\frac{4\sqrt{6}}{\pi} \frac{x_3( x ^2 - 1)}{( x ^2 + 1)^{5/2}}$
		1	$\frac{4\sqrt{6}}{\pi} \frac{x_1( x ^2 - 1)}{( x ^2 + 1)^{5/2}}$
		-1	$\frac{4\sqrt{6}}{\pi} \frac{x_2( x ^2 - 1)}{( x ^2 + 1)^{5/2}}$
	2	0	$\frac{4\sqrt{2}}{\pi} \frac{3x_3^2 -  x ^2}{( x ^2 + 1)^{5/2}}$
		1	$\frac{8\sqrt{6}}{\pi} \frac{x_1 x_3}{( x ^2 + 1)^{5/2}}$
		2	$\frac{4\sqrt{6}}{\pi} \frac{x_1^2 - x_2^2}{( x ^2 + 1)^{5/2}}$
		-1	$\frac{8\sqrt{6}}{\pi} \frac{x_2 x_3}{( x ^2 + 1)^{5/2}}$
		-2	$\frac{8\sqrt{6}}{\pi} \frac{x_1 x_2}{( x ^2 + 1)^{5/2}}$

TABLE 4. Explicit expressions of the first Arar-Boulmezaoud functions in  $\mathbb{R}^3$

#### REFERENCES

- [1] C. Abert, L. Exl, G. Selke, A. Drews, and Th. Schrefl. Numerical methods for the stray-field calculation: A comparison of recently developed algorithms. *Journal of magnetism and magnetic materials*, 326(326):176–185, 2013.
- [2] C. Abert, Gunnar Selke G., Kruger B., and A. Drews. A fast finite-difference method for micromagnetics using the magnetic scalar potential. *IEEE transactions on magnetics*, 38(3):1105–1109, 2011.
- [3] F. Alliot. *Etude des équations stationnaires de Stokes et Navier-Stokes dans des domaines extérieurs*. PhD Thesis, ENPC, Paris, 1998.
- [4] C. Amrouche, V. Girault, and J. Giroire. Weighted Sobolev spaces for Laplace’s equation in  $\mathbf{R}^n$ . *J. Math. Pures Appl. (9)*, 73(6):579–606, 1994.

- [5] N. Arar and T. Z. Boulmezaoud. Eigenfunctions of a weighted Laplace operator in the whole space. *J. Math. Anal. Appl.*, 400(1):161–173, 2013.
- [6] M. Aurada, J.M. Melenk, and Dirk Praetorius. Fem-bem coupling for the large-body limit in micromagnetics. *Journal of Computational and Applied Mathematics*, 281:10–31, 06 2015.
- [7] D. V. Berkov, K. Ramstöck, and A. Hubert. Solving micromagnetic problems: toward and optimal numerical method. *Phys. Stat. Sol (a)*, 137:207–225, 1993.
- [8] S. K. Bhowmik, R. Belbaki, Boulmezaoud T. Z., and S. Mziou. Solving two dimensional second order elliptic equations in exterior domains using the inverted finite elements method. *Computers & Mathematics with Applications*, 72(9):2315–2333, 2016.
- [9] J. L. Blue and M. R. Scheinfein. Using multipoles decreases computation time for magnetic self-energy. *IEEE Trans. Magn.*, 27:4778–4780, 1991.
- [10] T. Z. Boulmezaoud. Inverted finite elements: a new method for solving elliptic problems in unbounded domains. *M2AN Math. Model. Numer. Anal.*, 39(1):109–145, 2005.
- [11] T. Z. Boulmezaoud, N. Arar, N. Kerdid, and A. Kourta. Discretization by rational and quasi-rational functions of multi-dimensional elliptic problems in the whole space. *ESAIM Math. Model. Numer. Anal.*, 50(1):263–288, 2016.
- [12] T. Z. Boulmezaoud, Kaliche K., and N. Kerdid. Inverted finite elements for div-curl systems in the whole space. *Advances in Computational Mathematics*, 43(6):1469–1489, 2017.
- [13] T. Z. Boulmezaoud and K. Kaliche. Stray field computation by inverted finite elements: a new method in micromagnetic simulations. *arXiv:2301.10544*, 2023.
- [14] T. Z. Boulmezaoud, S. Mziou, B. Boudjedaa, and M. M. Babatin. Inverted finite elements for degenerate and radial elliptic problems in unbounded domains. *Japan Journal of Industrial and Applied Mathematics*, 32(1):237–261, 2015.
- [15] T. Z. Boulmezaoud, S. Mziou, and T. Boudjedaa. Numerical approximation of second-order elliptic problems in unbounded domains. *Journal of Scientific Computing*, 60(2):295–312, 2014.
- [16] W. F. Brown. *Micromagnetics. Interscience*. Interscience Publishers, 1963.
- [17] C. Carstensen and A. Prohl. Numerical analysis of relaxed micromagnetics by penalised finite elements. *Numer. Math.*, 90(1):65–99, 2001.
- [18] R. Engel-Herbert and T. Hesjedal. Calculation of the magnetic stray field of a uniaxial magnetic domain. *Journal of Applied Physics*, 97:074504–074504, 05 2005.
- [19] L. Exl, W. Auzinger, S. Bance, M. Gusenbauer, F. Reichel, and T. Schrefl. Fast stray field computation on tensor grids. *J. Comput. Phys.*, 231(7):2840–2850, 2012.
- [20] D. R. Fredkin and T. R. Koehler. Hybrid method for computing demagnetizing fields. *IEEE Trans. Magn.*, 26:415–417, 1990.
- [21] J. J. Friess, S. S. Gubser, G. Michalogiorgakis, and S. S. Pufu. Expanding plasmas and quasinormal modes of anti-de sitter black holes. *Journal of High Energy Physics*, (4), 2007.
- [22] V. Girault and P. A. Raviart. *Finite element methods for Navier-Stokes equations*, volume 5 of *Springer Series in Computational Mathematics*. Springer-Verlag, Berlin, 1986.
- [23] A. Hubert and R. Schäfer. *Magnetic domains: the analysis of magnetic microstructures*. Springer, 1998.
- [24] K. Kaliche. *Méthode des éléments finis inversés pour des domaines non bornés*. PhD Thesis, Paris-Saclay University, Paris, 2016.
- [25] T. R. Koehler and D. R. Fredkin. Finite element methods for micromagnetism. *IEEE Trans. Magn.*, 28:1239–1244, 1992.
- [26] M. Kravčik and A. Prohl. Recent developments in the modeling, analysis, and numerics of ferromagnetism. *Numer. Math.*, 48(3):439–483, 2006.
- [27] S. Labbé. Fast computation for large magnetostatic systems adapted for micromagnetism. *SIAM J. Sci. Comput.*, 26(6):2160–2175, 2005.
- [28] L. D. Landau and E. M. Lifschitz. On the theory of the dispersion of magnetic permeability in ferromagnetic bodies. *Phys. Z. Sowjetunion*, 8(8):153–69, 1935.



- [29] H. Long, E. Ong, Z. Liu, and E. Li. Fast fourier transform on multipoles for rapid calculation of magnetostatic fields. *IEEE Trans. Magn.*, 42:295–300, 2006.
- [30] T. Miyazaki and H. Jin. *The Physics of Ferromagnetism*. Springer, 2012.
- [31] C. Müller. *Spherical harmonics*, volume 17 of *Lecture Notes in Mathematics*. Springer-Verlag, Berlin, 1966.
- [32] N. Popović and D. Praetorius. Applications of  $H$ -matrix techniques in micromagnetics. *Computing*, 74(3):177–204, 2005.
- [33] A. Prohl. *Computational micromagnetism*. Advances in Numerical Mathematics. B. G. Teubner, Stuttgart, 2001.
- [34] R. T. Seeley. Spherical harmonics. *Amer. Math. Monthly*, 73(4, part II):115–121, 1966.
- [35] K. Tomita. Tensor spherical and pseudospherical harmonics in four-dimensional spaces. *RRK*, 82(3), 1982.
- [36] G. F. Torres del Castillo and J. L. Calvario Acócal. The hydrogen atom via the four-dimensional spherical harmonics. *Revista Mexicana de física*, 53(5):407–414, 2007.
- [37] J.C. Toussaint, A. Marty, N. Vukadinovic, and M. Youssef, J. B. and Labrune. A new technique for ferromagnetic resonance calculations. *Computational Materials Sciences*, 24:175–180, 2002.
- [38] A. Vansteenkiste, J. Leliaert, M. Dvornik, M. Helsen, F. Garcia-Sanchez, and B. Van Waeyenberge. The design and verification of mumax3. *AIP ADVANCES*, 4(10):22, 2014.