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Point vortex dynamics as zero-radius limit of the motion of a rigid body in an irrotational fluid

Olivier Glass · Alexandre Munnier ·
Franck Sueur

April 13, 2016

Abstract The vortex point system is usually considered as an idealized model where the vorticity of an ideal incompressible two-dimensional fluid is concentrated in a finite number of moving points.

In the case of a single vortex in an otherwise irrotational ideal fluid occupying a bounded and simply-connected two-dimensional domain the motion is given by the so-called Kirchhoff-Routh velocity which depends only on the domain. The main result of this paper establishes that this dynamics can also be obtained as the limit of the motion of a rigid body immersed in such a fluid when the body shrinks to a pointwise massless particle with fixed circulation. The rigid body is assumed to be only accelerated by the force exerted by the fluid pressure on its boundary, the fluid velocity and pressure being given by the incompressible Euler equations, with vanishing vorticity. The circulation of the fluid velocity around the particle is conserved as time proceeds according to Kelvin's theorem and gives the strength of the limit point vortex.

In order to understand the solid dynamics one has to evaluate the pressure field on the boundary of the solid, that is, where the singularity is concentrated at the limit. Moreover the Newton equations driving the particle's dynamics involve a singular perturbation problem in time whereas the fluid state may be seen as solving an auxiliary problem. The fluid velocity can indeed be recovered by an elliptic-type problem where time appears only as a parameter.

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Since the analysis involves singular perturbation problem both in space and in time and that nothing excludes a priori some sharp energy transfer from the fluid toward the rigid immersed particle it is crucial to recast the solid dynamics under normal forms. This form makes appear the gyroscopic structure of these singularities and allows one to obtain uniform estimates on the dynamics thanks to some energy-type quantities modulated by the limit dynamics, and therefore to describe the transition of the dynamics in the limit. In order to get such normal forms we first establish that the Newton equations can be seen as a geodesic equation, with a metric associated with the well-known “added inertia” phenomenon, under an applied force similar to the Lorentz force which can be seen as an extension of the celebrated Kutta-Joukowski lift force. Then, in the limit, surprising relations between leading and subprincipal orders of various terms and the modulation variables show up and allow us to establish a normal form with a more gyroscopic structure.

As a byproduct of our analysis we also prove that in the different regime where the body shrinks with a fixed mass the limit equation is a second-order differential equation involving a Kutta-Joukowski-type lift force, which extends the result of [9] to the case where the domain occupied by the solid-fluid system is bounded.

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1 Introduction

The point vortex system is a classical topic which originates from fluid mechanics and goes back to Helmholtz [14], Kirchhoff [20], Poincaré [33], Routh [36], Kelvin [19], and Lin [22, 23]. It appeared as an idealized model where the vorticity of an ideal incompressible two-dimensional fluid is concentrated in a finite number of points. Although it does not constitute a solution of the Euler equations in the sense of distributions, it is now understood that point vortices can be viewed as limits of concentrated smooth vortices which evolve according to the Euler equations. In the case of a single vortex moving in a bounded and simply-connected domain this was proved by Turkington in [38]. An extension to the case of several vortices was given by Marchioro and Pulvirenti, see [25]. Recently Gallay has proven in [8] that the point vortex system can also be obtained as vanishing viscosity limits of concentrated smooth vortices evolving according to the incompressible Navier-Stokes equations.

The main goal of this paper is to prove that the point vortex system can also be viewed as the limit of the dynamics of a solid, shrinking into a pointwise massless particle with fixed circulation, in free motion in an irrotational fluid. By free motion we mean that the rigid body is only accelerated by the force exerted by the fluid pressure on its boundary, the fluid velocity and pressure being given by the incompressible Euler equations, with vanishing vorticity. In a different regime, we also derive a different “massive” point vortex system. Let us mention that in [7, Chapter 3] Friedrichs already evoked such a massive point vortex system in the case of two point vortices in the whole plane under the terminology of *bound vortices* (as opposed to *free vortices*). Here we perform a rigorous analysis to derive the massive point vortex system in a cavity from the dynamics of a shrinking solid in a fluid and we even obtain the *classical point vortex system* as a limit in a certain inertia regime of the dynamics of a shrinking solid in a fluid.

The limit where the body has a diameter tending to zero corresponds to a singular perturbation problem (in space) for the fluid. The singular feature of the fluid perturbation by the vanishingly small particle is largely due to the circulation around the particle, a quantity which is conserved as time proceeds according to Kelvin’s theorem. The circulation somehow encodes the amount of vorticity hidden in the particle from the fluid viewpoint. The amount of circulation is supposed to be independent of the size of the body in our problem. It is well understood since the work [17], see also [24] that when a solid obstacle with a non vanishing given circulation is held fixed in a perfect incompressible fluid, with possibly non vanishing vorticity, then in the limit the obstacle shrinks into a fixed pointwise particle and the Euler equation driving the fluid evolution has to be modified: in the Biot-Savart law providing the fluid velocity generated by the fluid vorticity, a point vortex placed at the fixed position of the pointwise obstacle has to be added to the fluid vorticity, with an amplitude equal to the circulation previously mentioned.

Still the analysis of the dynamics of immersed rigid particles requires a more precise analysis, in particular because it is driven by the fluid pressure,

a quantity which depends in a non linear and non local way on the fluid velocity. Hence to understand the limit dynamics one has to precisely evaluate the pressure field on the boundary of the solid, that is, where the singularity is concentrated. Moreover the Newton equations driving the particles dynamics involve a singular perturbation problem in time (in addition to the singularity in space), in a particularly intricate way for asymmetric particles (actually for any other form than a disk) and even more so for light particles whose mass and moment of inertia go to zero. It is therefore crucial to understand the structure of these singularities in order to be able to describe the transition of the dynamics in the limit.

In order to deal with small light particles we look for normal forms of the Newton equations tailored to obtain uniform estimates on the dynamics thanks to some energy-type quantities modulated by the limit dynamics. Indeed even if the whole system is hamiltonian, cf. [11], nothing excludes a priori some sharp energy exchange between the fluid and the rigid immersed particles. The development of these normal forms requires to obtain an accurate description of the pressure, which itself requires to obtain asymptotic developments of various functions allowing to describe the fluid. Then the analysis will also require surprising relations between various terms appearing in the force term and the modulation variables.

Let us give more details. The setting that we consider in this paper is the same as Turkington in [38], that is we assume that the fluid is ideal, confined in a two-dimensional bounded domain and we consider the motion of a single solid immersed in it. Moreover the flow is supposed to be irrotational.

After the inspection of this situation one may forge the two following convictions regarding the nature of the fluid-rigid body interaction. On the one hand the rigid body drives the dynamics and the fluid may be seen as solving an auxiliary problem. The fluid velocity can indeed be recovered by an elliptic-type problem where time appears only as a parameter. On the other hand the force undergone by a rigid body highly depends on its position: It is not the same in the neighborhood of the fixed external boundary as in the bulk of the fluid. This is of course due to the incompressibility of the ambient fluid which allows such nonlocal effects.

More precisely the whole system may be recast as an ODE for the three degrees of freedom of the body (two for the translation and one for the rotation). This is the first result of this paper where the precise structure of this ODE is given; it can be seen as a geodesic equation with an applied force similar to the Lorentz force. See Theorem 2.2 below. One can also interpret this ODE as a lagrangian reformulation of the hamiltonian formulation of the whole system mentioned above, but projected in a way that it depends on the degrees of freedom corresponding to the immersed rigid particle only. Of course this cannot be done without altering the structure. Let us underline that this reformulation of the system plays a central role to establish our asymptotic results, when the solid shrinks to zero.

Three remarks on this ODE are in order.

- The metric associated with the geodesic part of the equation is given by the total inertia, that is the inertia of the solid to which one adds the so-called “added inertia”. This added mass effect is natural: a rigid body has to accelerate not only itself but also an amount of the incompressible fluid around it, this extra inertia being encoded into a symmetric nonnegative matrix which depends on the geometry of the whole system.
- The magnetic part of the Lorentz force is a gyroscopic force, proportional to the circulation around the body, which can be seen as an extension of the celebrated Kutta-Joukowski lift force revealed in the case of a single body in a irrotational unbounded flow at the beginning of the 20th century during the first mathematical investigations in aeronautics, see for example [21]. In particular the contribution of this force to the energy vanishes.
- Still the ODE also contains some electric-type force which may seem very damaging in order to obtain uniform estimates and to pass to the limit. The electric term will however have some structure which we can exploit.

Fortunately enough an intricate multi-scale analysis of the fluid state (a precise analysis of the singular perturbation problem given by the vanishingly small presence of the particle in the space domain of the elliptic-type problem satisfied by the fluid velocity) and some striking combinations of apparently unrelated terms (very loosely speaking some parts of the subprincipal order in the asymptotic expansion of two types of term can be nicely combined with the principal order in the expansion of a third type of term) allow to transfer this bad electric-type part of the force term as a modulation of the particle velocity appearing in the other terms of the equation (i.e. the geodesic and magnetic terms).

To summarize, our analysis allows to give a sharp asymptotic analysis of the remote interaction between one moving body (in particular its velocity and circulation) and an exterior boundary, and provides another derivation of the point vortex system as limit of dynamics of a rigid particle immersed in a perfect fluid as its radius and mass go to zero. The vorticity of the limit pointwise particle comes from the circulation (which is held fixed during the shrinking process) around the particle.

Besides, as already mentioned earlier, our analysis also allows to cover another asymptotic regime corresponding to the shrinking of a solid with fixed mass and circulation. In this case we obtain at the limit a second-order differential equation involving a Kutta-Joukowski-type lift force, which extends the result of [9] to the case where the solid-fluid system is bounded.

We are therefore interested in determining the limit of the dynamics of the solid when its size goes to 0, distinguishing two cases:

- Case (i): when the mass of the solid is fixed (and then the solid tends to a massive pointwise particle), and
- Case (ii): when the mass tends to 0 along with the size (and then the solid tends to a massless pointwise particle). This encompasses the case of fixed density.

The main results in this paper, rigorously stated in the following section, are Theorem 2.9 concerning Case (i) and Theorem 2.10 concerning Case (ii).

Case (i) is way easier to tackle and we will therefore present it first in the sequel as a warm-up before the difficulties appearing in the massless case. In particular only a rough asymptotic analysis of the body dynamics suffices in this case.

In the course of our analysis we will thus make use of three normal forms :

- a geodesic-Lorentz normal form for a fixed positive radius,
- a first gyroscopic asymptotic normal form for Case (i),
- a second geodesic-gyroscopic asymptotic normal form for Case (ii), where a modulation of the unknown is used.

2 Main results

2.1 Dynamics of a solid with fixed size and mass

To begin with, let us recall the dynamics of a solid with fixed size and mass in a perfect incompressible fluid. We denote by Ω the bounded open smooth and simply connected¹ domain of \mathbb{R}^2 occupied by the system fluid-solid. At the initial time, the domain of the solid is a non-empty closed smooth and simply connected domain $\mathcal{S}_0 \subset \Omega$ and $\mathcal{F}_0 := \Omega \setminus \mathcal{S}_0$ is the domain of the fluid. There is no loss of generality in assuming that the center of mass of the solid coincides at the initial time with the origin.

The rigid motion of the solid is described at every moment by a rotation matrix

$$R(\vartheta(t)) := \begin{bmatrix} \cos \vartheta(t) & -\sin \vartheta(t) \\ \sin \vartheta(t) & \cos \vartheta(t) \end{bmatrix},$$

describing the rotation of the solid with respect to its original position and a vector $h(t)$ in \mathbb{R}^2 describing the position of the center of mass. The domain of the solid at every time $t > 0$ is therefore $\mathcal{S}(t) := R(\vartheta(t))\mathcal{S}_0 + h(t)$, while the domain of the fluid is $\mathcal{F}(t) := \Omega \setminus \mathcal{S}(t)$ (see Fig. 1). The fluid-solid system is

¹ the simple connectedness is a simplifying assumption but is actually not essential in the analysis.

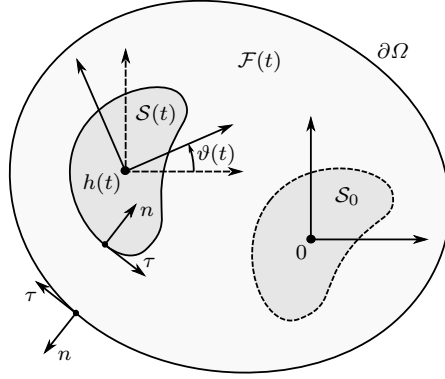


Fig. 1 The domains Ω , $\mathcal{S}(t)$ and $\mathcal{F}(t) := \Omega \setminus \overline{\mathcal{S}(t)}$ of the problem.

governed by the following set of coupled equations:

Fluid equations:

$$\frac{\partial u}{\partial t} + (u \cdot \nabla)u + \nabla \Pi = 0 \quad \text{in } \mathcal{F}(t), \quad (2.1a)$$

$$\operatorname{div} u = 0 \quad \text{in } \mathcal{F}(t). \quad (2.1b)$$

Solid equations:

$$mh'' = \int_{\partial \mathcal{S}(t)} \Pi n \, ds, \quad (2.1c)$$

$$\mathcal{J}\vartheta'' = \int_{\partial \mathcal{S}(t)} (x - h(t))^\perp \cdot \Pi n \, ds. \quad (2.1d)$$

Boundary conditions:

$$u \cdot n = (\vartheta'(\cdot - h)^\perp + h') \cdot n \quad \text{on } \partial \mathcal{S}(t), \quad (2.1e)$$

$$u \cdot n = 0 \quad \text{on } \partial \Omega. \quad (2.1f)$$

Initial data:

$$u|_{t=0} = u_0 \quad \text{in } \mathcal{F}_0, \quad (2.1g)$$

$$\vartheta(0) = 0, h(0) = 0, h'(0) = \ell_0, \vartheta'(0) = \omega_0. \quad (2.1h)$$

Above u and Π denote the velocity and pressure fields in the fluid, $m > 0$ and $\mathcal{J} > 0$ denote respectively the mass and the moment of inertia of the body while the fluid is supposed to be homogeneous of density 1, to simplify the notations. When $x = (x_1, x_2)$ the notation x^\perp stands for $x^\perp = (-x_2, x_1)$, n denotes the unit normal vector pointing outside of the fluid. Let us also emphasize that ds will stand for the arc length without any distinction between $\partial \Omega$, $\partial \mathcal{S}(t)$ and $\partial \mathcal{S}_0$.

Let us recall that if the flow is irrotational at the initial time, that is if $\operatorname{curl} u_0 = 0$ in \mathcal{F}_0 , it will remain irrotational for every time, that is

$$\operatorname{curl} u(t, \cdot) = 0 \text{ in } \mathcal{F}(t), \quad (2.2)$$

according to Helmholtz's third theorem. On the other hand the circulation around the body is constant in time:

$$\int_{\partial\mathcal{S}(t)} u(t) \cdot \tau \, ds = \gamma, \quad (2.3)$$

with

$$\gamma = \int_{\partial\mathcal{S}_0} u_0 \cdot \tau \, ds,$$

according to Kelvin's theorem. Here τ denotes the unit counterclockwise tangential vector so that $n = \tau^\perp$. Let us mention here that we will also use the notation τ on $\partial\Omega$ such that $n := \tau^\perp$ so that it is clockwise in this case (see Fig. 1).

In the irrotational case, the system (2.1) can be recast as an ODE whose unknowns are the degrees of freedom of the solid, namely ϑ and h . In particular, given (2.3), the motion of the fluid is completely determined by the solid position and velocity. In order to state this rigorously, let us introduce the variables

$$h := (h_1, h_2) \quad \text{and} \quad q := (\vartheta, h_1, h_2) \in \mathbb{R}^3.$$

Since the domains $\mathcal{S}(t)$ and $\mathcal{F}(t)$ depend on q only, we will rather denote them $\mathcal{S}(q)$ and $\mathcal{F}(q)$ in the rest of the paper. Since throughout this paper we will not consider any collision, we introduce:

$$\mathcal{Q} := \{q \in \mathbb{R}^3 : d(\mathcal{S}(q), \partial\Omega) > 0\}, \quad (2.4)$$

where $d(A, B)$ denotes the minimal distance between the two sets A and B in the plane

$$d(A, B) := \inf \{|x - y|_{\mathbb{R}^2}, x \in A, y \in B\}. \quad (2.5)$$

Above and all along the paper we will use the notation $|\cdot|_{\mathbb{R}^d}$ for the Euclidean norm in \mathbb{R}^d . Since \mathcal{S}_0 is a closed subset in the open set Ω , the initial position $q(0) = 0$ of the solid belongs to \mathcal{Q} .

Now we need to introduce various objects depending on the geometry and on the constants m, \mathcal{J}, γ , in order to make the aforementioned ODE explicit.

Kirchhoff potentials. Consider the functions ξ_j , for $j = 1, 2, 3$, defined for (q, x) in $\cup_{q \in \mathcal{Q}} (\{q\} \times \mathcal{F}(q))$, by the formula

$$\xi_1(q, x) := (x - h)^\perp \quad \text{and} \quad \xi_j(q, x) := e_{j-1}, \quad \text{for } j = 2, 3, \quad (2.6)$$

where e_1 and e_2 are the unit vectors of the canonical basis. For any $j = 1, 2, 3$, for any q in \mathcal{Q} , we denote by $K_j(q, \cdot)$ the normal trace of ξ_j on $\partial\Omega \cup \partial\mathcal{S}(q)$, that is:

$$K_j(q, \cdot) := n \cdot \xi_j(q, \cdot) \quad \text{on } \partial\Omega \cup \partial\mathcal{S}(q). \quad (2.7)$$

Now fixed q in \mathcal{Q} , we introduce the Kirchhoff's potentials $\varphi_j(q, \cdot)$, for $j = 1, 2, 3$, as the unique (up to an additive constant) solutions in $\mathcal{F}(q)$ of the following Neumann problem:

$$\Delta\varphi_j(q, \cdot) = 0 \quad \text{in } \mathcal{F}(q), \quad (2.8a)$$

$$\frac{\partial\varphi_j}{\partial n}(q, \cdot) = K_j(q, \cdot) \quad \text{on } \partial\mathcal{S}(q), \quad (2.8b)$$

$$\frac{\partial\varphi_j}{\partial n}(q, \cdot) = 0 \quad \text{on } \partial\Omega. \quad (2.8c)$$

We concatenate the K_j and φ_j into the vectors:

$$\mathbf{K}(q, \cdot) := (K_1(q, \cdot), K_2(q, \cdot), K_3(q, \cdot))^t \text{ and} \quad (2.9a)$$

$$\boldsymbol{\varphi}(q, \cdot) := (\varphi_1(q, \cdot), \varphi_2(q, \cdot), \varphi_3(q, \cdot))^t, \quad (2.9b)$$

where the exponent t denotes the transpose of the vector.

Stream function for the circulation term. For every q in \mathcal{Q} , there exists a unique $C(q) \in \mathbb{R}$ such that the unique solution $\psi(q, \cdot)$ of the Dirichlet problem:

$$\Delta\psi(q, \cdot) = 0 \quad \text{in } \mathcal{F}(q) \quad (2.10a)$$

$$\psi(q, \cdot) = C(q) \quad \text{on } \partial\mathcal{S}(q) \quad (2.10b)$$

$$\psi(q, \cdot) = 0 \quad \text{on } \partial\Omega, \quad (2.10c)$$

satisfies

$$\int_{\partial\mathcal{S}(q)} \frac{\partial\psi}{\partial n}(q, \cdot) \, ds = -1. \quad (2.10d)$$

This can be seen easily by defining the corresponding harmonic function $\tilde{\psi}(q, \cdot)$ with $\tilde{\psi}(q, \cdot) = -1$ on $\partial\mathcal{S}(q)$ and $\tilde{\psi}(q, \cdot) = 0$ on $\partial\Omega$ and renormalizing it. Indeed the strong maximum principle ensures that $\frac{\partial\tilde{\psi}}{\partial n}(q, \cdot) < 0$ on $\partial\mathcal{S}(q)$, so that $\int_{\partial\mathcal{S}(q)} \frac{\partial\tilde{\psi}}{\partial n}(q, \cdot) \, ds < 0$.

The function $C(q)$ is actually the opposite of the inverse of the condenser capacity of $\mathcal{S}(q)$ in Ω , that is, of $\int_{\mathcal{F}(q)} |\nabla\tilde{\psi}(q, \cdot)|^2 \, dx$. Observe that

$$\forall q \in \mathcal{Q}, \quad C(q) = - \int_{\mathcal{F}(q)} |\nabla\tilde{\psi}(q, \cdot)|^2 \, dx < 0, \quad (2.11)$$

$$C \in C^\infty(\mathcal{Q}; (-\infty, 0)) \text{ and depends on } \mathcal{S}_0 \text{ and } \Omega. \quad (2.12)$$

Concerning (2.12) and other similar properties below about the regularity with respect to the shape, we refer to [15] or [37].

Decomposition of the fluid velocity. The fluid velocity u satisfies a div-curl type system in the doubly-connected domain $\mathcal{F}(q)$, constituted of (2.1b), (2.2), (2.1e), (2.1f), and of (2.3). When the solid position q in \mathcal{Q} , and the right hand sides of these equations, including the solid velocity $q' = (q'_1, q'_2, q'_3) \in \mathbb{R}^3$ are given, the fluid velocity u is determined in a unique way and we will therefore

denote it by $u(q, \cdot)$. Moreover, using (2.8), (2.9) and (2.10), the solution $u(q, \cdot)$ takes the form:

$$u(q, \cdot) = u_{q'}(q, \cdot) + u_\gamma(q, \cdot), \quad (2.13)$$

with

$$u_{q'}(q, \cdot) := \nabla(\varphi(q, \cdot) \cdot q') = \nabla \left(\sum_{j=1}^3 \varphi_j(q, \cdot) q'_j \right) \quad (2.14a)$$

$$\text{and } u_\gamma(q, \cdot) := \gamma \nabla^\perp \psi(q, \cdot). \quad (2.14b)$$

So besides the dependence with respect to \mathcal{S}_0 , to Ω and to the space variable, $u_{q'}$ depends on q and linearly on q' while u_γ depends on q and linearly on γ . Notice that in the system of equations (2.1), the initial data (2.1g) for the fluid is no longer required since it can be deduced from the given circulation γ and the initial data of the solid through the functions $\varphi(0, \cdot)$ and $\psi(0, \cdot)$.

Inertia matrices. We can now define the mass matrices

$$M_g := \text{diag}(\mathcal{J}, m, m), \quad (2.15a)$$

$$M_a(q) := \int_{\partial\mathcal{S}(q)} \varphi(q, \cdot) \otimes \frac{\partial\varphi}{\partial n}(q, \cdot) ds = \left(\int_{\mathcal{F}(q)} \nabla\varphi_i \cdot \nabla\varphi_j dx \right)_{1 \leq i, j \leq 3}, \quad (2.15b)$$

$$M(q) := M_g + M_a(q). \quad (2.15c)$$

The matrix $M(q)$ corresponds to the sum of the genuine inertia M_g of the body and the so-called added inertia $M_a(q)$, which, loosely speaking, measures how much the surrounding fluid resists the acceleration of the body motion (since the fluid undergoes an acceleration as well). Both M_g and $M_a(q)$ are symmetric and positive-semidefinite, and M_g is positive definite.

Christoffel symbols. A bilinear symmetric mapping $\Gamma(q)$ associated with $M(q)$ can be defined as follows:

$$\langle \Gamma(q), p, p \rangle := \left(\sum_{1 \leq i, j \leq 3} \Gamma_{i,j}^k(q) p_i p_j \right)_{1 \leq k \leq 3} \in \mathbb{R}^3, \quad (2.16a)$$

where, for every $i, j, k \in \{1, 2, 3\}$, we set

$$\Gamma_{i,j}^k(q) := \frac{1}{2} \left((M_a)_{k,j}^i + (M_a)_{k,i}^j - (M_a)_{i,j}^k \right)(q). \quad (2.16b)$$

In this identity, the notation $(M_a)_{i,j}^k$ stands for the partial derivative of the entry of indices (i, j) of the matrix M_a with respect to q_k , that is

$$(M_a)_{i,j}^k := \frac{\partial (M_a)_{i,j}}{\partial q_k}. \quad (2.16c)$$

With a slight imprecision, we call the coefficients $\Gamma_{i,j}^k$ the *Christoffel symbols* associated with the mass matrix. Usually, one should multiply by $M(q)^{-1}$ the right hand side (2.16b) considered as a column vector indexed by k to get the standard Christoffel symbols.

We underline that since the genuine inertia M_g of the body is independent of the position q of the solid, only the added inertia is involved in the Christoffel symbols.

Force term. Eventually, we also define the column vectors:

$$B(q) := \int_{\partial\mathcal{S}(q)} \left(\frac{\partial\psi}{\partial n} \left(\frac{\partial\varphi}{\partial n} \times \frac{\partial\varphi}{\partial\tau} \right) \right) (q, \cdot) ds, \quad (2.17a)$$

$$E(q) := -\frac{1}{2} \int_{\partial\mathcal{S}(q)} \left(\left| \frac{\partial\psi}{\partial n} \right|^2 \frac{\partial\varphi}{\partial n} \right) (q, \cdot) ds, \quad (2.17b)$$

and for p in \mathbb{R}^3 the force term

$$F(q, p) := \gamma^2 E(q) + \gamma p \times B(q). \quad (2.17c)$$

We recall that γ denotes the circulation around the body.

Remark 2.1. The notations E and B are chosen on purpose to highlight the analogy with the Lorentz force acting on a charged particle moving under the influence of an electromagnetic field (E, B) . This force vanishes if $\gamma = 0$.

It can be checked that

$$M \in C^\infty(\mathcal{Q}; S_3^{++}(\mathbb{R})) \text{ and depends on } \mathcal{S}_0, m, \mathcal{J} \text{ and } \Omega, \quad (2.18a)$$

$$F \in C^\infty(\mathcal{Q} \times \mathbb{R}^3; \mathbb{R}^3) \text{ and depends on } \mathcal{S}_0, \gamma \text{ and } \Omega, \\ \text{and vanishes when } \gamma = 0, \quad (2.18b)$$

$$\Gamma \in C^\infty(\mathcal{Q}; \mathcal{BL}(\mathbb{R}^3 \times \mathbb{R}^3; \mathbb{R}^3)) \text{ and depends on } \mathcal{S}_0 \text{ and } \Omega. \quad (2.18c)$$

Above $S_3^{++}(\mathbb{R})$ denotes the set of real symmetric positive-definite 3×3 matrices, $\mathcal{BL}(\mathbb{R}^3 \times \mathbb{R}^3; \mathbb{R}^3)$ denotes the space of bilinear mappings from $\mathbb{R}^3 \times \mathbb{R}^3$ to \mathbb{R}^3 .

We stress that M does not depend on the circulation γ whereas F does not depend on m and \mathcal{J} and Γ does not depend on m, γ and \mathcal{J} . In the following, when specifying these dependences is relevant, we will denote

$$M[\mathcal{S}_0, m, \mathcal{J}, \Omega], \Gamma[\mathcal{S}_0, \Omega] \text{ and } F[\mathcal{S}_0, \gamma, \Omega] \text{ instead of } M, \Gamma \text{ and } F. \quad (2.19)$$

Now our first result is a rephrasing of System (2.1) as an ordinary differential equation.

Theorem 2.2. *Up to the first collision, System (2.1) is equivalent to the second order ODE*

$$M(q)q'' + \langle \Gamma(q), q', q' \rangle = F(q, q'), \quad (2.20)$$

with Cauchy data

$$q(0) = 0, \quad q'(0) = (\omega_0, \ell_0) \in \mathbb{R} \times \mathbb{R}^2.$$

Let us emphasize that (2.20) only determines the body motion. The fluid velocity $u(q, \cdot)$ is then deduced by (2.13) and (2.14).

The proof of Theorem 2.2 is postponed to Section 5.

Remark 2.3. In the potential case, i.e. $\gamma = 0$, the ODE (2.20) means that the particle is moving along the geodesics associated with the Riemann metric induced on \mathcal{Q} by the matrix $M(q)$, cf. [30].

According to classical ODE theory, given γ , there exists a maximal time $T > 0$ and a unique maximal solution q in $C^\infty([0, T]; \mathcal{Q})$ to (2.20) with Cauchy data $q(0) = (0, 0)$, $q'(0) = (\omega_0, \ell_0)$.

Moreover, T is the time of the first collision of the solid with the outer boundary of the fluid domain. If there is no collision, then $T = +\infty$. This follows from Corollary 2.6 below, which itself relies on an energy argument. Indeed an important feature of the system (2.20) is that it is conservative. More precisely, let us defined for (q, p) in $\mathcal{Q} \times \mathbb{R}^3$,

$$\mathcal{E}(q, p) := \frac{1}{2}M(q)p \cdot p + U(q), \quad (2.21)$$

where the potential energy $U(q)$ is given by

$$U(q) := -\frac{1}{2}\gamma^2 C(q),$$

with $C(q)$ given by (2.10). We will prove the following in Section 5.5.

Lemma 2.4. *The derivative $DC(q)$ of $C(q)$ with respect to q satisfies:*

$$\forall q \in \mathcal{Q}, \quad E(q) = \frac{1}{2}DC(q). \quad (2.22)$$

As a corollary of Lemma 2.4 we have the following result regarding the energy conservation.

Proposition 2.5. *For any q in $C^\infty([0, T]; \mathcal{Q})$ satisfying (2.20) one has*

$$\frac{d}{dt}\mathcal{E}(q, q') = 0. \quad (2.23)$$

We will prove Proposition 2.5 in Section 5.6.

Let us emphasize that the energy function \mathcal{E} belongs to $C^\infty(\mathcal{Q} \times \mathbb{R}^3; \mathbb{R})$ and is the sum of two positive terms, see (2.12) and (2.18a). In addition to its dependence on q and p , the energy \mathcal{E} depends on $\mathcal{S}_0, m, \mathcal{J}, \gamma$ and Ω .

If we assume that the body remains at a distance greater than $\delta > 0$ from the boundary we may infer a bound for the body velocity. Indeed we have the following immediate corollary of Proposition 2.5.

Corollary 2.6. *Let \mathcal{S}_0 a subset of Ω , p_0 in \mathbb{R}^3 and (γ, m, \mathcal{J}) in $\mathbb{R} \times (0, +\infty) \times (0, +\infty)$. Let $\delta > 0$ and q in $C^\infty([0, T]; \mathcal{Q} \times \mathbb{R}^3)$ satisfying (2.20) with the Cauchy data $(q, q')(0) = (0, p_0)$ and such that $d(\mathcal{S}(q(t)), \partial\Omega) > \delta$ in $[0, T]$. Then there exists $K > 0$ depending only on $\mathcal{S}_0, \Omega, p_0, \gamma, m, \mathcal{J}$ and δ such that $|q'|_{\mathbb{R}^3} \leq K$ on $[0, T]$.*

Let us refer here to [16] and [31] for a study of the collision of a solid moving in a potential flow (that is in the case where $\gamma = 0$) with the fixed boundary of the fluid domain.

Let us now turn our attention back to the Christoffel symbols defined in (2.16). They actually can be split into two parts: one taking into account the effect of the solid rotation and the other part encoding the effect of the exterior boundary. First, we introduce the impulses ρ_g, ρ_a, ρ in \mathbb{R} and P_g, P_a and P in \mathbb{R}^2 by the following relations

$$\begin{pmatrix} \rho_g \\ P_g \end{pmatrix} := M_g p, \quad \begin{pmatrix} \rho_a \\ P_a \end{pmatrix} := M_a(q) p, \quad \begin{pmatrix} \rho \\ P \end{pmatrix} := \begin{pmatrix} \rho_g + \rho_a \\ P_g + P_a \end{pmatrix}. \quad (2.24)$$

Then for every q in \mathcal{Q} and for every $p = (\omega, \ell)$ in \mathbb{R}^3 , we let:

$$\langle \Gamma^{\text{rot}}(q), p, p \rangle := - \begin{pmatrix} 0 \\ P_a \end{pmatrix} \times p - \omega M_a(q) \begin{pmatrix} 0 \\ \ell^\perp \end{pmatrix} \in \mathbb{R}^3. \quad (2.25)$$

We can notice that one also has

$$\langle \Gamma^{\text{rot}}(q), p, p \rangle = - \begin{pmatrix} 0 \\ P \end{pmatrix} \times p - \omega M(q) \begin{pmatrix} 0 \\ \ell^\perp \end{pmatrix} \in \mathbb{R}^3,$$

since the extra terms cancel out.

Next, for every $j, k, l \in \{1, 2, 3\}$, we set

$$\begin{aligned} (\Gamma^{\partial\Omega})_{k,l}^j(q) &:= \\ &\frac{1}{2} \int_{\partial\Omega} \left(\frac{\partial\varphi_j}{\partial\tau} \frac{\partial\varphi_k}{\partial\tau} K_l + \frac{\partial\varphi_j}{\partial\tau} \frac{\partial\varphi_l}{\partial\tau} K_k - \frac{\partial\varphi_k}{\partial\tau} \frac{\partial\varphi_l}{\partial\tau} K_j \right) (q, \cdot) \, ds, \end{aligned} \quad (2.26)$$

and we associate correspondingly $\Gamma^{\partial\Omega}(q)$ a symmetric form in $\mathcal{BL}(\mathbb{R}^3 \times \mathbb{R}^3; \mathbb{R}^3)$ so that for $p = (p_1, p_2, p_3) \in \mathbb{R}^3$:

$$\langle \Gamma^{\partial\Omega}(q), p, p \rangle := \left(\sum_{1 \leq k, l \leq 3} (\Gamma^{\partial\Omega})_{k,l}^j(q) p_k p_l \right)_{1 \leq j \leq 3} \in \mathbb{R}^3.$$

The Christoffel symbols satisfy the following relation.

Proposition 2.7. *For every q in \mathcal{Q} and for every p in \mathbb{R}^3 we have:*

$$\langle \Gamma(q), p, p \rangle = \langle \Gamma^{\text{rot}}(q), p, p \rangle + \langle \Gamma^{\partial\Omega}(q), p, p \rangle. \quad (2.27)$$

The proof of Proposition 2.7 is given in Section 5.4. We emphasize that in (2.27) and the expressions above, unlike (2.16c), there is no derivative with respect to q , that is, no more shape derivative.

We will see that in the decomposition (2.27), the term $\Gamma^{\partial\Omega}$ obeys a scaling law with respect to ε different from the one of Γ^{rot} (compare (9.4) and (9.6) below), which makes it of lower order.

The case where \mathcal{S}_0 is a disk is peculiar, and we focus on it for the rest of Subsection 2.1. In particular several degeneracies appear in this case:

- the added mass matrix $M_a(q)$ degenerates (it becomes of rank 2),
- the potentials φ_2 , φ_3 and ψ depend on q only through the position h_c of the center of the disk $\mathcal{S}(q)$, and in particular so do E_2 , E_3 , B_1 and $(M_{a,i,j})_{i,j=2,3}$,
- the dynamics of the solid also degenerates in the sense that it satisfies $\mathcal{J}\vartheta'' = mh'' \cdot (h_c - h)^\perp$.

Above and in the sequel we denote with an index (in normal font type) the coordinates of E or B (we will sometimes use italic type indices for other purposes, in a way that should not be ambiguous). Note that in particular if the solid is homogeneous, $h = h_c$ and ϑ' is constant over time.

As a consequence, in this case where \mathcal{S}_0 is a disk, we establish a particular reduction of the dynamics. We will use a block decomposition of the matrix $M_a = M_a(q)$:

$$M_a =: \begin{pmatrix} m_\# & \mu^t \\ \mu & M_b \end{pmatrix}, \quad (2.28)$$

where M_b is a symmetric 2×2 matrix. This matrix is useful in the case where \mathcal{S}_0 is a disk of center ζ (we will use this notation later in a broader context). Of course the position h_c of the center of $\mathcal{S}(t)$ is related to $q = (h, \vartheta)$ by

$$h_c = h + R(\vartheta)\zeta. \quad (2.29)$$

It is easy to see that $M_b(q)$ depends on q only through h_c so that we define

$$\tilde{M}_b(h_c) := M_b(q). \quad (2.30)$$

We associate with the matrix field $\tilde{M}_b(h_c)$ a bilinear symmetric mapping $\Gamma_b(h_c)$ defined as follows: for $p_b \in \mathbb{R}^2$,

$$\langle \Gamma_b(h_c), p_b, p_b \rangle := \left(\sum_{1 \leq i, j \leq 2} (\Gamma_b)_{i,j}^k(h_c) p_{b,i} p_{b,j} \right)_{1 \leq k \leq 2} \in \mathbb{R}^2, \quad (2.31a)$$

where, for every $i, j, k \in \{1, 2\}$, we set

$$(\Gamma_b)_{i,j}^k(h_c) := \frac{1}{2} \left((\tilde{M}_b)_{k,j}^i + (\tilde{M}_b)_{k,i}^j - (\tilde{M}_b)_{i,j}^k \right) (h_c). \quad (2.31b)$$

In this identity, the notation $(\tilde{M}_b)_{i,j}^k$ stands for the partial derivative of the entry of indices (i, j) of the matrix \tilde{M}_b with respect to $(h_c)_k$, that is

$$(\tilde{M}_b)_{i,j}^k := \frac{\partial (\tilde{M}_b)_{i,j}}{\partial (h_c)_k}. \quad (2.31c)$$

The field $E(q)$ also depends on q only through h_c and we define $E_b(h_c)$ in \mathbb{R}^2 by

$$E_b(h_c) := (E_2, E_3)^t(q). \quad (2.32)$$

In the same way, $B(q)$ depends on q only through h_c and we define $\tilde{B}_1(h_c)$ in \mathbb{R} by the relation

$$\tilde{B}_1(h_c) = B_1(q). \quad (2.33)$$

With these settings, the dynamics can be described as follows.

Theorem 2.8. *In the case where \mathcal{S}_0 is a disk of center ζ , System (2.1) is equivalent, up to the first collision, to the following differential system*

$$mh'' + \tilde{M}_b(h_c)h_c'' + \langle \Gamma_b(h_c), h'_c, h'_c \rangle = \gamma^2 E_b(h_c) - \gamma \tilde{B}_1(h_c)(h'_c)^\perp, \quad (2.34a)$$

$$h_c - h = R(\vartheta)\zeta, \quad (2.34b)$$

$$\mathcal{J}\vartheta'' = (h_c - h)^\perp \cdot mh'', \quad (2.34c)$$

with Cauchy data $(h, \vartheta)(0) = 0$, $(\vartheta, h)'(0) = (\omega_0, \ell_0)$ in $\mathbb{R} \times \mathbb{R}^2$.

As for Theorem 2.2, the fluid velocity $u(q, \cdot)$ is then deduced by (2.13) and (2.14).

The proof of Theorem 2.8 is given in Section 5.7.

We have now at our disposal all the material to deal with the limit of the dynamics when the size of the solid goes to 0. As mentioned above, we distinguish two cases:

- Case (i): the mass of the solid is fixed (and then the solid tends to a massive pointwise particle), and
- Case (ii): the mass tends to 0 along with the size (and then the solid tends to a massless pointwise particle).

2.2 Case (i): Dynamics of a solid shrinking to a pointwise massive particle

From now on, we suppose that $0 \in \Omega$ and scale \mathcal{S}_0 around 0. Precisely, for every ε in $(0, 1]$, we denote

$$\mathcal{S}_{0,\varepsilon} := \varepsilon\mathcal{S}_0, \quad (2.35)$$

and for every $q = (\vartheta, h) \in \mathbb{R}^3$,

$$\mathcal{S}_\varepsilon(q) := R(\vartheta)\mathcal{S}_{0,\varepsilon} + h \quad \text{and} \quad \mathcal{F}_\varepsilon(q) = \Omega \setminus \mathcal{S}_\varepsilon(q). \quad (2.36)$$

We recall that $h(0) = 0$ so that (2.35) represents a homothety centered at $h(0)$. Without loss of generality, we suppose that for any $\varepsilon \in (0, 1]$, one has $\mathcal{S}_{0,\varepsilon} \subset \Omega$ (considering some $\mathcal{S}_{0,\varepsilon}$ as the initial solid \mathcal{S}_0 if necessary).

In Case (i) the solid occupying the domain $\mathcal{S}_\varepsilon(q)$ is assumed to have a mass and a moment of inertia of the form

$$m_\varepsilon = m \quad \text{and} \quad \mathcal{J}_\varepsilon = \varepsilon^2 \mathcal{J}_1, \quad (2.37)$$

where $m > 0$ and $\mathcal{J}_1 > 0$ are fixed.

With these settings, we denote by q_ε the solution to the ODE (2.20) associated with $M_\varepsilon := M[\mathcal{S}_{0,\varepsilon}, m_\varepsilon, \mathcal{J}_\varepsilon, \Omega]$, $\Gamma_\varepsilon := \Gamma[\mathcal{S}_{0,\varepsilon}, \Omega]$ and $F_\varepsilon := F[\mathcal{S}_{0,\varepsilon}, \gamma, \Omega]$ in place of M , Γ and F , respectively, defined on the maximal time interval $[0, T^\varepsilon)$. As before we decompose q_ε into

$$q_\varepsilon = (\vartheta_\varepsilon, h_\varepsilon) \in \mathbb{R} \times \mathbb{R}^2.$$

We emphasize that γ and the Cauchy data (q_0, p_0) do not depend on ε . The latter are decomposed into

$$q_0 = (0, 0) \text{ and } p_0 = (\omega_0, \ell_0).$$

Our first result is the convergence, in this setting, of h_ε to the solution of a massive point vortex equation. Let us introduce this limit equation. Let (\bar{h}, \bar{T}) be the maximal solution of the ODE:

$$m\bar{h}'' = \gamma \left(\bar{h}' - \gamma u^\Omega(\bar{h}) \right)^\perp \quad \text{for } t \in [0, \bar{T}), \quad (2.38a)$$

$$\text{with } \bar{h}(0) = 0 \text{ and } \bar{h}'(0) = \ell_0, \quad (2.38b)$$

where u^Ω is the Kirchhoff-Routh velocity defined as follows. Consider first $\psi_0^\Omega(h, \cdot)$, the solution to the following Dirichlet problem:

$$\Delta \psi_0^\Omega(h, \cdot) = 0 \text{ in } \Omega, \quad \psi_0^\Omega(h, \cdot) = G(\cdot - h) \text{ on } \partial\Omega, \quad (2.39)$$

where

$$G(r) := -\frac{1}{2\pi} \ln |r|. \quad (2.40)$$

The Kirchhoff-Routh stream function ψ^Ω is defined as

$$\psi^\Omega(x) := \frac{1}{2} \psi_0^\Omega(x, x), \quad x \in \Omega, \quad (2.41)$$

and the Kirchhoff-Routh velocity u^Ω is defined by

$$u^\Omega := \nabla^\perp \psi^\Omega, \quad (2.42)$$

where $\nabla^\perp := (-\partial_2, \partial_1)$.

The existence of (\bar{h}, \bar{T}) follows from classical ODE theory. Moreover it follows from the energy conservation stated below in (3.1) and from the regularity of the Kirchhoff-Routh stream function ψ^Ω in Ω that \bar{T} is the time of the first collision of \bar{h} with the outer boundary $\partial\Omega$ of the fluid domain. If there is no collision, then $\bar{T} = +\infty$.

The precise statement of our first convergence result is as follows.

Theorem 2.9. *Let \mathcal{S}_0 a subset of Ω as above, p_0 in \mathbb{R}^3 and (γ, m, \mathcal{J}) in $\mathbb{R} \times (0, +\infty) \times (0, +\infty)$. Let (\bar{h}, \bar{T}) be the maximal solution of (2.38). For every ε in $(0, 1]$, let $(q_\varepsilon, T_\varepsilon)$ be the maximal solution of (2.20) with respectively $M_\varepsilon = M[\mathcal{S}_{0,\varepsilon}, m_\varepsilon, \mathcal{J}_\varepsilon, \Omega]$, $\Gamma_\varepsilon = \Gamma[\mathcal{S}_{0,\varepsilon}, \Omega]$ and $F_\varepsilon = F[\mathcal{S}_{0,\varepsilon}, \gamma, \Omega]$ in place of M , Γ and F , respectively, where $m_\varepsilon, \mathcal{J}_\varepsilon$ are given by (2.37), and with the initial data $q_\varepsilon(0) = 0$ and $q'_\varepsilon(0) = p_0$. Then, as $\varepsilon \rightarrow 0^+$, one has $\liminf T_\varepsilon \geq \bar{T}$, and for all T in $(0, \bar{T})$, one has $h_\varepsilon \rightharpoonup \bar{h}$ in $W^{2,\infty}([0, T]; \mathbb{R}^2)$ weak- \star and $\varepsilon \vartheta_\varepsilon \rightharpoonup 0$ in $W^{2,\infty}([0, T]; \mathbb{R})$ weak- \star .*

2.3 Case (ii): Dynamics of a solid shrinking to a pointwise massless particle

In this section the solid is still assumed to occupy initially the domain $\mathcal{S}_{0,\varepsilon}$ given by (2.35) (satisfying the same assumptions as above) but we assume now that it has a mass and a moment of inertia given by

$$m_\varepsilon = \alpha_\varepsilon m_1 \text{ and } \mathcal{J}_\varepsilon = \alpha_\varepsilon \varepsilon^2 \mathcal{J}_1, \quad (2.43)$$

where $\alpha_1 = 1$ and $\alpha_\varepsilon \rightarrow 0^+$ as $\varepsilon \rightarrow 0^+$, and where $m_1 > 0$ and $\mathcal{J}_1 > 0$ are fixed. In order to simplify the notations we will assume that α_ε is of the form

$$\alpha_\varepsilon = \varepsilon^\alpha, \quad (2.44)$$

with $\alpha > 0$. The particular case where $\alpha = 2$ corresponds to the case of a fixed solid density. Case (i) corresponded to the case where $\alpha = 0$.

In this setting, we denote by $q_\varepsilon = (\vartheta_\varepsilon, h_\varepsilon)$ in $\mathbb{R} \times \mathbb{R}^2$ the solution to the ODE (2.20) associated with $M_\varepsilon := M[\mathcal{S}_{0,\varepsilon}, m_\varepsilon, \mathcal{J}_\varepsilon, \Omega]$, $\Gamma_\varepsilon := \Gamma[\mathcal{S}_{0,\varepsilon}, \Omega]$ and $F_\varepsilon := F[\mathcal{S}_{0,\varepsilon}, \gamma, \Omega]$ in place of M , Γ and F , respectively, defined on the maximal time interval $[0, T^\varepsilon)$. We stress that the circulation γ and the Cauchy data are still assumed independent of ε . Moreover we will assume here that

$$\gamma \neq 0.$$

Our second result is the convergence of h_ε as $\varepsilon \rightarrow 0^+$ to the solution of the point vortex equation:

$$\bar{h}' = \gamma u^\Omega(\bar{h}) \text{ for } t > 0, \text{ with } \bar{h}(0) = 0. \quad (2.45)$$

It is well-known that the solution \bar{h} is global in time, and in particular that there is no collision of the point vortex with the external boundary $\partial\Omega$. This follows from (3.2) below and the fact that $\psi^\Omega(h) \rightarrow +\infty$ when h gets close to $\partial\Omega$, see for instance [38, Eq. (1.27)].

Our result in this situation is the following.

Theorem 2.10. *Let \mathcal{S}_0 a subset of Ω as above, different from a disk. Let $\gamma \neq 0$, p_0 in \mathbb{R}^3 and (m_1, \mathcal{J}_1) in $(0, +\infty) \times (0, +\infty)$. Let us consider the global solution \bar{h} to (2.45) and for every ε in $(0, 1]$, let $(q_\varepsilon, T_\varepsilon)$ the maximal solution to (2.20) with respectively $M_\varepsilon = M[\mathcal{S}_{0,\varepsilon}, m_\varepsilon, \mathcal{J}_\varepsilon, \Omega]$, $\Gamma_\varepsilon = \Gamma[\mathcal{S}_{0,\varepsilon}, \Omega]$ and $F_\varepsilon = F[\mathcal{S}_{0,\varepsilon}, \gamma, \Omega]$ in place of M , Γ and F , where $m_\varepsilon, \mathcal{J}_\varepsilon$ are given by (2.43), and with the initial data $q_\varepsilon(0) = 0$ and $q_\varepsilon'(0) = p_0$. Then, as $\varepsilon \rightarrow 0^+$, one has $T_\varepsilon \rightarrow +\infty$ and $h_\varepsilon \rightharpoonup \bar{h}$ in $W^{1,\infty}([0, T]; \mathbb{R}^2)$ weak- \star for all $T > 0$.*

In the case where \mathcal{S}_0 is a disk, the statement needs a slight modification.

Theorem 2.11. *Let \mathcal{S}_0 a disk in Ω , with center ζ and center of mass 0. Let p_0 in \mathbb{R}^3 and (γ, m, \mathcal{J}) in $\mathbb{R} \times (0, +\infty) \times (0, +\infty)$. Let us consider the global solution \bar{h} to (2.45) and for every ε in $(0, 1]$, let $(q_\varepsilon, T_\varepsilon)$ be as in Theorem 2.10. Let $h_{c,\varepsilon} = h_\varepsilon + \varepsilon R(\vartheta_\varepsilon)\zeta$ the center of the disk. Then, as $\varepsilon \rightarrow 0^+$, one has $T_\varepsilon \rightarrow +\infty$ and $h_{c,\varepsilon} \rightharpoonup \bar{h}$ in $W^{1,\infty}([0, T]; \mathbb{R}^2)$ weak- \star for all $T > 0$.*

It is straightforward to check that this involves the convergence of h_ε to \bar{h} in L^∞ weak- \star for all $T > 0$. Actually, one even gets the convergence in $W^{\beta,\infty}(0,T)$ weak- \star for all $T > 0$ with $\beta := \min(1, \frac{2}{\alpha})$, see Section 7.5. As we will see during the proofs, if either \mathcal{S}_0 is homogeneous (so that the center of the disk and the center of gravity coincide) or if $\alpha \leq 2$, Theorem 2.10 is actually valid without change.

3 Comments and organization of the paper

3.1 Comments

Energy for the limit equations. Let us first turn to the energy conservations for the limit equations. It is elementary to see that for any h in $C^\infty([0,T];\Omega)$ satisfying (2.38) one has

$$\frac{d}{dt}\mathcal{E}_{(i)}(h, h') = 0, \quad \text{with } \mathcal{E}_{(i)}(h, h') := \frac{1}{2}mh' \cdot h' - \gamma^2\psi^\Omega(h). \quad (3.1)$$

Similarly for any h in $C^\infty([0,T];\Omega)$ satisfying (2.45) one has

$$\frac{d}{dt}\mathcal{E}_{(ii)}(h) = 0, \quad \text{with } \mathcal{E}_{(ii)}(h) := \gamma^2\psi^\Omega(h). \quad (3.2)$$

Theorem 2.9 and Theorem 2.10 therefore respectively assert the convergence of the trajectories of the system (2.20) associated with the energy $\mathcal{E}(q, p)$ given by (2.23) to those of the Hamiltonian system (2.38) associated with the energy (3.1) in Case (i) and to the Hamiltonian system (2.45) associated with the energy (3.2) in Case (ii). For further Hamiltonian aspects related to Systems (2.1) and (2.45), we refer for instance to [25,32,39].

Modulated energy argument. In our analysis of the massless Case (ii) we use a modulated energy argument in a crucial way. This type of argument was also used in [29] for a related issue but we would like to emphasize that the modulation occurs for a different part of the energy. More precisely, in the paper [29] an Euler-Vlasov system is introduced as a mean-field model for massive pointwise particles moving in a perfect incompressible fluid according to a system of ODEs which is believed to extend the equation (2.38) of Case (i). The last section of [29] deals with the limit where the individual mass of the particles converges to 0. Then a modulated energy, inspired by the paper [4] of Brenier regarding yet another issue: the gyrokinetic limit for the Vlasov-Poisson system, is used in order to obtain a hydrodynamic convergence, that is the convergence of the macroscopic mixture velocity associated by the Biot-Savart law with both fluid vorticity and particles circulations. Referring to the decomposition of the energy in (2.21) or to Eq. (78) in [29] we sketch the following opposition. There the potential part of the energy or more exactly of its modulation is crucial. On the contrary we will use here a modulation

of the kinetic part of the energy in order to deduce the limit of the particle dynamics, the fluid dynamics being subjugated to the particle one.

Miscellaneous remarks. The limit systems obtained in Cases (i) and (ii) do not depend on the body shape nor on the value of $\alpha > 0$. Still the proof is simpler in the case where the body is a homogeneous disk. Indeed if \mathcal{S}_0 is a disk, in both Cases (i) and (ii), it follows directly from (2.1d) that the rotation ϑ_ε satisfies, for any ε in $(0, 1)$, $\vartheta_\varepsilon(t) = t\omega_0$ as long as the solution exists.

We notice that the convergences in Theorem 2.9 and Theorem 2.10 cannot be improved unless some compatibility condition holds at initial time. One may however wonder if these convergences could be improved in the open interval $(0, T)$. It seems that some strong oscillations in time show up when $\varepsilon \rightarrow 0^+$ which prevent strong convergence from happening. We plan to study this phenomenon by a multi-scale approach of the solution of the ODE (2.20) in a forthcoming work. Once again the case where the body is a homogeneous disk is likely to simplify the discussion.

Let us also underline that the convergence of the dynamics of the solid involves some convergence in the fluid. It is not difficult to check that for any $k \in \mathbb{N}$ one has the weak- \star convergence, in $W^{1,\infty}(0, T; C^k(K))$ for any $T < \bar{T}$ in Case (i) and in $L^\infty(0, T; C^k(K))$ for any $T > 0$ in Case (ii), of the fluid velocity u_ε toward $\gamma u^\Omega(\cdot) + \frac{\gamma}{2\pi} \frac{(-\bar{h}(t))^\perp}{|-\bar{h}(t)|^2}$, for any compact set $K \subset \Omega$ not intersecting $\{\bar{h}(t), t \in [0, T]\}$.

In Case (i), one may also raise the question whether it is possible that $\liminf T_\varepsilon > \bar{T}$. This problem should be connected to the behaviour of the potentials and stream functions as the body approaches the boundary; see for instance [3], [5], [27] and [31] and references therein for this question.

The analysis performed in this paper can be easily adapted in order to cover the case where the circulation γ depends on ε under the form $\gamma^\varepsilon = \varepsilon^\beta \gamma^1$ with $\beta > 0$ in Case (i) and β in $(0, 1)$ in Case (ii). One obtains respectively at the limit the trivial equations $\bar{h}'' = 0$ and $\bar{h}' = 0$.

Our analysis should hold as well in the case of several bodies moving in the full plane or in a multiply-connected domain, as long as there is no collision. This will be tackled in a forthcoming work.

Another natural question is whether or not one may extend the results of Theorem 2.9 and Theorem 2.10 to rotational flows. This issue is tackled in a more restricted geometric set-up in the work [10].

3.2 Organization of the paper

The paper is organized as follows.

In Section 4 we deal with the simpler case when there is no external boundary. This case is well-known in the literature and has been in particular studied with complex analysis, using Blasius' lemma (see e.g. [25]). We shall use an alternative approach based on a lemma due to Lamb, cf. Lemma 4.3, allowing

one to exchange some normal and tangential components in some trilinear integrals over the body boundary. The case with no outer boundary is actually very important to tackle the general one.

Then Theorem 2.2, Proposition 2.5 and Proposition 2.7, which are independent of ε , are proved in Section 5.

In Sections 6 and 7, respectively, we prove Theorem 2.9 concerning the limit of a massive particle (Case (i)) and Theorem 2.10 concerning the limit of a massless particle (Case (ii)). These proofs rely on some asymptotic normal forms (6.7) and (7.10), respectively. These normal forms are the key point of the demonstration and allow to establish some renormalized and modulated energy estimates and prove the passage to the limit. They are established in the last two sections.

In Section 8 we establish some asymptotic expansions of stream and potential functions with respect to ε . These expansions involve two scales corresponding respectively to variations over length $O(1)$ and $O(\varepsilon)$ respectively on $\partial\Omega$ and $\partial\mathcal{S}_\varepsilon(q)$. The profiles appearing in these expansions are obtained by successive corrections, considering alternatively at their respective scales the body boundary from which the external boundary seems remote and the external boundary from which the body seems tiny, so that good approximations are given respectively by the case without external boundary and without the body.

Then in Section 9, we prove the normal forms (6.7) and (7.10). In order to do so we plug the expansions obtained in Section 8 into the expressions of the inertia matrix $M_\varepsilon(q)$, of the Christoffel symbol $\langle \Gamma_\varepsilon(q), p, p \rangle$ and of the force fields E_ε and $p \times B_\varepsilon$ and compute the leading terms of the resulting expansions. These expansions can themselves be plugged into the ODE (2.20) of Theorem 2.2. In particular, thanks to Lamb's lemma we will make appear in several terms of the expansions of E_ε and B_ε some coefficients of the added inertia of the solid as if the external boundary was not there. Strikingly this allows to combine the subprincipal terms of the expansions of E and B with the leading term of the expansion of Γ , see Lemma 9.9. This fact is essential in the proof.

4 Case without external boundary

In this section, we consider a simplified version of the problem that we are interested in. The simplification consists in assuming that the domain occupied by the fluid-solid system is the whole plane, i.e. $\Omega = \mathbb{R}^2$, the fluid being at rest at infinity.

The purpose of this section is not only to provide a “warm-up” for the much more involved “bounded” configuration. It turns out that several objects that will come up in the analysis will be also of central importance in the sequel.

We aim to give (in this simplified unbounded configuration) the counterparts of Theorems 2.2, 2.9 and 2.10.

The rephrasing of the equations driving the dynamics as an ODE (that is a result similar to Theorem 2.2), has been known since the investigations of Blasius, Kutta, Joukowski, Chaplygin and Sedov. Their analysis, relying on a complex-analytic approach, was then revisited following another approach which seems to date back to Lamb. Since we will elaborate on the latter in order to deal with the bounded case, we will first establish the counterpart of Theorem 2.2 relying on Lamb's analysis.

We will deduce from this ODE an energy conservation providing an analogous to Proposition 2.5.

Then we will investigate the passage to the limit of the dynamics when the size of the solid goes to 0 in both Cases (i) and (ii). Hence we will establish the counterparts of Theorem 2.9 and Theorem 2.10.

4.1 Recasting of the system as an ODE

In the case where $\Omega = \mathbb{R}^2$, the fluid-solid system is governed by the following set of coupled equations, quite similar to System (2.1):

$$\frac{\partial u}{\partial t} + (u \cdot \nabla)u + \nabla H = 0 \quad \text{and} \quad \operatorname{div} u = 0 \quad \text{in } \mathbb{R}^2 \setminus \mathcal{S}(t), \quad (4.1a)$$

$$(\mathcal{J}\vartheta'', mh'') = \left(\int_{\partial\mathcal{S}(t)} (x - h(t))^\perp \cdot Hn \, ds, \int_{\partial\mathcal{S}(t)} Hn \, ds \right), \quad (4.1b)$$

$$u \cdot n = (\vartheta'(\cdot - h)^\perp + h') \cdot n \quad \text{on } \partial\mathcal{S}(t) \quad \text{and} \quad u(x) \rightarrow 0 \quad \text{when } |x| \rightarrow +\infty, \quad (4.1c)$$

$$u|_{t=0} = u_0 \quad \text{in } \mathbb{R}^2 \setminus \mathcal{S}_0 \quad \text{and} \quad (\vartheta, h)(0) = (0, 0), \quad (\vartheta', h')(0) = (\omega_0, \ell_0). \quad (4.1d)$$

In order to state the aforementioned reformulation, we introduce the Kirchhoff potentials, the inertia matrices, the Christoffel symbols and the force term corresponding to this simplified case.

Kirchhoff potential. Let us first denote by $\varphi_j^{\mathcal{S}_0}$, for $j = 1, 2, 3$, the Kirchhoff's potentials in $\mathbb{R}^2 \setminus \mathcal{S}_0$ which are the functions that satisfy the following Neumann problem:

$$\Delta \varphi_j^{\mathcal{S}_0} = 0 \quad \text{in } \mathbb{R}^2 \setminus \mathcal{S}_0, \quad (4.2a)$$

$$\frac{\partial \varphi_j^{\mathcal{S}_0}}{\partial n} = \begin{cases} x^\perp \cdot n & \text{for } j = 1 \\ e_{j-1} \cdot n, & \text{for } j = 2, 3 \end{cases} \quad \text{on } \partial\mathcal{S}_0, \quad (4.2b)$$

$$\nabla \varphi_j^{\mathcal{S}_0}(x) \rightarrow 0 \quad \text{at infinity.} \quad (4.2c)$$

We also denote

$$\boldsymbol{\varphi}^{\mathcal{S}_0} := (\varphi_1^{\mathcal{S}_0}, \varphi_2^{\mathcal{S}_0}, \varphi_3^{\mathcal{S}_0})^t. \quad (4.3)$$

Stream function for the circulation term. In the same spirit as (2.10), we first introduce the function $\psi_{-1}^{S_0}$ as the solution of

$$-\Delta\psi_{-1}^{S_0} = 0 \quad \text{in } \mathbb{R}^2 \setminus \mathcal{S}_0, \quad (4.4a)$$

$$\psi_{-1}^{S_0} = C^{S_0} \quad \text{on } \partial\mathcal{S}_0, \quad (4.4b)$$

$$\psi_{-1}^{S_0} = O(\ln|x|) \quad \text{at infinity}, \quad (4.4c)$$

where the constant C^{S_0} is such that:

$$\int_{\partial\mathcal{S}_0} \frac{\partial\psi_{-1}^{S_0}}{\partial n} ds = -1. \quad (4.4d)$$

The existence and uniqueness of $\psi_{-1}^{S_0}$ will be recalled below in Proposition 8.1; one can still identify $\frac{\partial\psi_{-1}^{S_0}}{\partial n}$ as being the equilibrium measure of $\partial\mathcal{S}_0$.

Change of frame and decomposition of the fluid velocity. The vector field \underline{u} defined from the fluid velocity u by

$$\underline{u}(t, x) := R(\vartheta(t))^t u(t, R(\vartheta(t))x + h(t)), \quad (4.5)$$

satisfies the following div-curl type system in the doubly-connected domain \mathcal{F}_0 :

$$\operatorname{div} \underline{u} = 0 \text{ and } \operatorname{curl} \underline{u} = 0 \text{ in } \mathbb{R}^2 \setminus \mathcal{S}_0, \quad (4.6a)$$

$$\underline{u} \cdot n = (\underline{\ell} + \vartheta' x^\perp) \cdot n \quad \text{on } \partial\mathcal{S}_0, \quad (4.6b)$$

$$\underline{u} \rightarrow 0 \quad \text{at infinity}, \quad (4.6c)$$

$$\int_{\partial\mathcal{S}_0} \underline{u} \cdot \tau ds = \gamma, \quad (4.6d)$$

where $\gamma := \int_{\partial\mathcal{S}_0} u_0 \cdot \tau ds$ and $\underline{\ell}(t) := R(\vartheta(t))^t h'(t)$.

When the right hand sides of these equations are given the auxiliary velocity field \underline{u} is determined in a unique way and takes the form:

$$\underline{u} = \underline{u}_{(\vartheta', \underline{\ell})} + \underline{u}_\gamma \quad (4.7a)$$

where, for $\underline{\ell} = (\underline{\ell}_1, \underline{\ell}_2)$,

$$\underline{u}_{(\vartheta', \underline{\ell})} := \nabla(\varphi_1^{S_0} \vartheta' + \varphi_2^{S_0} \underline{\ell}_1 + \varphi_3^{S_0} \underline{\ell}_2) \quad \text{and} \quad \underline{u}_\gamma := \gamma \nabla^\perp \psi_{-1}^{S_0}. \quad (4.7b)$$

Inertia matrices. We define the added mass matrix

$$M_a^{S_0} := \int_{\partial\mathcal{S}_0} \varphi^{S_0} \otimes \frac{\partial\varphi^{S_0}}{\partial n} ds = \left(\int_{\mathbb{R}^2 \setminus \mathcal{S}_0} \nabla\varphi_i^{S_0} \cdot \nabla\varphi_j^{S_0} dx \right)_{1 \leq i, j \leq 3}. \quad (4.8)$$

This matrix is symmetric positive-semidefinite and depends only on \mathcal{S}_0 . Actually it is positive definite if and only if \mathcal{S}_0 is not a disk. Moreover, when \mathcal{S}_0 is a disk centered at the center of mass, this matrix reads $M_a^{S_0} =: \operatorname{diag}(0, m_a^{S_0}, m_a^{S_0})$ with $m_a^{S_0} > 0$.

Then we can introduce the mass matrix $M_{a,\vartheta}^{S_0}$ taking the rotation into account by

$$M_{a,\vartheta}^{S_0} := \mathcal{R}(\vartheta) M_a^{S_0} \mathcal{R}(\vartheta)^t. \quad (4.9)$$

Above we used the notation

$$\mathcal{R}(\vartheta) := \begin{pmatrix} 1 & 0 \\ 0 & R(\vartheta) \end{pmatrix} \in \text{SO}(3). \quad (4.10)$$

Christoffel symbols. We define for all $p = (\omega, \ell) \in \mathbb{R} \times \mathbb{R}^2$,

$$\langle \Gamma_\vartheta^{S_0}, p, p \rangle := - \begin{pmatrix} 0 \\ P_{a,\vartheta}^{S_0} \end{pmatrix} \times p - \omega M_{a,\vartheta}^{S_0} \begin{pmatrix} 0 \\ \ell^\perp \end{pmatrix}, \quad (4.11)$$

where $P_{a,\vartheta}^{S_0}$ denotes the last two coordinates of the impulse $M_{a,\vartheta}^{S_0} p$. Of course no $\Gamma^{\partial\Omega}$ is needed here.

Force term. Let us define the geometric constant

$$\zeta := - \int_{\partial S_0} x \frac{\partial \psi_{-1}^{S_0}}{\partial n} ds \quad (4.12)$$

which is a vector of \mathbb{R}^2 , depending only on S_0 , usually referred to as the conformal center of S_0 . Remark that if S_0 is a disk, $\psi_{-1}^{S_0}(x) = \frac{1}{2\pi} \ln |x - h_c(0)|$ where $h_c(0)$ is the center of the disk (recalling that the initial position of the center of mass is $h(0) = 0$); it follows that $\zeta = h_c(0)$ and the definition of ζ in (4.12) is coherent with the notation in Theorem 2.8.

Next we define the force term, for all $p = (\omega, \ell) \in \mathbb{R}^3$, as being:

$$F_\vartheta^{S_0}(p) := \gamma \begin{pmatrix} \zeta_\vartheta \cdot \ell \\ \ell^\perp - \omega \zeta_\vartheta \end{pmatrix}, \quad (4.13)$$

where

$$\zeta_\vartheta := R(\vartheta) \zeta. \quad (4.14)$$

An important feature of the force vector field $(\vartheta, p) \mapsto F_\vartheta^{S_0}(p)$ is that it is gyroscopic, in the sense of the following definition, see for instance [1, p. 428].

Definition 4.1. We say that a vector field F in $C^\infty(\mathbb{R}^3 \times \mathbb{R}^3; \mathbb{R}^3)$ is gyroscopic if for any (q, p) in $\mathbb{R}^3 \times \mathbb{R}^3$, $p \cdot F(q, p) = 0$.

Indeed, for any (ϑ, p) in $\mathbb{R} \times \mathbb{R}^3$, the force $F_\vartheta^{S_0}(p)$ can be written as

$$F_\vartheta^{S_0}(p) = \gamma p \times B_\vartheta^{S_0} \text{ with } B_\vartheta^{S_0} := \begin{pmatrix} -1 \\ \zeta_\vartheta^\perp \end{pmatrix}. \quad (4.15)$$

The next result is a reformulation of System (4.1) as an ordinary differential equation.

Theorem 4.2. *System (4.1) is equivalent to the second order ODE*

$$(M_g + M_{a,\vartheta}^{S_0}) q'' + \langle \Gamma_\vartheta^{S_0}, q', q' \rangle = F_\vartheta^{S_0}(q'), \quad (4.16)$$

with Cauchy data $q(0) = 0$, $q'(0) = (\omega_0, \ell_0)$ in $\mathbb{R} \times \mathbb{R}^2$.

The proof of Theorem 4.2 is postponed to Section 4.3. Let us underline that it is understood that we associate with the body equation given by (4.16) the fluid velocity u such that the vector field \underline{u} given by (4.5) satisfies (4.7). Theorem 4.2 is therefore a statement similar to Theorem 2.2 with $M_g + M_{a,\vartheta}^{S_0} = M[\mathcal{S}_0, m^1, \mathcal{J}^1, \mathbb{R}^2](q)$, $\Gamma_\vartheta^{S_0} = \Gamma[\mathcal{S}_0, \mathbb{R}^2](q)$ and $F_\vartheta^{S_0}(p) = F[\mathcal{S}_0, \gamma, \mathbb{R}^2](q, p)$. Observe in particular that the dependence on q of M , Γ and F reduces to a dependence on the rotation $R(\vartheta)$ only; from now on, we will mention this dependence on ϑ through an index, so q does no longer appear as an argument. Moreover the dependence of M , Γ and F on \mathcal{S}_0 , m_1 , \mathcal{J}_1 and γ is rather explicit.

According to classical ODE theory, given γ and \mathcal{S}_0 and some Cauchy data, there exists a unique global smooth solution to (4.16).

4.2 Lamb's lemma

In order to prove Theorem 4.2, we start with recalling a technical result borrowed from [21, Article 134a. (3) and (7)] and which is a cornerstone of our analysis. We recall that ξ_j and K_j , for $j = 1, 2, 3$, were defined in (2.6) and (2.7) respectively.

Lemma 4.3. *For any pair of vector fields u, v in $C^\infty(\overline{\mathbb{R}^2 \setminus \mathcal{S}_0}; \mathbb{R}^2)$ satisfying*

- $\operatorname{div} u = \operatorname{div} v = \operatorname{curl} u = \operatorname{curl} v = 0$,
- $u(x) = O(1/|x|)$ and $v(x) = O(1/|x|)$ as $|x| \rightarrow +\infty$,

one has, for any $j = 1, 2, 3$,

$$\int_{\partial \mathcal{S}_0} (u \cdot v) K_j(0, \cdot) \, ds = \int_{\partial \mathcal{S}_0} \xi_j(0, \cdot) \cdot \left((u \cdot n)v + (v \cdot n)u \right) \, ds. \quad (4.17)$$

Proof Let us start with the case where $j = 2$ or 3 . Then

$$\int_{\partial \mathcal{S}_0} (u \cdot v) K_j(0, \cdot) \, ds = \int_{\mathbb{R}^2 \setminus \mathcal{S}_0} \operatorname{div} \left((u \cdot v) \xi_j(0, \cdot) \right) \, dx, \quad (4.18)$$

by using that $u(x) = O(1/|x|)$ and $v(x) = O(1/|x|)$ when $|x| \rightarrow +\infty$. Therefore

$$\begin{aligned} \int_{\partial \mathcal{S}_0} (u \cdot v) K_j(0, \cdot) \, ds &= \int_{\mathbb{R}^2 \setminus \mathcal{S}_0} \xi_j(0, \cdot) \cdot \nabla (u \cdot v) \, dx \\ &= \int_{\mathbb{R}^2 \setminus \mathcal{S}_0} \xi_j(0, \cdot) \cdot (u \cdot \nabla v + v \cdot \nabla u) \, dx, \end{aligned} \quad (4.19)$$

using that $\operatorname{curl} u = \operatorname{curl} v = 0$. Now, integrating by parts, using that $\operatorname{div} u = \operatorname{div} v = 0$ and once again that $u(x) = O(1/|x|)$ and $v(x) = O(1/|x|)$ as $|x| \rightarrow +\infty$, we obtain (4.17) when $j = 2$ or 3 .

We now tackle the case where $j = 1$. We follow the same lines as above, with two precisions. First we observe that there is no contribution at infinity in (4.18) and (4.19) when $j = 1$ as well. Indeed ξ_1 and the normal to a centered circle are orthogonal. Moreover there is no additional distributed term coming from the integration by parts in (4.19) when $j = 1$ since

$$\int_{\mathbb{R}^2 \setminus \mathcal{S}_0} v \cdot (u \cdot \nabla_x \xi_j(0, \cdot)) + u \cdot (v \cdot \nabla_x \xi_j(0, \cdot)) \, dx = \int_{\mathbb{R}^2 \setminus \mathcal{S}_0} (v \cdot u^\perp + u \cdot v^\perp) \, dx = 0.$$

The proof is then complete. \square

4.3 Proof of Theorem 4.2

In order to transfer the equations in the body frame we recast the equations in terms of the vector field \underline{u} (defined from the fluid velocity u by (4.5)), of $\underline{\ell}(t) = R(\vartheta(t))^t h'(t)$ and of the auxiliary pressure $\underline{\Pi}$ (defined from the fluid pressure Π by $\underline{\Pi}(t, x) = \Pi(t, R(\vartheta(t))x + h(t))$). Equations (4.1) become

$$\frac{\partial \underline{u}}{\partial t} + [(\underline{u} - \underline{\ell} - \vartheta' x^\perp) \cdot \nabla] \underline{u} + \vartheta' \underline{u}^\perp + \nabla \underline{\Pi} = 0 \text{ for } x \in \mathbb{R}^2 \setminus \mathcal{S}_0 \quad (4.20a)$$

$$\operatorname{div} \underline{u} = 0 \text{ for } x \in \mathbb{R}^2 \setminus \mathcal{S}_0, \quad (4.20b)$$

$$\underline{u} \cdot n = (\underline{\ell} + \vartheta' x^\perp) \cdot n \text{ for } x \in \partial \mathcal{S}_0, \quad (4.20c)$$

$$(\mathcal{J}\vartheta'', m\underline{\ell}') = \left(\int_{\partial \mathcal{S}_0} x^\perp \cdot \underline{\Pi} n \, ds, \int_{\partial \mathcal{S}_0} \underline{\Pi} n \, ds - m\vartheta' \underline{\ell}^\perp \right). \quad (4.20d)$$

Using that the vector field \underline{u} is irrotational we infer that the equation (4.20a) can be written as

$$\frac{\partial \underline{u}}{\partial t} + \nabla(Q + \underline{\Pi}) = 0, \text{ with } Q := \frac{1}{2} |\underline{u}|^2 - (\underline{\ell} + \vartheta' x^\perp) \cdot \underline{u}.$$

On the other hand using an integration by parts we obtain

$$\left(\int_{\partial \mathcal{S}_0} \underline{\Pi} x^\perp \cdot n \, ds, \int_{\partial \mathcal{S}_0} \underline{\Pi} n \, ds \right) = \left(\int_{\mathbb{R}^2 \setminus \mathcal{S}_0} \nabla \underline{\Pi} \cdot \nabla \varphi_i^{\mathcal{S}_0} \, dx \right)_{i=1,2,3}.$$

Integrating by parts twice and taking into account the boundary condition (4.20c), one observes that the contribution of $\frac{\partial \underline{u}}{\partial t}$ is

$$\int_{\mathbb{R}^2 \setminus \mathcal{S}_0} \frac{\partial \underline{u}}{\partial t} \cdot \nabla \varphi_i^{\mathcal{S}_0}(x) \, dx = M_a^{\mathcal{S}_0} \underline{p}', \text{ where } \underline{p}' := \begin{pmatrix} \vartheta' \\ \underline{\ell} \end{pmatrix}.$$

Thus we obtain that the solid equations (4.20d) can be rewritten in the form

$$\begin{aligned} (M_g + M_a^{\mathcal{S}_0}) \underline{p}' + m\vartheta' \begin{pmatrix} 0 \\ \underline{\ell}^\perp \end{pmatrix} \\ = - \left(\frac{1}{2} \int_{\partial \mathcal{S}_0} |\underline{u}|^2 K_i \, ds - \int_{\partial \mathcal{S}_0} (\underline{\ell} + \vartheta' x^\perp) \cdot \underline{u} K_i \, ds \right)_{i=1,2,3}, \end{aligned}$$

where we simplified the notation by writing K_i instead of $K_i(0, \cdot)$ (defined in (2.7)), which corresponds to the initial position of the solid, that is $q = 0$. In the same way, we will write ξ_i for the vector fields $\xi_i(0, \cdot)$ (defined in (2.6)) in the computations below.

We compute the first term in the right hand side by using Lamb's lemma and the boundary conditions. For $i = 1, 2, 3$, we get

$$\begin{aligned} \frac{1}{2} \int_{\partial S_0} |\underline{u}|^2 K_i \, ds &= \int_{\partial S_0} ((\underline{\ell} + \vartheta' x^\perp) \cdot n) (\underline{u} \cdot n) K_i \, ds \\ &\quad + \int_{\partial S_0} ((\underline{\ell} + \vartheta' x^\perp) \cdot n) (\underline{u} \cdot \tau) (\xi_i \cdot \tau) \, ds, \end{aligned}$$

so that

$$\begin{aligned} (M_g + M_a^{S_0}) \underline{p}' + m \vartheta' \begin{pmatrix} 0 \\ \underline{\ell}^\perp \end{pmatrix} \\ = - \left(\sum_{k=1}^3 \underline{p}_k \int_{\partial S_0} (\underline{u} \cdot \tau) [(\xi_i \cdot \tau) K_k - (\xi_k \cdot \tau) K_i] \, ds \right)_{i=1,2,3} \end{aligned} \quad (4.21)$$

We observe that the brackets above are either vanishing (for $i = k$) or given by the following identities:

$$\begin{aligned} (\xi_2 \cdot \tau) K_3 - (\xi_3 \cdot \tau) K_2 &= 1, \quad (\xi_2 \cdot \tau) K_1 - (\xi_1 \cdot \tau) K_2 = x^\perp \cdot e_2, \\ \text{and } (\xi_3 \cdot \tau) K_1 - (\xi_1 \cdot \tau) K_3 &= -x^\perp \cdot e_1. \end{aligned} \quad (4.22)$$

Thanks to these equalities we compute the previous integrals in terms of the entries of the matrix $M_a^{S_0}$. Precisely we decompose $M_a^{S_0}$ into

$$M_a^{S_0} =: \begin{pmatrix} m_{\#}^{S_0} & (\mu^{S_0})^t \\ \mu^{S_0} & M_b^{S_0} \end{pmatrix}, \quad (4.23)$$

where $M_b^{S_0}$ is a symmetric 2×2 matrix. Then we have the following result.

Lemma 4.4. *There holds:*

$$\int_{\partial S_0} (\underline{u} \cdot \tau) x^\perp \, ds = \gamma \zeta^\perp - \vartheta' \mu - M_b^{S_0} \underline{\ell}. \quad (4.24)$$

Proof We first use (4.7) and (4.12) to get

$$\int_{\partial S_0} (\underline{u} \cdot \tau) x^\perp \, ds = \gamma \zeta^\perp + \sum_{j=1}^3 \underline{p}_j \int_{\partial S_0} (\nabla \varphi_j^{S_0} \cdot \tau) x^\perp \, ds.$$

Therefore, it only remains to observe that an integration by parts yields:

$$\int_{\partial S_0} (\nabla \varphi_j^{S_0} \cdot \tau) x^\perp \, ds = - \int_{\partial S_0} \varphi_j^{S_0} n \, ds = - ((M_a^{S_0})_{i,j})_{i=2,3}.$$

□

Combining now (4.21), (4.22), (4.6d) and (4.24) we end up with the identity:

$$[M_g + M_a^{S_0}] \underline{p}' + m \vartheta' \begin{pmatrix} 0 \\ \underline{\ell}^\perp \end{pmatrix} + \begin{pmatrix} \underline{\ell}^\perp \cdot M_b^{S_0} \underline{\ell} \\ \vartheta' (M_b^{S_0} \underline{\ell})^\perp \end{pmatrix} + \vartheta' \underline{p} \times \begin{pmatrix} 0 \\ \mu \end{pmatrix} = \gamma \underline{p} \times \begin{pmatrix} -1 \\ \zeta^\perp \end{pmatrix}.$$

Going back to the original frame leads to (4.16). This concludes the proof of Theorem 4.2.

4.4 Energy conservation

Let us first justify that $\Gamma_\vartheta^{S_0}$ defined in (4.11) are Christoffel symbols in the sense of (2.16b) (see the comment below (2.16b) regarding the slight modification of the usual definition of the Christoffel symbols). We develop, for any ϑ in \mathbb{R} , for any p in \mathbb{R}^3 , the bilinear symmetric mapping $\Gamma_\vartheta^{S_0}$ into

$$\langle \Gamma_\vartheta^{S_0}, p, p \rangle := \left(\sum_{1 \leq i, j \leq 3} (\Gamma_\vartheta^{S_0})_{i,j}^k p_i p_j \right)_{1 \leq k \leq 3} \in \mathbb{R}^3.$$

Then a tedious computation reveals that for every $i, j, k \in \{1, 2, 3\}$, we have:

$$(\Gamma_\vartheta^{S_0})_{i,j}^k = \frac{1}{2} \left((M_{a,\vartheta}^{S_0})_{k,j}^i + (M_{a,\vartheta}^{S_0})_{k,i}^j - (M_{a,\vartheta}^{S_0})_{i,j}^k \right),$$

where $(M_{a,\vartheta}^{S_0})_{i,j}^k$ denotes the partial derivative with respect to q_k of the entry of indexes (i, j) of the matrix $M_{a,\vartheta}^{S_0}$.

For any ϑ in \mathbb{R} , for any p in \mathbb{R}^3 , we define the matrix

$$S_\vartheta^{S_0}(p) := \left(\sum_{1 \leq i \leq 3} (\Gamma_\vartheta^{S_0})_{i,j}^k p_i \right)_{1 \leq k, j \leq 3} \quad \text{so that } \langle \Gamma_\vartheta^{S_0}, p, p \rangle = S_\vartheta^{S_0}(p)p.$$

Then, an explicit computation (another method, more theoretical, is given in the proof of Proposition 2.5, see Lemma 5.12) proves that for any ϑ in \mathbb{R} , for any p in \mathbb{R}^3 ,

$$\frac{1}{2} \frac{\partial M_{a,\vartheta}^{S_0}}{\partial q}(\vartheta) \cdot p - S_\vartheta^{S_0}(p) \text{ is skew-symmetric.} \quad (4.25)$$

Then the following energy conservation property straightforwardly follows from (4.25) and from the fact that the force $F_\vartheta^{S_0}(p)$ is gyroscopic in the sense of Definition 4.1.

Proposition 4.5. *Let q be a solution to (4.16), then the quantity $(M_g + M_{a,\vartheta}^{S_0}) q' \cdot q'$ is conserved.*

The quantity above corresponds to twice the sum of the kinetic energy of the solid associated to the genuine inertia and of the one associated to the added inertia.

Now in the remaining subsections will make use of the reformulation of the system established in Theorem 4.2 to investigate the passage to the limit of the dynamics when the size of the solid goes to 0^+ in both Cases (i) and (ii). We will establish the counterparts of Theorem 2.9 and Theorem 2.10.

4.5 Scaling with respect to ε

In this subsection we investigate the scaling of several objects with respect to ε in the absence of outer boundary. We will treat at once both Cases (i) and (ii). Recall that the scaling of the inertia is given by the relations (2.43) and (2.44), that is to say, Case (i) corresponds to $\alpha = 0$ in (2.44) while Case (ii) occurs when $\alpha > 0$. We consider for ε in $(0, 1]$, the scaled solid occupying initially the domain $\mathcal{S}_{0,\varepsilon}$ given by (2.35). We recall that γ and the Cauchy data $p_0 = (\omega_0, \ell_0)$ are supposed to be independent of ε whereas $q_0 = 0$. We denote by $q_\varepsilon := (\vartheta_\varepsilon, h_\varepsilon)$ in $\mathbb{R} \times \mathbb{R}^2$ the solution of the ODE obtained from (4.16) by rescaling the coefficients following the relations (2.35), (2.43) and (2.44).

Genuine inertia matrix and kinetic energy. Under the relation (2.43) and (2.44), the matrix of genuine inertia reads

$$M_{g,\varepsilon} = \varepsilon^\alpha I_\varepsilon M_g I_\varepsilon, \quad (4.26)$$

where we define

$$I_\varepsilon := \text{diag}(\varepsilon, 1, 1) \text{ and } M_g := \text{diag}(\mathcal{J}_1, m_1, m_1) \quad (4.27)$$

Therefore, it is natural to introduce the vector

$$p_\varepsilon := I_\varepsilon q'_\varepsilon = \begin{pmatrix} \varepsilon \vartheta'_\varepsilon \\ h'_\varepsilon \end{pmatrix}. \quad (4.28)$$

In particular the solid kinetic energy of the solid can be recast as

$$\frac{1}{2} M_{g,\varepsilon} q'_\varepsilon \cdot q'_\varepsilon = \frac{1}{2} \varepsilon^\alpha M_g p_\varepsilon \cdot p_\varepsilon. \quad (4.29)$$

Hence the natural counterpart to h'_ε for what concerns the angular velocity is rather $\varepsilon \vartheta'_\varepsilon$ than ϑ'_ε . This can also be seen on the boundary condition (2.1e): when x belongs to $\partial \mathcal{S}_\varepsilon(t)$, the term $\vartheta'_\varepsilon (x - h_\varepsilon)^\perp$ is of order $\varepsilon \vartheta'_\varepsilon$ and is added to h'_ε .

Added inertia matrix. We recall that $M_a^{\mathcal{S}_0}$ is the added inertia matrix for \mathcal{S}_0 defined in (4.8), while $M_{a,\vartheta}^{\mathcal{S}_0}$ is the one corresponding to $\mathcal{S}(q) = R(\vartheta)\mathcal{S}_0 + h$ (with $q = (\vartheta, h)$), given in (4.9). The scaled version of these matrices are $M_{a,\varepsilon}^{\mathcal{S}_0}$

corresponding to the solid $\mathcal{S}_{0,\varepsilon}$ and $M_{a,\vartheta,\varepsilon}^{\mathcal{S}_0}$ corresponding to $\mathcal{S}_\varepsilon(q)$. Then one easily sees after suitable scaling/rotation arguments that

$$M_{a,\varepsilon}^{\mathcal{S}_0} = \varepsilon^2 I_\varepsilon M_a^{\mathcal{S}_0} I_\varepsilon \text{ and } M_{a,\vartheta,\varepsilon}^{\mathcal{S}_0} = \varepsilon^2 I_\varepsilon M_{a,\vartheta}^{\mathcal{S}_0} I_\varepsilon. \quad (4.30)$$

So the dependence of the added inertia matrices with respect to ε is quite simple. It will not be the case any longer in the case Ω bounded, where this dependence will be much more intricate, see Proposition 9.1 below.

Other terms and scaled equation. The other terms in (4.16) have also a simple scaling with respect to ε in the unbounded case. Concerning the Christoffel symbols (4.11), it is not hard to check that $\langle \Gamma_{\vartheta,\varepsilon}^{\mathcal{S}_0}, q'_\varepsilon, q'_\varepsilon \rangle = \varepsilon I_\varepsilon \langle \Gamma_{\vartheta}^{\mathcal{S}_0}, p_\varepsilon, p_\varepsilon \rangle$, and concerning the force term (4.13) that $F_{\vartheta,\varepsilon}^{\mathcal{S}_0}(q'_\varepsilon) = I_\varepsilon F_{\vartheta}^{\mathcal{S}_0}(p_\varepsilon)$. The counterpart of Equation (4.16) for a shrinking solid is therefore

$$\left(\varepsilon^\alpha M_g + \varepsilon^2 M_{a,\vartheta,\varepsilon}^{\mathcal{S}_0} \right) p'_\varepsilon + \varepsilon \langle \Gamma_{\vartheta,\varepsilon}^{\mathcal{S}_0}, p_\varepsilon, p_\varepsilon \rangle = F_{\vartheta,\varepsilon}^{\mathcal{S}_0}(p_\varepsilon). \quad (4.31)$$

As we can see here, a difficulty is that the term $F_{\vartheta,\varepsilon}^{\mathcal{S}_0}$ depends on the unknown $\varepsilon\vartheta_\varepsilon$ through ϑ_ε , that is singularly. This difficulty, which is still present in the general case, will be overcome by using some averaging effect; see (4.36) and Lemma 4.8 below.

Rescaled energy. Let, for any ϑ in \mathbb{R} and for any ε in $(0, 1)$,

$$M_\vartheta(\varepsilon) := \begin{cases} M_g + \varepsilon^{2-\alpha} M_{a,\vartheta}^{\mathcal{S}_0} & \text{if } \alpha \leq 2, \\ M_{a,\vartheta}^{\mathcal{S}_0} + \varepsilon^{\alpha-2} M_g & \text{if } \alpha > 2. \end{cases} \quad (4.32)$$

Observe that $\varepsilon^\alpha M_g + \varepsilon^2 M_{a,\vartheta}^{\mathcal{S}_0} = \varepsilon^{\min(2,\alpha)} M_\vartheta(\varepsilon)$, and that this matrix is of order $O(1)$ with respect to ε . Using (4.28), (4.29) and (4.30) we obtain that the total energy is $\varepsilon^{\min(2,\alpha)} \mathcal{E}_\vartheta(\varepsilon, p_\varepsilon)$, where for any p in \mathbb{R}^3 ,

$$\mathcal{E}_\vartheta(\varepsilon, p) := \frac{1}{2} M_\vartheta(\varepsilon) p \cdot p. \quad (4.33)$$

The use of $\mathcal{E}_\vartheta(\varepsilon, p)$ is motivated by the following elementary result.

Lemma 4.6. *Suppose that \mathcal{S}_0 is not a disk or $\alpha \leq 2$. There exists $K > 0$ depending only on \mathcal{S}_0 , m_1 and \mathcal{J}_1 such that, for any $(\varepsilon, \vartheta, p)$ in $(0, 1) \times \mathbb{R} \times \mathbb{R}^3$,*

$$K |p|_{\mathbb{R}^3}^2 \leq \mathcal{E}_\vartheta(\varepsilon, p) \leq K^{-1} |p|_{\mathbb{R}^3}^2.$$

Using the energy conservation provided by Proposition 4.5 we deduce that, unless \mathcal{S}_0 is a disk and $\alpha > 2$, $(|p_\varepsilon|)_{\varepsilon \in (0,1)}$ is bounded uniformly on $[0, +\infty)$ in both Cases (i) and (ii). (Obtaining such an a priori estimate when Ω is a bounded domain will be much more involved in particular in Case (ii).)

In the degenerate case when \mathcal{S}_0 is a disk and $\alpha > 2$ (hence in Case (ii)), the problem is that $M_{a,\vartheta}^{\mathcal{S}_0}$ is both the principal part of $M_\vartheta(\varepsilon)$ and degenerate. However using (5.66) below, one can check that for $p = (\omega, \ell)$,

$$M_{a,\vartheta}^{\mathcal{S}_0} p \cdot p = M_b^{\mathcal{S}_0} (\ell + \omega \zeta_\vartheta^\perp) \cdot (\ell + \omega \zeta_\vartheta^\perp),$$

where $M_b^{S_0}$ is given in (4.23). We obtain that

$$\mathcal{E}_\vartheta(\varepsilon, p) = \frac{1}{2} \left(\varepsilon^{\alpha-2} m_1 |h'_\varepsilon|^2 + \varepsilon^\alpha \mathcal{J}_1 |\vartheta'_\varepsilon|^2 + M_b^{S_0} h'_{c,\varepsilon} \cdot h'_{c,\varepsilon} \right).$$

We deduce in this case that $(|h'_{c,\varepsilon}|)_{\varepsilon \in (0,1)}$, $(\varepsilon^{\frac{\alpha}{2}-1} |h'_\varepsilon|)_{\varepsilon \in (0,1)}$ and $(\varepsilon^{\frac{\alpha}{2}} |\vartheta'_\varepsilon|)_{\varepsilon \in (0,1)}$ are bounded uniformly on $[0, +\infty)$.

Now our goal is to pass to the limit in each term of (4.31). We distinguish Case (i) (with m_ε , \mathcal{J}_ε given by (2.37)) and Case (ii) (with m_ε , \mathcal{J}_ε given by (2.43)).

4.6 Case (i) without external boundary

In Case (i) we show that h_ε converges to the solution of a massive point vortex equation similar to (2.38) with a vanishing Kirchhoff-Routh velocity. For $\gamma \neq 0$ we let for $t \geq 0$,

$$\bar{h}(t) := \frac{m}{\gamma} \left[R \left(\frac{\gamma t}{m} - \frac{\pi}{2} \right) \ell_0 + \ell_0^\perp \right],$$

and $\bar{h}(t) := \ell_0 t$ for $\gamma = 0$. Of course \bar{h} satisfies

$$m \bar{h}'' = \gamma (\bar{h}')^\perp \quad \text{for } t \geq 0, \quad \text{with } \bar{h}(0) = 0 \quad \text{and } \bar{h}'(0) = \ell_0. \quad (4.34)$$

The precise statement of our first convergence result is as follows.

Theorem 4.7. *Let \mathcal{S}_0 a subset of \mathbb{R}^2 , p_0 in \mathbb{R}^3 and (γ, m, \mathcal{J}) in $\mathbb{R} \times (0, +\infty) \times (0, +\infty)$. Let \bar{h} be the global solution of (4.34). For every ε in $(0, 1]$, let q_ε be the global solution of (4.28) and (4.31) where m_ε and \mathcal{J}_ε are given by (2.37), and with the initial data $q_\varepsilon(0) = 0$ and $q'_\varepsilon(0) = p_0$. Then, as $\varepsilon \rightarrow 0^+$, for all $T > 0$, $h_\varepsilon \rightharpoonup \bar{h}$ in $W^{2,\infty}([0, T]; \mathbb{R}^2)$ weak- \star and $\varepsilon \vartheta_\varepsilon \rightharpoonup 0$ in $W^{2,\infty}([0, T]; \mathbb{R})$ weak- \star .*

Proof Let $T > 0$. Using that $(|p_\varepsilon|)_{\varepsilon \in (0,1)}$ is bounded uniformly on $[0, T]$ and the equation (4.31), we deduce some uniform $W^{2,\infty}$ bounds on h_ε and $\varepsilon \vartheta_\varepsilon$. This entails the existence of a converging subsequence $(h_{\varepsilon_n}, \varepsilon_n \vartheta_{\varepsilon_n})$ of $(h_\varepsilon, \varepsilon \vartheta_\varepsilon)$:

$$(h_{\varepsilon_n}, \varepsilon_n \vartheta_{\varepsilon_n}) \rightharpoonup (h_*, \Theta_*) \quad \text{in } W^{2,\infty} \text{ weak-}\star. \quad (4.35)$$

We now aim at characterizing the limit. First it is clear that the left hand side of (4.31) converges to $M_g(\Theta_*'', h_*'')^t$ in L^∞ weak- \star . Now consider the force term $F_\vartheta^{S_0}(p_{\varepsilon_n}) = \gamma(\zeta_{\vartheta_{\varepsilon_n}} \cdot h'_{\varepsilon_n}, (h'_{\varepsilon_n})^\perp - \varepsilon_n \vartheta'_{\varepsilon_n} \zeta_{\vartheta_{\varepsilon_n}})^t$ as defined in (4.13). On the one hand using (4.14) we see that

$$\varepsilon_n \vartheta'_{\varepsilon_n} \zeta_{\vartheta_{\varepsilon_n}} \rightharpoonup 0 \quad \text{in } W^{-1,\infty} \text{ weak-}\star. \quad (4.36)$$

On the other hand since the weak- \star convergence in $W^{2,\infty}$ entails the strong $W^{1,\infty}$ one, we get that $h'_{\varepsilon_n} \rightarrow h'_*$ in L^∞ and $(h_*(0), h'_*(0)) = (0, \ell_0)$. Hence

we infer from the two last components of System (4.31) that $m_1 h_*'' = \gamma [h_*']^\perp$. Due to the uniqueness of the solution of (4.34), this proves that $h_* = \bar{h}$ and the convergence as $\varepsilon \rightarrow 0^+$ of the whole sequence (not merely the convergence along a subsequence). This concludes the proof of Theorem 4.7 for the part concerning the position of the center of mass.

We now turn to the part concerning the angle, that is the convergence of $\varepsilon \vartheta_\varepsilon$. We will use the following lemma (see [9]).

Lemma 4.8. *Let $(\omega_n)_{n \in \mathbb{N}}$ in $W^{1,\infty}(0, T; \mathbb{R})^\mathbb{N}$, $(\varepsilon_n)_{n \in \mathbb{N}}$ in $(0, +\infty)^\mathbb{N}$ with $\varepsilon_n \rightarrow 0^+$ as $n \rightarrow +\infty$, such that $\varepsilon_n \omega_n \rightharpoonup \bar{\rho}$ in $W^{1,\infty}(0, T; \mathbb{R})$ weak- \star as $n \rightarrow +\infty$. Let $(w_n)_{n \in \mathbb{N}}$ in $L^\infty(0, T; \mathbb{C})^\mathbb{N}$ such that $w_n \rightarrow w$ in $L^\infty(0, T; \mathbb{C})$ as $n \rightarrow +\infty$. Let $\vartheta_n := \int_0^t \omega_n$. Suppose that, on $(0, T)$, $\varepsilon_n \omega_n'(t) = \operatorname{Re} [w_n(t) \exp(-i\vartheta_n(t))]$. Then $\bar{\rho}$ is constant on $[0, T]$.*

We consider the first component of the system (4.31). We apply Lemma 4.8 to $\omega_n = \vartheta_{\varepsilon_n}'$, $\vartheta_n = \vartheta_{\varepsilon_n}$ and $w_n := \frac{\gamma}{\mathcal{J}_1} [(\zeta \cdot h_{\varepsilon_n}') - i(\zeta \cdot (h_{\varepsilon_n}')^\perp)]$. The assumptions are satisfied thanks to (4.35). Using the initial data, we infer that $\Theta_* = 0$. This concludes the proof of Theorem 4.7. \square

4.7 Case (ii) without external boundary

In this situation we show that the limit system is the point vortex system with a vanishing Kirchhoff-Routh velocity field, that is the trivial system $\bar{h}' = 0$. Our result is the following.

Theorem 4.9. *Let S_0 a subset of \mathbb{R}^2 different from a disk (or $\alpha \leq 2$), $\gamma \neq 0$, p_0 in \mathbb{R}^3 and (γ, m, \mathcal{J}) in $\mathbb{R} \times (0, +\infty) \times (0, +\infty)$. For every ε in $(0, 1]$, let q_ε be the global solution of (4.28) and (4.31) where m_ε and \mathcal{J}_ε are given by (2.43), and with the initial data $q_\varepsilon(0) = 0$ and $q_\varepsilon'(0) = p_0$. Then, as $\varepsilon \rightarrow 0^+$, $h_\varepsilon \rightharpoonup 0$ in $W^{1,\infty}([0, T]; \mathbb{R}^2)$ weak- \star for all $T > 0$.*

Proof First, thanks to the energy estimate, $\varepsilon M_{a, \vartheta_\varepsilon}^{S_0}$ is bounded in $W^{1,\infty}$. Since moreover M_g is constant and $(p_\varepsilon)'$ is bounded in $W^{-1,\infty}$, we can conclude that the left hand side of (4.31) converges to 0 in $W^{-1,\infty}$ (due to the extra powers of ε). Next, concerning the right hand side, the term $\varepsilon \langle \Gamma_{\vartheta_\varepsilon}^{S_0}, p_\varepsilon, p_\varepsilon \rangle$ converges to 0 in L^∞ since the terms inside the brackets are bounded. Now let us consider the remaining terms in the two last lines of the equation (4.31), that is, $\gamma (h_\varepsilon')^\perp - \gamma \varepsilon \vartheta_\varepsilon' \zeta_{\vartheta_\varepsilon}$. The last term converges weakly to 0 in $W^{-1,\infty}$ as seen in Case (i), see (4.36). Hence we infer that h_ε' converges weakly- \star to 0 in $W^{-1,\infty}$. Due to the a priori estimate, this convergence occurs in L^∞ weak- \star . Again this is sufficient to deduce the strong convergence of h_ε towards some h_* in L^∞ , and that $h_*' = 0$ and $h_*(0) = 0$. \square

In the case of a non-homogeneous disk (or more precisely a disk for which the center of gravity is not at the center of the disk) and for $\alpha > 2$, we have the equivalent convergence for the center of the disk $h_{c,\varepsilon}$. It suffices to pass to the weak limit in (2.34a) where the right hand side is simplified as in (4.16).

Note that in that case \tilde{M}_b is constant, so $\Gamma_b = 0$ (while the “original” Γ is not, since M_a has one more dimension and a different set of variables). We omit the details.

5 Recasting the system: Proofs of Theorem 2.2, Lemma 2.4, Proposition 2.5 and Proposition 2.7

In this Section, we prove the results of Subsection 2.1 concerning the dynamics of a solid with fixed size and mass.

5.1 Splitting the proof of Theorem 2.2

The pressure Π can be recovered by means of Bernoulli's formula which is obtained by combining (2.1a) and (2.2), and which reads:

$$\nabla \Pi = - \left(\frac{\partial u}{\partial t} + \frac{1}{2} \nabla |u|^2 \right) \quad \text{in } \mathcal{F}(q). \quad (5.1)$$

Given q , p and γ , the pair (u, Π) where u is given by (2.13) and Π by (5.1) yields a solution to (2.1a-d).

Now, equations (2.1g-h) can be summarized in the variational form:

$$mh'' \cdot \ell^* + \mathcal{J}\vartheta''\omega^* = \int_{\partial S(q)} \Pi(\omega^*(x-h)^\perp + \ell^*) \cdot n \, ds, \quad (5.2)$$

for all $p^* := (\omega^*, \ell^*)$ in $\mathbb{R} \times \mathbb{R}^2$. Let us associate with (q, p^*) in $\mathcal{Q} \times \mathbb{R}^3$ the potential vector field

$$u^* := \nabla(\varphi(q, \cdot) \cdot p^*), \quad (5.3)$$

which is defined on $\mathcal{F}(q)$. According to Bernoulli's formula (5.1) and upon an integration by parts, identity (5.2) can be turned into:

$$mh'' \cdot \ell^* + \mathcal{J}\vartheta''\omega^* = - \int_{\mathcal{F}(q)} \left(\frac{\partial u}{\partial t} + \frac{1}{2} \nabla |u|^2 \right) \cdot u^* \, dx, \quad (5.4)$$

for all $p^* := (\omega^*, \ell^*)$ in $\mathbb{R} \times \mathbb{R}^2$. Therefore plugging the decomposition (2.13) into (5.4) leads to

$$\begin{aligned} mh'' \cdot \ell^* + \mathcal{J}\vartheta''\omega^* + \int_{\mathcal{F}(q)} \left(\frac{\partial u_{q'}}{\partial t} + \frac{1}{2} \nabla |u_{q'}|^2 \right) \cdot u^* \, dx = \\ - \int_{\mathcal{F}(q)} \left(\frac{1}{2} \nabla |u_\gamma|^2 \right) \cdot u^* \, dx - \int_{\mathcal{F}(q)} \left(\frac{\partial u_\gamma}{\partial t} + \frac{1}{2} \nabla (u_{q'} \cdot u_\gamma) \right) \cdot u^* \, dx, \end{aligned} \quad (5.5)$$

for all $p^* := (\omega^*, \ell^*)$ in $\mathbb{R} \times \mathbb{R}^2$.

Then the reformulation of Equations (2.1g-h) mentioned in Theorem 2.2 will follow from the three following lemmas which deal respectively with the left hand side of (5.5) and the two terms in the right hand side.

Lemma 5.1. *For any smooth curve $t \mapsto q(t)$ in \mathcal{Q} and every $p^* = (\omega^*, \ell^*)$ in \mathbb{R}^3 , the following identity holds:*

$$\begin{aligned} mh'' \cdot \ell^* + \mathcal{J}\vartheta''\omega^* + \int_{\mathcal{F}(q)} \left(\frac{\partial u_{q'}}{\partial t} + \frac{1}{2} \nabla |u_{q'}|^2 \right) \cdot u^* dx \\ = M(q)q'' \cdot p^* + \langle \Gamma(q), q', q' \rangle \cdot p^*, \end{aligned} \quad (5.6)$$

where $q = (\vartheta, h)$, u^* is given by (5.3), $u_{q'}$ is given by (2.14), $M(q)$ and $\Gamma(q)$ are defined in (2.15) and (2.16).

Lemma 5.2. *For every $q \in \mathcal{Q}$ and every $p^* = (\omega^*, \ell^*) \in \mathbb{R}^3$, the following identity holds:*

$$- \int_{\mathcal{F}(q)} \left(\frac{1}{2} \nabla |u_\gamma|^2 \right) \cdot u^* dx = \gamma^2 E(q) \cdot p^*, \quad (5.7)$$

where u^* is given by (5.3), $u_{q'}$ is given by (2.14), $E(q)$ is defined in (2.17).

Lemma 5.3. *For any smooth curve $t \mapsto q(t)$ in \mathcal{Q} and every $p^* = (\omega^*, \ell^*)$ in \mathbb{R}^3 , the following identity holds:*

$$- \int_{\mathcal{F}(q)} \left(\frac{\partial u_\gamma}{\partial t} + \nabla(u_{q'} \cdot u_\gamma) \right) \cdot u^* dx = \gamma(q' \times B(q)) \cdot p^*, \quad (5.8)$$

where $q = (\vartheta, h)$, u^* is given by (5.3), u_γ and $u_{q'}$ are given by (2.14), $B(q)$ is defined in (2.17).

Lemma 5.2 simply follows from an integration by parts. Let us consider Lemmas 5.1 and 5.3 as granted. Then gathering the results of Lemmas 5.1, 5.2 and 5.3 with (5.5), the conclusion of Theorem 2.2 follows. \square

5.2 Reformulation of the potential part: Proof of Lemma 5.1

We start with observing that

$$mh'' \cdot \ell^* + \mathcal{J}\vartheta''\omega^* = M_g q'' \cdot p^*. \quad (5.9)$$

Now in order to deal with the last term of the left hand side of (5.6) we use a Lagrangian strategy. For any q in \mathcal{Q} and every p in \mathbb{R}^3 , let us denote

$$\mathcal{E}_1(q, p) := \frac{1}{2} \int_{\mathcal{F}(q)} |\nabla(\varphi(q, \cdot) \cdot p)|^2 dx. \quad (5.10)$$

Thus $\mathcal{E}_1(q, q')$ denotes the kinetic energy of the potential part $u_{q'}$ of the flow. It follows from classical shape derivative theory that \mathcal{E}_1 is in $C^\infty(\mathcal{Q} \times \mathbb{R}^3; [0, +\infty))$.

Now the crucial quantity here is the Euler-Lagrange function:

$$\mathcal{E}\mathcal{L} = \left(\frac{d}{dt} \frac{\partial \mathcal{E}_1}{\partial p}(q, q') - \frac{\partial \mathcal{E}_1}{\partial q}(q, q') \right) \cdot p^*, \quad (5.11)$$

associated with any smooth curve $t \mapsto q(t)$ in \mathcal{Q} and any $p^* \in \mathbb{R}^3$.

Lemma 5.4. *For any smooth curve $t \mapsto q(t)$ in \mathcal{Q} , for every $p^* \in \mathbb{R}^3$, we have:*

$$\int_{\mathcal{F}(q)} \left(\frac{\partial u_{q'}}{\partial t} + \frac{1}{2} \nabla |u_{q'}|^2 \right) \cdot u^* dx = \mathcal{E}\mathcal{L} \quad (5.12)$$

where $u_{q'}$ is given by (2.14), u^* is given by (5.3) and $\mathcal{E}\mathcal{L}$ is given by (5.11).

Proof Let us make use of the following slight abuse of notations which simplifies the presentation of the proof of Lemma 5.4. For a smooth function $I(q, p)$, where (q, p) is running into $\mathcal{Q} \times \mathbb{R}^3$, and a smooth curve $t \mapsto q(t)$ in \mathcal{Q} let us denote

$$\left(\frac{\partial}{\partial q} \frac{d}{dt} I(q, q') \right) (t) := \left(\frac{\partial}{\partial q} J \right) (q(t), q'(t), q''(t)),$$

where, for (q, p, r) in $\mathcal{Q} \times \mathbb{R}^3 \times \mathbb{R}^3$,

$$J(q, p, r) = p \frac{\partial I}{\partial q}(q, p) + r \frac{\partial I}{\partial p}(q, p). \quad (5.13)$$

Observe in particular that

$$\frac{d}{dt} (I(q(t), q'(t))) = J(q(t), q'(t), q''(t)),$$

and

$$\frac{d}{dt} \left(\frac{\partial I}{\partial q}(q(t), q'(t)) \right) = \left(\frac{\partial}{\partial q} \frac{d}{dt} I(q, q') \right) (t). \quad (5.14)$$

Below, in such circumstances, it will be comfortable to write

$$\frac{\partial}{\partial q} [J(q(t), q'(t), q''(t))] \text{ instead of } \left(\frac{\partial J}{\partial q} \right) (q(t), q'(t), q''(t)),$$

and it will be understood that J is extended from $(q(t), q'(t), q''(t))$ to general (q, p, r) by (5.13).

We start with manipulating the first term of $\mathcal{E}\mathcal{L}$. An integration by parts leads to:

$$\frac{\partial \mathcal{E}_1}{\partial p} \cdot p^* = \int_{\mathcal{F}(q)} u_{q'} \cdot u^* dx = \int_{\partial \mathcal{S}(q)} (\varphi \cdot q') (u^* \cdot n) ds.$$

Then, invoking the Reynolds transport theorem (see [35, pages 12–13]), we get:

$$\frac{\partial \mathcal{E}_1}{\partial p} \cdot p^* = \frac{\partial}{\partial q} \left(\int_{\mathcal{F}(q)} (\varphi \cdot q') dx \right) \cdot p^* - \int_{\mathcal{F}(q)} \left(\frac{\partial \varphi}{\partial q} \cdot q' \right) \cdot p^* dx. \quad (5.15)$$

Differentiating (5.15) with respect to t , we obtain:

$$\begin{aligned} \frac{d}{dt} \frac{\partial \mathcal{E}_1}{\partial p} \cdot p^* &= \frac{d}{dt} \frac{\partial}{\partial q} \left(\int_{\mathcal{F}(q)} (\varphi \cdot q') \, dx \right) \cdot p^* \\ &\quad - \frac{d}{dt} \left(\int_{\mathcal{F}(q)} \left(\frac{\partial \varphi}{\partial q} \cdot q' \right) \cdot p^* \, dx \right). \end{aligned} \quad (5.16)$$

With the abuse of notations mentioned above we exchange the derivatives involved in the first term of the right hand side, so that the identity (5.16) can be rewritten as follows:

$$\begin{aligned} \frac{d}{dt} \frac{\partial \mathcal{E}_1}{\partial p} \cdot p^* &= \frac{\partial}{\partial q} \frac{d}{dt} \left(\int_{\mathcal{F}(q)} (\varphi \cdot q') \, dx \right) \cdot p^* \\ &\quad - \frac{d}{dt} \left(\int_{\mathcal{F}(q)} \left(\frac{\partial \varphi}{\partial q} \cdot q' \right) \cdot p^* \, dx \right). \end{aligned} \quad (5.17)$$

Moreover, using again the Reynolds transport theorem, we have:

$$\begin{aligned} \frac{d}{dt} \left(\int_{\mathcal{F}(q)} (\varphi \cdot q') \, dx \right) &= \int_{\mathcal{F}(q)} \frac{\partial}{\partial t} (\varphi \cdot q') \, dx + \int_{\partial \mathcal{S}(q)} (\varphi \cdot q') (u_{q'} \cdot n) \, ds \\ &= \int_{\mathcal{F}(q)} \frac{\partial}{\partial t} (\varphi \cdot q') \, dx + 2\mathcal{E}_1(q, q'), \end{aligned} \quad (5.18)$$

by integration by parts.

We infer from (5.17) and (5.18), again with the abuse of notations mentioned above, that:

$$\begin{aligned} \mathcal{E}\mathcal{L} &= \frac{\partial \mathcal{E}_1}{\partial q} + \frac{\partial}{\partial q} \left[\int_{\mathcal{F}(q)} \frac{\partial}{\partial t} (\varphi \cdot q') \, dx \right] \cdot p^* \\ &\quad - \frac{d}{dt} \left(\int_{\mathcal{F}(q)} \left(\frac{\partial \varphi}{\partial q} \cdot q' \right) \cdot p^* \, dx \right). \end{aligned} \quad (5.19)$$

Thanks to the Reynolds transport theorem, we get for the second term of the right hand side

$$\begin{aligned} \frac{\partial}{\partial q} \left[\int_{\mathcal{F}(q)} \frac{\partial}{\partial t} (\varphi \cdot q') \, dx \right] \cdot p^* &= \int_{\mathcal{F}(q)} \frac{\partial}{\partial q} \left(\frac{\partial}{\partial t} (\varphi \cdot q') \right) \cdot p^* \, dx \\ &\quad + \int_{\partial \mathcal{S}(q)} \frac{\partial}{\partial t} (\varphi \cdot q') (u^* \cdot n) \, ds, \end{aligned} \quad (5.20)$$

and for the last one:

$$\begin{aligned} \frac{d}{dt} \left(\int_{\mathcal{F}(q)} \left(\frac{\partial \varphi}{\partial q} \cdot q' \right) \cdot p^* dx \right) &= \int_{\mathcal{F}(q)} \frac{\partial}{\partial t} \left(\left(\frac{\partial \varphi}{\partial q} \cdot q' \right) \cdot p^* \right) dx \\ &\quad + \int_{\partial \mathcal{S}(q)} \left(\left(\frac{\partial \varphi}{\partial q} \cdot q' \right) \cdot p^* \right) (u_{q'} \cdot n) ds. \end{aligned} \quad (5.21)$$

Using again (5.14) for the first term and integrating by parts the second one, we obtain:

$$\begin{aligned} \frac{d}{dt} \left(\int_{\mathcal{F}(q)} \left(\frac{\partial \varphi}{\partial q} \cdot q' \right) \cdot p^* dx \right) \\ = \int_{\mathcal{F}(q)} \frac{\partial}{\partial q} \left(\frac{\partial}{\partial t} (\varphi \cdot q') \right) \cdot p^* dx + \int_{\mathcal{F}(q)} \left(\frac{\partial u_{q'}}{\partial q} \cdot p^* \right) \cdot u_{q'} dx. \end{aligned} \quad (5.22)$$

On the other hand, again using the Reynolds transport theorem, we have:

$$\frac{\partial \mathcal{E}_1}{\partial q} \cdot p^* = \int_{\mathcal{F}(q)} \left(\frac{\partial u_{q'}}{\partial q} \cdot p^* \right) \cdot u_{q'} dx + \frac{1}{2} \int_{\partial \mathcal{S}(q)} |u_{q'}|^2 (u^* \cdot n) ds. \quad (5.23)$$

Plugging the expressions (5.23), (5.20) and (5.22) into (5.19) and simplifying, we end up with:

$$\mathcal{E} \mathcal{L} = \int_{\partial \mathcal{S}(q)} \left[\frac{\partial}{\partial t} (\varphi \cdot q') + \frac{1}{2} |u_{q'}|^2 \right] (u^* \cdot n) ds.$$

Upon an integration by parts, we recover (5.12) and the proof of Lemma 5.4 is completed. \square

Now, we observe that the kinetic energy $\mathcal{E}_1(q, q')$ associated to the potential part of the flow, as defined by (5.10), can be rewritten as:

$$\mathcal{E}_1(q, q') = \frac{1}{2} M_a(q) q' \cdot q', \quad (5.24)$$

where $M_a(q)$ is defined by (2.15). This allows us to prove the following result.

Lemma 5.5. *For any smooth curve $t \mapsto q(t)$ in \mathcal{Q} , for every $p^* \in \mathbb{R}^3$, we have:*

$$\mathcal{E} \mathcal{L} = M_a(q) q'' \cdot p^* + \langle \Gamma(q), q', q' \rangle \cdot p^*, \quad (5.25)$$

with $\mathcal{E} \mathcal{L}$ is given by (5.11), $M_a(q)$ defined by (2.15) and $\Gamma(q)$ associated with $M(q)$ by the Christoffel formula (2.16a)-(2.16b).

Proof of Lemma 5.5. Using (5.24) in the definition (5.11) of \mathcal{EL} we have

$$\mathcal{EL} = M_a(q)q'' \cdot p^* + \left((DM_a(q) \cdot q')q' \right) \cdot p^* - \frac{1}{2} \left((DM_a(q) \cdot p^*)q' \right) \cdot q'.$$

Let us recall the notation $(M_a)_{i,j}^k(q)$ in (2.16c) and let the notation \sum stand for $\sum_{1 \leq i,j,k \leq 3}$ for the rest of this proof. Then

$$\mathcal{EL} = M_a q'' \cdot p^* + \sum (M_a)_{i,j}^k q'_k q'_j p_i^* - \frac{1}{2} \sum (M_a)_{i,j}^k q'_i q'_j p_k^*.$$

A symmetrization of the second term of the right hand side above leads to

$$\begin{aligned} \mathcal{EL} &= M_a q'' \cdot p^* \\ &+ \frac{1}{2} \left(\sum (M_a)_{i,j}^k q'_k q'_j p_i^* + \sum (M_a)_{i,k}^j q'_k q'_j p_i^* - \sum (M_a)_{i,j}^k q'_i q'_j p_k^* \right), \end{aligned}$$

Then the result follows by exchanging i and k in the last sum. \square

Finally Lemma 5.1 straightforwardly results from the combination of (5.9), Lemmas 5.4 and 5.5. This concludes the proof of Lemma 5.1.

5.3 Reformulation of the cross part: Proof of Lemma 5.3

Write $q := (\vartheta, h)$ and recall that

$$u_\gamma := \gamma \nabla^\perp \psi(q, \cdot), \quad u_{q'} := \nabla(\varphi(q, \cdot) \cdot q') \text{ and } u^* := \nabla(\varphi(q, \cdot) \cdot p^*).$$

We first observe that

$$\int_{\mathcal{F}(q)} \left(\frac{\partial u_\gamma}{\partial t} \right) \cdot u^* dx = -\gamma \int_{\partial \mathcal{S}(q)} \left(\frac{\partial}{\partial t} (\psi(q, \cdot)) \right) \left(\frac{\partial \varphi}{\partial \tau} (q, \cdot) \cdot p^* \right) ds. \quad (5.26)$$

Let us emphasize that there is no contribution on $\partial \Omega$ since $\psi(q, \cdot)$ is vanishing there. Now we have the following result.

Lemma 5.6. *On $\partial \mathcal{S}(q)$, we have*

$$\frac{\partial}{\partial t} (\psi(q, \cdot)) = -\frac{\partial \psi}{\partial n} (q, \cdot) \left(\frac{\partial \varphi}{\partial n} (q, \cdot) \cdot q' \right) + DC(q) \cdot q'. \quad (5.27)$$

Proof We start with the observation that

$$\frac{\partial}{\partial t} (\psi(q, \cdot)) = \frac{\partial \psi}{\partial q} (q, \cdot) \cdot q', \quad (5.28)$$

is the derivative of the function $\psi(q, \cdot)$ when the boundary $\partial \mathcal{S}(q)$ undergoes a rigid displacement of velocity $w = \omega(x - h)^\perp + \ell$ where $q' = (\omega, \ell)$. Then we differentiate the identity:

$$\psi(q, R(\vartheta)X + h) = C(q), \quad \text{for } X \in \partial \mathcal{S}_0,$$

with respect to q in the direction q' . We obtain:

$$\frac{\partial \psi}{\partial q}(q, x) \cdot q' + \nabla \psi(q, x) \cdot w = DC(q) \cdot q', \quad \text{for } x \in \partial \mathcal{S}(q). \quad (5.29)$$

Since $\psi(q, \cdot)$ is constant on $\partial \mathcal{S}(q)$, its tangential derivative is zero. Besides, on $\partial \mathcal{S}(q)$ we have $w \cdot n = u_{q'} \cdot n = \frac{\partial \varphi}{\partial n}(q, \cdot) \cdot q'$ so we get

$$\nabla \psi(q, x) \cdot w = \frac{\partial \psi}{\partial n}(q, x) \left(\frac{\partial \varphi}{\partial n}(q, x) \cdot q' \right) \quad \text{for } x \in \partial \mathcal{S}(q). \quad (5.30)$$

Gathering (5.28), (5.29) and (5.30) we obtain (5.27). \square

Plugging now (5.27) into (5.26) we deduce that:

$$\int_{\mathcal{F}(q)} \left(\frac{\partial u_\gamma}{\partial t} \right) \cdot u^* dx = \gamma \int_{\partial \mathcal{S}(q)} \frac{\partial \psi}{\partial n} \left(\frac{\partial \varphi}{\partial n} \cdot q' \right) \left(\frac{\partial \varphi}{\partial \tau} \cdot p^* \right) ds. \quad (5.31)$$

On the other hand, integrating by parts and using that $u_\gamma = -\gamma \frac{\partial \psi}{\partial n} \tau$ on $\partial \mathcal{S}(q)$, we get:

$$\int_{\mathcal{F}(q)} \nabla(u_{q'} \cdot u_\gamma) \cdot u^* dx = \gamma \int_{\partial \mathcal{S}(q)} (u_{q'} \cdot u_\gamma)(u^* \cdot n) ds. \quad (5.32)$$

Adding (5.31) and (5.32), we get:

$$\begin{aligned} \int_{\mathcal{F}(q)} \left(\frac{\partial u_\gamma}{\partial t} + \nabla(u_{q'} \cdot u_\gamma) \right) \cdot u^* dx = \\ \gamma \int_{\partial \mathcal{S}(q)} \frac{\partial \psi}{\partial n} \left[\left(\frac{\partial \varphi}{\partial n} \cdot q' \right) \left(\frac{\partial \varphi}{\partial \tau} \cdot p^* \right) - \left(\frac{\partial \varphi}{\partial n} \cdot p^* \right) \left(\frac{\partial \varphi}{\partial \tau} \cdot q' \right) \right] ds, \end{aligned}$$

which is (5.8). This ends the proof of Lemma 5.3.

5.4 Decomposition of the Christoffel symbols: Proof of Proposition 2.7

This section is devoted to the proof of Proposition 2.7. We will use the matrix $\underline{M}_a(q)$ given by

$$\underline{M}_a(q) := \mathcal{R}(\vartheta)^t M_a(q) \mathcal{R}(\vartheta), \quad (5.33)$$

where we recall that $\mathcal{R}(\vartheta)$ is the rotation matrix defined by (4.10). We also introduce the following real-valued functions depending on the variables $q = (\vartheta, h) \in \mathcal{Q}$, p in \mathbb{R}^3 and p^* in \mathbb{R}^3 :

$$\begin{aligned} \Xi_1(q, p, p^*) &:= \left[\left(\frac{\partial \underline{M}_a}{\partial q}(q) \cdot p^* \right) \mathcal{R}(\vartheta)^t p \right] \cdot \mathcal{R}(\vartheta)^t p, \\ \Xi_3(q, p, p^*) &:= \left[\left(\frac{\partial \underline{M}_a}{\partial q}(q) \cdot p \right) \mathcal{R}(\vartheta)^t p \right] \cdot \mathcal{R}(\vartheta)^t p^*. \end{aligned}$$

The indices in Ξ_1 and Ξ_3 above are chosen in order to recall the position where p^* appears (the “first p ” in the expression of Ξ_1 is p^* and so on.) Similarly we define, for $p = (\omega, \ell)$ and $p^* = (\omega^*, \ell^*)$ the functions:

$$\begin{aligned} \Upsilon_1(q, p, p^*) &:= \omega^* M_a(q) p \cdot \begin{pmatrix} 0 \\ \ell^\perp \end{pmatrix}, & \Upsilon_2(q, p, p^*) &:= \omega M_a(q) p^* \cdot \begin{pmatrix} 0 \\ \ell^\perp \end{pmatrix}, \\ & & \text{and } \Upsilon_3(q, p, p^*) &:= \omega M_a(q) p \cdot \begin{pmatrix} 0 \\ \ell^{*\perp} \end{pmatrix}. \end{aligned}$$

The proof of Proposition 2.7 is then split into the proof of the following three lemmas.

Lemma 5.7. *For any (q, p, p^*) in $\mathcal{Q} \times \mathbb{R}^3 \times \mathbb{R}^3$, we have:*

$$\begin{aligned} \langle \Gamma(q), p, p \rangle \cdot p^* &= \Upsilon_1(q, p, p^*) - \Upsilon_2(q, p, p^*) - \Upsilon_3(q, p, p^*) \\ &\quad + \Xi_3(q, p, p^*) - \frac{1}{2} \Xi_1(q, p, p^*). \end{aligned} \quad (5.34)$$

Lemma 5.8. *For any (q, p, p^*) in $\mathcal{Q} \times \mathbb{R}^3 \times \mathbb{R}^3$, we have:*

$$\Upsilon_1(q, p, p^*) - \Upsilon_2(q, p, p^*) - \Upsilon_3(q, p, p^*) = \langle \Gamma^{\text{rot}}(q), p, p \rangle \cdot p^*,$$

where $\Gamma^{\text{rot}}(q)$ is defined in (2.25).

Lemma 5.9. *For any (q, p, p^*) in $\mathcal{Q} \times \mathbb{R}^3 \times \mathbb{R}^3$, we have:*

$$\Xi_3(q, p, p^*) - \frac{1}{2} \Xi_1(q, p, p^*) = \langle \Gamma^{\partial\Omega}(q), p, p \rangle \cdot p^*, \quad (5.35)$$

where $\Gamma^{\partial\Omega}(q)$ is defined in (2.26).

Before proving these three lemmas, we introduce a few notations. For every $q = (\vartheta, h) \in \mathcal{Q}$, we define the change of variables $y = R(\vartheta)^t(x-h)$, the domains

$$\underline{\Omega}(q) := R(\vartheta)^t(\Omega - h), \quad \underline{\mathcal{F}}(q) := R(\vartheta)^t(\mathcal{F}(q) - h) = \underline{\Omega}(q) \setminus \mathcal{S}_0,$$

and the functions $\underline{\varphi}_i(q, y)$ such that, denoting $\underline{\varphi}(q, \cdot) := (\underline{\varphi}_1(q, \cdot), \underline{\varphi}_2(q, \cdot), \underline{\varphi}_3(q, \cdot))$, we have

$$\underline{\varphi}(q, y) := \mathcal{R}(\vartheta)^t \varphi(q, x), \quad y \in \underline{\mathcal{F}}(q).$$

For every $j = 1, 2, 3$, the functions $\underline{\varphi}_j(q, \cdot)$ are harmonic in $\underline{\mathcal{F}}(q)$ and satisfy:

$$\frac{\partial \underline{\varphi}_j}{\partial n}(q, y) = \begin{cases} y^\perp \cdot n & j = 1; \\ n_{j-1} & j = 2, 3 \end{cases} \quad \text{on } \partial \mathcal{S}_0, \quad (5.36a)$$

$$\frac{\partial \underline{\varphi}_j}{\partial n}(q, y) = 0 \quad (j = 1, 2, 3) \quad \text{on } \partial \underline{\Omega}(q). \quad (5.36b)$$

Therefore the matrix $\underline{M}_a(q)$ defined in (5.33) can be recast as

$$\underline{M}_a(q) = \int_{\partial \mathcal{S}_0} \underline{\varphi}(q) \otimes \frac{\partial \underline{\varphi}}{\partial n}(q) \, ds. \quad (5.37)$$

Proof of Lemma 5.7. Let $t \mapsto q(t)$ be a smooth curve in \mathcal{Q} , defined in a neighborhood of 0 such that $q(0) = q$ and $q'(0) = p$. For every $p^* = (\omega^*, \ell^*) \in \mathbb{R}^3$, using the expression of $\mathcal{E}_1(q, p)$ given by (5.24) leads to, on the one hand:

$$\begin{aligned} \frac{d}{dt} \frac{\partial \mathcal{E}_1}{\partial p} \cdot p^* &= \frac{d}{dt} (\underline{M}_a(q) \mathcal{R}(\vartheta)^t p) \cdot \mathcal{R}(\vartheta)^t p^* - \Upsilon_3(q, p, p^*) \\ &= M_a(q) p' \cdot p^* - \Upsilon_2(q, p, p^*) + \Xi_3(q, p, p^*) - \Upsilon_3(q, p, p^*). \end{aligned} \quad (5.38)$$

On the other hand, we get:

$$\frac{\partial \mathcal{E}_1}{\partial q} \cdot p^* = -\Upsilon_1(q, p, p^*) + \frac{1}{2} \Xi_1(q, p, p^*). \quad (5.39)$$

Gathering (5.38) and (5.39) and (5.25) we deduce (5.34). \square

Proof of Lemma 5.8. On the one hand, invoking the symmetry of $M(q)$, we get:

$$\Upsilon_2(q, p, p^*) = \omega M_a(q) \begin{pmatrix} 0 \\ \ell^\perp \end{pmatrix} \cdot p^*. \quad (5.40)$$

On the other hand, recalling the notation (2.24), one has:

$$-\Upsilon_3(q, p, p^*) + \Upsilon_1(q, p, p^*) = \begin{pmatrix} -P_a^\perp \cdot \ell \\ \omega P_a^\perp \end{pmatrix} \cdot p^* = - \left[\begin{pmatrix} 0 \\ P_a \end{pmatrix} \times p \right] \cdot p^*. \quad (5.41)$$

Now gathering (5.40) and (5.41) concludes the proof of Lemma 5.8. \square

Proof of Lemma 5.9. We will use the following lemma.

Lemma 5.10. *For $i, j = 1, 2, 3$, for every $\hat{q} = (\hat{\vartheta}, \hat{h})$ in \mathcal{Q} and every $p^* = (\omega^*, \ell^*)$ in \mathbb{R}^3 , we have:*

$$\frac{\partial}{\partial q} \left(\int_{\underline{\mathcal{F}}(q)} \nabla \varphi_i(q) \cdot \nabla \varphi_j(q) dy \right) \Big|_{q=\hat{q}} \cdot p^* = - \int_{\partial \underline{\Omega}(\hat{q})} \frac{\partial \varphi_i}{\partial \tau} \frac{\partial \varphi_j}{\partial \tau} (\hat{w}^* \cdot n) ds, \quad (5.42)$$

with

$$\hat{w}^*(\hat{q}, p^*, \cdot) := -\omega^* \cdot^\perp - R(\hat{\vartheta})^t \ell^*. \quad (5.43)$$

Let us first take Lemma 5.10 for granted and conclude the proof of Lemma 5.9. Applying the change of variables $x = R(\hat{\vartheta})y + \hat{h}$, we deduce that:

$$\begin{aligned} \frac{\partial \underline{M}_a}{\partial q}(\hat{q}) \cdot p^* &= \left(\frac{\partial}{\partial q} \left(\int_{\underline{\mathcal{F}}(q)} \nabla \varphi_i(q) \cdot \nabla \varphi_j(q) dy \right) \Big|_{q=\hat{q}} \cdot p^* \right)_{1 \leq i, j \leq 3} \\ &= - \left(\int_{\partial \underline{\Omega}(\hat{q})} \frac{\partial \varphi_i}{\partial \tau} \frac{\partial \varphi_j}{\partial \tau} (\hat{w}^* \cdot n) ds \right)_{1 \leq i, j \leq 3} \\ &= \mathcal{R}(\hat{\vartheta})^t \left(\int_{\partial \Omega} \frac{\partial \varphi_i}{\partial \tau}(\hat{q}) \frac{\partial \varphi_j}{\partial \tau}(\hat{q}) (w^* \cdot n) ds \right)_{1 \leq i, j \leq 3} \mathcal{R}(\hat{\vartheta}), \end{aligned}$$

with $w^* := \omega^*(x - h)^\perp + \ell^*$. Therefore, applying this with $(\hat{q}, p^*) = (q, p^*)$ and with $(\hat{q}, p^*) = (q, p)$, we get

$$\Xi_3(q, p, p^*) - \frac{1}{2}\Xi_1(q, p, p^*) = \sum \left[A_{ij}^l(q) p_i p_j - \frac{1}{2} A_{jl}^i(q) p_i p_j \right] p_i^*,$$

where, as before, the notation \sum stands for $\sum_{1 \leq i, j, l \leq 3}$ for the rest of this proof and for every $k = 1, 2, 3$ the matrices $A^k(\hat{q})$ are given by:

$$A^k(\hat{q}) = \left(\int_{\partial\Omega} \frac{\partial \varphi_i}{\partial \tau}(\hat{q}) \frac{\partial \varphi_j}{\partial \tau}(\hat{q}) K_k(\hat{q}, \cdot) ds \right)_{1 \leq i, j \leq 3},$$

where we recall that $K_1(\hat{q}, \cdot) = (x - \hat{h})^\perp \cdot n$ and $K_j(\hat{q}, \cdot) = n_{j-1}$ ($j = 2, 3$) on $\partial\Omega$. The quadratic form in p can be symmetrized as follows:

$$\sum \left[A_{ij}^l(q) p_i p_j - \frac{1}{2} A_{jl}^i(q) p_i p_j \right] p_i^* = \frac{1}{2} \sum \left[A_{ij}^l + A_{il}^j - A_{jl}^i \right] (q) p_i p_j p_i^*,$$

which leads to (5.35). This concludes the proof of Lemma 5.9. \square

Now we give the proof of Lemma 5.10.

Proof of Lemma 5.10. The quantity

$$\frac{\partial}{\partial q} \left(\int_{\mathcal{F}(q)} \nabla \underline{\varphi}_i(q) \cdot \nabla \underline{\varphi}_j(q) dy \right) \Big|_{q=\hat{q}} \cdot p^* \quad (5.44)$$

can be interpreted as the time derivative of the quantity between parentheses, when the fluid outer boundary $\partial\Omega(\hat{q})$ undergoes a rigid displacement of velocity w^* .

More precisely, denote by χ a cut-off function, compactly supported, valued in $[0, 1]$ and such that $\chi = 1$ in a neighborhood of $\partial\Omega(\hat{q})$ and $\chi = 0$ in a neighborhood of \mathcal{S}_0 . Then, denote by $\xi(t, \cdot)$ the flow associated with the ODE:

$$\xi'(t, y) = \chi(\xi(t, y)) \hat{w}^*(t, \xi(t, y)), \quad \text{for } t > 0, \text{ with } \xi(0, y) = y, \quad (5.45)$$

where \hat{w}^* was introduced in (5.43). Notice that:

$$\xi(t, y) = R(-t\omega^*)y - tR(\hat{v})^t \ell^*,$$

in a neighborhood of $\partial\Omega(\hat{q})$ and $\xi(t, y) = y$ in a neighborhood of $\partial\mathcal{S}_0$.

For every t small, define

$$\Omega_t := \xi(t, \Omega(\hat{q})) \text{ and } \mathcal{F}_t := \xi(t, \mathcal{F}(\hat{q})).$$

For $j = 1, 2, 3$, let $\underline{\varphi}_{j,t}$ be harmonic in \mathcal{F}_t and satisfy the Neumann boundary conditions:

$$\frac{\partial \underline{\varphi}_{j,t}}{\partial n} = \begin{cases} (y - \hat{h})^\perp \cdot n & j = 1; \\ n_{j-1} & j = 2, 3 \end{cases} \quad \text{on } \partial\mathcal{S}_0, \quad (5.46a)$$

$$\frac{\partial \underline{\varphi}_{j,t}}{\partial n} = 0 \quad (j = 1, 2, 3) \quad \text{on } \partial\Omega_t. \quad (5.46b)$$

With these settings, the quantity (5.44) can be rewritten as:

$$\frac{d}{dt} \left(\int_{\mathcal{F}_t} \nabla \underline{\varphi}_{i,t} \cdot \nabla \underline{\varphi}_{j,t} \, dx \right) \Big|_{t=0}.$$

According to the Reynolds transport theorem, it can be expanded as follows:

$$\begin{aligned} \frac{d}{dt} \left(\int_{\mathcal{F}_t} \nabla \underline{\varphi}_{i,t} \cdot \nabla \underline{\varphi}_{j,t} \, dx \right) \Big|_{t=0} &= \int_{\mathcal{F}(\hat{q})} \nabla \underline{\varphi}'_i \cdot \nabla \underline{\varphi}_j \, dx + \int_{\mathcal{F}(\hat{q})} \nabla \underline{\varphi}_i \cdot \nabla \underline{\varphi}'_j \, dx \\ &\quad + \int_{\partial \underline{\Omega}(\hat{q})} (\nabla \underline{\varphi}_i \cdot \nabla \underline{\varphi}_j) (\hat{w}^* \cdot n) \, ds, \end{aligned} \quad (5.47)$$

where

$$\underline{\varphi}'_j := \frac{\partial \underline{\varphi}_{j,t}}{\partial t} \Big|_{t=0}.$$

Lemma 5.11. *For $j = 1, 2, 3$, the function $\underline{\varphi}'_j$ is harmonic in $\mathcal{F}(\hat{q})$, satisfies*

$$\frac{\partial \underline{\varphi}'_j}{\partial n} = 0 \text{ on } \partial \mathcal{S}_0 \quad (5.48)$$

and

$$\frac{\partial \underline{\varphi}'_j}{\partial n} = \frac{\partial}{\partial \tau} \left((\hat{w}^* \cdot n) \frac{\partial \underline{\varphi}_j}{\partial \tau} \right) \text{ on } \partial \underline{\Omega}(\hat{q}) = \partial \underline{\Omega}(\hat{q}). \quad (5.49)$$

Once Lemma 5.11 is proved, (5.42) follows from (5.47) and an integration by parts. \square

Proof of Lemma 5.11. The function $\underline{\varphi}'_j$ is defined and harmonic in $\mathcal{F}(\hat{q})$ and the boundary conditions are obtained by differentiating with respect to t , at $t = 0$, the identities on the fixed boundaries $\partial \mathcal{S}_0$ and $\partial \underline{\Omega}(\hat{q})$:

$$\frac{\partial \underline{\varphi}_{j,t}}{\partial n} (\xi(t, \cdot)) = \begin{cases} (y - \hat{h})^\perp \cdot n & j = 1; \\ n_{j-1} & j = 2, 3; \end{cases} \text{ on } \partial \mathcal{S}_0, \quad (5.50a)$$

$$\frac{\partial \underline{\varphi}_{j,t}}{\partial n} (\xi(t, \cdot)) = 0 \quad (j = 1, 2, 3) \text{ on } \partial \underline{\Omega}(\hat{q}). \quad (5.50b)$$

Let us focus on the proof of (5.49), the proof of (5.48) being quite similar with some simplifications. On $\partial \underline{\Omega}(\hat{q})$, using (5.45), we can write that:

$$\frac{d}{dt} \left(\frac{\partial \underline{\varphi}_{j,t}}{\partial n} (\xi(t, \cdot)) \right) \Big|_{t=0} = \frac{\partial \underline{\varphi}'_j}{\partial n} + \langle D^2 \underline{\varphi}_j, \hat{w}^*, n \rangle + \omega^* \frac{\partial \underline{\varphi}_j}{\partial \tau}, \quad (5.51)$$

where the last term is obtained by noticing that $n(\xi(t, \cdot)) = R(-t\omega^*)n$. Therefore by taking the derivative at $t = 0$ of the identity (5.50b) and using (5.51) we obtain

$$\begin{aligned} \frac{\partial \varphi'_j}{\partial n} &= -\langle D^2 \varphi_j, \hat{w}^*, n \rangle - \omega^* \frac{\partial \varphi_j}{\partial \tau} \\ &= -\frac{\partial^2 \varphi_j}{\partial n^2} (\hat{w}^* \cdot n) - \langle D^2 \varphi_j, \tau, n \rangle (\hat{w}^* \cdot \tau) - \omega^* \frac{\partial \varphi_j}{\partial \tau}, \end{aligned} \quad (5.52)$$

by decomposing w^* into normal and tangential parts. Taking the tangential derivative of the identity (5.36b), we get:

$$\langle D^2 \varphi_j, \tau, n \rangle + \kappa \frac{\partial \varphi_j}{\partial \tau} = 0 \quad \text{on } \partial \underline{\Omega}(\hat{q}), \quad (5.53)$$

where we used the relation $\frac{\partial n}{\partial \tau} = \kappa \tau$ with κ the local curvature of $\partial \underline{\Omega}(\hat{q})$. Plugging (5.53) into (5.52) yields the identity:

$$\begin{aligned} \frac{\partial \varphi'_j}{\partial n} &= -\frac{\partial^2 \varphi_j}{\partial n^2} (\hat{w}^* \cdot n) + \left(\kappa (\hat{w}^* \cdot \tau) - \omega^* \right) \frac{\partial \varphi_j}{\partial \tau} \\ &= -\frac{\partial^2 \varphi_j}{\partial n^2} (\hat{w}^* \cdot n) + \left(\frac{\partial}{\partial \tau} (\hat{w}^* \cdot n) \right) \frac{\partial \varphi_j}{\partial \tau}. \end{aligned} \quad (5.54)$$

On $\partial \underline{\Omega}(\hat{q})$, we have with local coordinates:

$$\Delta \varphi_j = \frac{\partial^2 \varphi_j}{\partial \tau^2} - \kappa \frac{\partial \varphi_j}{\partial n} + \frac{\partial^2 \varphi_j}{\partial n^2}.$$

Since φ_j is harmonic and $\frac{\partial \varphi_j}{\partial n} = 0$ on $\partial \underline{\Omega}(\hat{q})$, we deduce that $\frac{\partial^2 \varphi_j}{\partial n^2} = -\frac{\partial^2 \varphi_j}{\partial \tau^2}$ on $\partial \underline{\Omega}(\hat{q})$, and therefore

$$\frac{\partial \varphi'_j}{\partial n} = (\hat{w}^* \cdot n) \frac{\partial^2 \varphi_j}{\partial \tau^2} + \left(\frac{\partial}{\partial \tau} (\hat{w}^* \cdot n) \right) \frac{\partial \varphi_j}{\partial \tau},$$

which is (5.49). □

5.5 Expression of the potential energy: Proof of Lemma 2.4

By definition we have

$$C(q) = - \int_{\mathcal{F}(q)} \nabla \psi(q, \cdot) \cdot \nabla \psi(q, \cdot) \, dx.$$

Thus, by the Reynolds transport theorem we infer that for every p in \mathbb{R}^3 ,

$$DC(q) \cdot p = -2 \int_{\mathcal{F}(q)} \nabla \left(\frac{\partial \psi}{\partial q} \cdot p \right) \cdot \nabla \psi \, dx - \int_{\partial \mathcal{S}(q)} |\nabla \psi|^2 u_{q'} \cdot n \, ds, \quad (5.55)$$

where $u_{q'}$ was defined in (2.14). Upon an integration by parts, we get

$$\int_{\mathcal{F}(q)} \nabla \left(\frac{\partial \psi}{\partial q} \cdot p \right) \cdot \nabla \psi \, dx = \int_{\partial \mathcal{S}(q)} \left(\frac{\partial \psi}{\partial q} \cdot p \right) \frac{\partial \psi}{\partial n} \, ds. \quad (5.56)$$

Moreover gathering (5.29) and (5.30) yields

$$\frac{\partial \psi}{\partial q}(q, x) \cdot p = DC(q) \cdot p - \frac{\partial \psi}{\partial n}(q, x) \frac{\partial \varphi}{\partial n}(q, x) \cdot p. \quad \text{for } x \in \partial \mathcal{S}(q). \quad (5.57)$$

Now combining (5.56) and (5.57) we obtain

$$\begin{aligned} \int_{\mathcal{F}(q)} \nabla \left(\frac{\partial \psi}{\partial q} \cdot p \right) \cdot \nabla \psi \, dx &= \int_{\partial \mathcal{S}(q)} (DC(q) \cdot p) \frac{\partial \psi}{\partial n} \, ds - \int_{\partial \mathcal{S}(q)} \left| \frac{\partial \psi}{\partial n} \right|^2 \frac{\partial \varphi}{\partial n} \cdot p \, ds \\ &= -DC(q) \cdot p - \int_{\partial \mathcal{S}(q)} \left| \frac{\partial \psi}{\partial n} \right|^2 \frac{\partial \varphi}{\partial n} \cdot p \, ds, \end{aligned} \quad (5.58)$$

thanks to (2.10d). On the other hand since $\psi(q, \cdot)$ is constant on $\partial \mathcal{S}(q)$, we get

$$\int_{\partial \mathcal{S}(q)} |\nabla \psi|^2 u_{q'} \cdot n \, ds = \int_{\partial \mathcal{S}(q)} \left| \frac{\partial \psi}{\partial n} \right|^2 \frac{\partial \varphi}{\partial n} \cdot p \, ds. \quad (5.59)$$

Gathering (5.55), (5.58), (5.59) and (2.17b) leads to the result. This concludes the proof of Lemma 2.4.

5.6 Conservation of energy: Proof of Proposition 2.5

We start with the observation that

$$(\mathcal{E}(q, q'))' = M(q)q'' \cdot q' + \frac{1}{2}(DM(q) \cdot q')q' \cdot q' - \frac{1}{2}\gamma^2 DC(q) \cdot q'. \quad (5.60)$$

Now, thanks to (2.20) and (2.17c), we have

$$M(q)q'' \cdot q' = -\langle \Gamma(q), q', q' \rangle \cdot q' + F(q, q') \cdot q', \quad (5.61)$$

and

$$F(q, q') \cdot q' = \gamma^2 E(q) \cdot q'. \quad (5.62)$$

We introduce the matrix for any (q, p) in $\mathcal{Q} \times \mathbb{R}^3$,

$$S(q, p) := \left(\sum_{1 \leq i \leq 3} \Gamma_{i,j}^k(q) p_i \right)_{1 \leq k, j \leq 3}, \quad (5.63)$$

so that

$$\langle \Gamma(q), p, p \rangle = S(q, p)p. \quad (5.64)$$

Combining (5.60), (5.61), (5.62), (5.63) and (5.64) we obtain

$$(\mathcal{E}(q, q'))' = \gamma^2 \left(E(q) - \frac{1}{2} DC(q) \right) \cdot q' + \left(\frac{1}{2} DM(q) \cdot q' - S(q, q') \right) q' \cdot q'. \quad (5.65)$$

The first term of the right hand side vanishes thanks to (2.22). The proof of Proposition 2.5 follows then from the following result.

Lemma 5.12. *For any (q, p) in $\mathcal{Q} \times \mathbb{R}^3$, the matrix $\frac{1}{2} DM(q) \cdot p - S(q, p)$ is skew-symmetric.*

Proof We start with the observation that $DM(q) \cdot p$ is the 3×3 matrix containing the entries

$$\sum_{1 \leq k \leq 3} (M_a)_{i,j}^k(q) p_k, \text{ for } 1 \leq i, j \leq 3,$$

where $(M_a)_{i,j}^k(q)$ is defined in (2.16c). On the other hand, the matrix $S(q, p)$ contains the entries

$$\frac{1}{2} \sum_{1 \leq k \leq 3} \left((M_a)_{i,j}^k + (M_a)_{i,k}^j - (M_a)_{k,j}^i \right) (q) p_k,$$

for $1 \leq i, j \leq 3$. Therefore, the matrix $DM(q) \cdot p - S(q, p)$ contains the entries

$$c_{ij}(q, p) = -\frac{1}{2} \sum_{1 \leq k \leq 3} \left((M_a)_{i,k}^j - (M_a)_{k,j}^i \right) (q) p_k,$$

for $1 \leq i, j \leq 3$. Using that the matrix $M(q)$ is symmetric, we get that $c_{ij}(q, p) = -c_{ji}(q, p)$ for $1 \leq i, j \leq 3$. \square

5.7 The case of a disk: proof of Theorem 2.8

In this subsection, we suppose that \mathcal{S}_0 is a disk, say of center ζ and radius r_0 . We start by observing that, as noticed in Subsection 4.1, $\zeta = h_c(0) = h_c(0) - h(0)$, so with (4.14) we deduce $\zeta_\vartheta = h_c - h$, that is, (2.34b). Relation (2.34c) is an immediate consequence of the decomposition $x - h = (x - h_c) + (h_c - h)$, the fact that $(x - h_c) \cdot n = 0$ on $\partial \mathcal{S}_0$, (2.1d) and (2.34b). Hence only (2.34a) needs proving.

Concerning the added mass matrix, due to (2.8), one has in this case $\partial_n \varphi_1(q, \cdot) = (h_c - h)^\perp \cdot n(\cdot)$, and consequently,

$$\varphi_1(q, \cdot) = -\zeta_{\vartheta,2} \varphi_2(q, \cdot) + \zeta_{\vartheta,1} \varphi_3(q, \cdot). \quad (5.66)$$

We underline that φ_2 and φ_3 depend merely on h_c while φ_1 depends also on ϑ . From (5.66), one deduces that for any q in \mathcal{Q} such that $d(B(h, r_0), \partial \Omega) > 0$ and any $p = (\omega, \ell)$:

$$M_a(q)p \cdot p = \tilde{M}_b(h_c)p_b \cdot p_b \quad (5.67)$$

with

$$p_b = p_b(\vartheta, \omega, \ell) := \ell + \omega \zeta_\vartheta^\perp. \quad (5.68)$$

Notice that for a solution $(p, q) : [0, T] \rightarrow \mathcal{Q} \times \mathbb{R}^3$ of the system with $p = q'$ one has for all times

$$p_b(\vartheta(t), \vartheta'(t), h'(t)) = h'_c(t). \quad (5.69)$$

Now to establish (2.34a) we rely on the following adaptation of Lemma 5.5.

Lemma 5.13. *For any smooth curve $t \mapsto q(t)$ in \mathcal{Q} , for every $p^* = (\omega^*, \ell^*) \in \mathbb{R}^3$, we have:*

$$\mathcal{E}\mathcal{L} = \tilde{M}_b(h_c)p'_b \cdot p_b^* + \langle \Gamma_b(h_c), p_b, p_b \rangle \cdot p_b^*, \quad (5.70)$$

where

$$p_b^* = p_b^*(\vartheta, \omega^*, \ell^*) := \ell^* + \omega^* \zeta_\vartheta^\perp, \quad (5.71)$$

with $\mathcal{E}\mathcal{L}$ is given by (5.11), $\tilde{M}_b(h_c)$ is defined by (2.28)-(2.30) and $\Gamma_b(h_c)$ is defined in (2.31).

Proof Mimicking the proof of Lemma 5.5, using $\frac{d}{dt}p_b^* = -\vartheta'\omega^*\zeta_\vartheta$, (2.30) and (5.67) we obtain

$$\frac{d}{dt} \frac{\partial \mathcal{E}_1}{\partial p} \cdot p^* = \tilde{M}_b p'_b \cdot p_b^* + (D_{h_c} \tilde{M}_b \cdot h'_c) p_b \cdot p_b^* - \tilde{M}_b p_b \cdot (\vartheta' \omega^* \zeta_\vartheta).$$

On the other side, one has

$$\frac{\partial \mathcal{E}_1}{\partial q} \cdot p^* = \frac{1}{2} (D_{h_c} \tilde{M}_b \cdot \omega^* \zeta_\vartheta^\perp) p_b \cdot p_b + \tilde{M}_b (\vartheta' \omega^* \zeta_\vartheta) \cdot p_b.$$

One gets

$$\mathcal{E}\mathcal{L} = \tilde{M}_b p'_b \cdot p_b^* + (D_{h_c} \tilde{M}_b \cdot h'_c) p_b \cdot p_b^* - \frac{1}{2} (D_{h_c} \tilde{M}_b \cdot \omega^* \zeta_\vartheta^\perp) p_b \cdot p_b,$$

and one concludes as in Lemma 5.5, by symmetrizing the second term. \square

Relying on Lemmas 5.1 and 5.2, Theorem 2.8 finally follows from a rewriting of $E(q)$ and $B(q)$ taking (5.66) into account. Details are left to the reader. This ends the proof of Theorem 2.8.

6 Convergence to the massive point vortex system in Case (i): Proof of Theorem 2.9

In this section we prove Theorem 2.9 which corresponds to Case (i), relying on results (Proposition 6.3, Lemma 6.4 and Proposition 9.1) whose proof is postponed to the last sections. We work on Equations (2.20) with a shrunk solid, that is

$$M_\varepsilon(q_\varepsilon)q_\varepsilon'' + \langle \Gamma_\varepsilon(q_\varepsilon), q'_\varepsilon, q'_\varepsilon \rangle = F_\varepsilon(q_\varepsilon, q'_\varepsilon), \quad (6.1)$$

where

$$M_\varepsilon := M[\mathcal{S}_{0,\varepsilon}, m_\varepsilon, \mathcal{J}_\varepsilon, \Omega], \quad \Gamma_\varepsilon := \Gamma[\mathcal{S}_{0,\varepsilon}, \Omega] \quad \text{and} \quad F_\varepsilon := F[\mathcal{S}_{0,\varepsilon}, \gamma, \Omega],$$

with $m_\varepsilon, \mathcal{J}_\varepsilon$ given by (2.37). The functions $M_\varepsilon(q), \langle \Gamma_\varepsilon(q), p, p \rangle$ and $F_\varepsilon(q, p)$ are defined for q in \mathcal{Q}^ε and for p in \mathbb{R}^3 .

We begin by introducing some notations. Given $\delta > 0$ and ε_0 in $(0, 1)$, we let

$$\mathfrak{Q} := \{(\varepsilon, q) \in (0, 1) \times \mathbb{R}^3 : d(\mathcal{S}_\varepsilon(q), \partial\Omega) > 0\}, \quad (6.2)$$

$$\mathfrak{Q}^\delta := \{(\varepsilon, q) \in (0, 1) \times \mathbb{R}^3 : d(\mathcal{S}_\varepsilon(q), \partial\Omega) > \delta\}, \quad (6.3)$$

$$\mathfrak{Q}_{\delta, \varepsilon_0} := \{(\varepsilon, q) \in (0, \varepsilon_0) \times \mathbb{R}^3 : d(\mathcal{S}_\varepsilon(q), \partial\Omega) > \delta\}. \quad (6.4)$$

We will also make use, for $\delta > 0$, of

$$\Omega_\delta := \{x \in \Omega : d(x, \partial\Omega) > \delta\}. \quad (6.5)$$

Observe that despite the fact that the center of mass h_ε does not necessarily belong to $\mathcal{S}_\varepsilon(q)$, we have the following elementary result whose proof is left to the reader.

Lemma 6.1. *Let $\delta > 0$. There exists δ_0 in $(0, \delta)$ and ε_0 in $(0, 1]$ such that for any (ε, q) in $\mathfrak{Q}_{\delta, \varepsilon_0}$, with $q = (\vartheta, h)$, necessarily h belongs to Ω_{δ_0} .*

6.1 Normal form

In the Case (i) considered in this section, we will rephrase the normal form (6.1) suitably, in order to pass to the limit as ε goes to 0. The following definition will be useful to deal with the remainder.

Definition 6.2. Let $\delta > 0$ and $\varepsilon_0 \in (0, 1)$ be given. We say that a vector field F in $L^\infty(\mathfrak{Q}_{\delta, \varepsilon_0} \times \mathbb{R}^3; \mathbb{R}^3)$ is weakly nonlinear if there exists $K > 0$ depending on $\mathcal{S}_0, m, \mathcal{J}, \gamma, \Omega$ and δ such that for any (ε, q, p) in $\mathfrak{Q}_{\delta, \varepsilon_0} \times \mathbb{R}^3$,

$$|F(\varepsilon, q, p)|_{\mathbb{R}^3} \leq K(1 + |p|_{\mathbb{R}^3} + \varepsilon|p|_{\mathbb{R}^3}^2). \quad (6.6)$$

The normal form is as follows.

Proposition 6.3. *Let $\delta > 0$. There exists ε_0 in $(0, 1)$ and F_r in $L^\infty(\mathfrak{Q}_{\delta, \varepsilon_0} \times \mathbb{R}^3; \mathbb{R}^3)$ depending on \mathcal{S}_0, γ and Ω , weakly nonlinear in the sense of Definition 6.2 such that Equation (6.1) can be recast as*

$$M_q p'_\varepsilon = F_{\vartheta'_\varepsilon}^{\mathcal{S}_0}(\varepsilon \vartheta'_\varepsilon, h'_\varepsilon - \gamma u^\Omega(h_\varepsilon)) + \varepsilon F_r(\varepsilon, q_\varepsilon, p_\varepsilon). \quad (6.7)$$

Recall that u^Ω was defined in (2.42), that the force term $F_{\vartheta'_\varepsilon}^{\mathcal{S}_0}(p)$ was defined in the case without outer boundary, see (4.13), and that $p_\varepsilon = (\varepsilon \vartheta'_\varepsilon, h'_\varepsilon)^t$ was defined in (4.28).

The normal form (6.7) will be useful in order to pass to the limit. To get Proposition 6.3, we will perform expansions of the inertia matrix, of the Christoffel symbols and of the force terms with respect to ε . Roughly speaking the leading terms coming from the force terms will be gathered into the first term of the right hand side of (6.7), see (9.60).

6.2 Renormalized energy estimates

We will of course need uniform estimates as $\varepsilon \rightarrow 0^+$ in order to pass to the limit in (6.7). The energy is the natural candidate to yield such estimates. Hence we are led to consider the behavior of the energy with respect to ε . We index the energy as follows:

$$\mathcal{E}_\varepsilon(q, p) := \frac{1}{2} M_\varepsilon(q) p \cdot p + U_\varepsilon(q), \quad (6.8)$$

where the potential energy U_ε is given by

$$U_\varepsilon(q) := -\frac{1}{2} \gamma^2 C_\varepsilon(q). \quad (6.9)$$

Of course Proposition 2.5 can be applied for each $\varepsilon \in (0, 1)$ so that the energy associated with a solution q_ε as in Theorem 2.9 is conserved along time until its maximal time of existence T_ε . We will establish in Section 8.3 the following result regarding the expansion of $C_\varepsilon(q)$ with respect to ε . The expansion is uniform, in the sense that the remainder is uniformly bounded, as long as the solid stays at a positive distance from the external boundary. Let us recall that the Newtonian potential G was introduced in (2.40), the Kirchhoff-Routh stream function ψ^Ω was defined in (2.41) and the constant C^{S_0} in (4.4). We will also use the scalar field defined for $q := (\vartheta, h)$ in $\mathbb{R} \times \Omega$,

$$\psi_c(q) := D_h \psi^\Omega(h) \cdot \zeta_\vartheta. \quad (6.10)$$

Above D_h denotes the derivative with respect to h .

Lemma 6.4. *Let $\delta > 0$. There exists ε_0 in $(0, 1)$ and a function C_r in the space $L^\infty(\mathfrak{Q}_{\delta, \varepsilon_0}; \mathbb{R})$ such that for any (ε, q) in $\mathfrak{Q}_{\delta, \varepsilon_0}$,*

$$C_\varepsilon(q) = -G(\varepsilon) + C^{S_0} + 2\psi^\Omega(h) + 2\varepsilon\psi_c(q) + \varepsilon^2 C_r(\varepsilon, q). \quad (6.11)$$

We recall that C^{S_0} was introduced in (4.4). This result establishes that the potential energy $U_\varepsilon(q)$ diverges logarithmically as $\varepsilon \rightarrow 0^+$. However the contributions of the two first terms of (6.11) in (6.8) can be discarded from the energy (6.8) since they do not depend on the solid position and velocity.

On the other hand an immediate consequence of Proposition 9.1 below (on the asymptotic development of the added mass matrix with respect to ε) is the following result regarding the kinetic energy part.

Lemma 6.5. *Let $\delta > 0$. There exists ε_0 in $(0, 1)$ and a function M_r in the space $L^\infty(\mathfrak{Q}_{\delta, \varepsilon_0}; \mathbb{R}^{3 \times 3})$ depending on S_0 and Ω , such that, for all (ε, q) in $\mathfrak{Q}_{\delta, \varepsilon_0}$, for all p in \mathbb{R}^3 ,*

$$\frac{1}{2} M_\varepsilon(q) p \cdot p = \frac{1}{2} ((M_g + \varepsilon^2 M_{a, \vartheta}^{S_0}) I_\varepsilon p) \cdot (I_\varepsilon p) + \frac{1}{2} \varepsilon^4 (M_r(\varepsilon, q) I_\varepsilon p) \cdot (I_\varepsilon p).$$

Recall that $M_{a, \vartheta}^{S_0}$ was defined in (4.9) and I_ε in (4.27). Combining Lemma 6.4 and Lemma 6.5 we obtain the following.

Corollary 6.6. *Let q_ε and T_ε be as in Theorem 2.9. Then the renormalized energy*

$$\begin{aligned} & \frac{1}{2}((M_g + \varepsilon^2 M_{a,\vartheta}^{\mathcal{S}_0})p_\varepsilon) \cdot p_\varepsilon - \gamma^2 \psi^\Omega(h_\varepsilon) \\ & \quad + \frac{1}{2}\varepsilon^4 M_r(\varepsilon, q_\varepsilon)p_\varepsilon \cdot p_\varepsilon + \frac{1}{2}\varepsilon\gamma^2(\psi_c(q_\varepsilon) + \varepsilon C_r(\varepsilon, q_\varepsilon)), \end{aligned}$$

is constant in time until T_ε .

The two most important terms in the renormalized energy above are the first and second ones which are respectively of order $O(|p_\varepsilon|_{\mathbb{R}^3}^2)$ and $O(1)$ as long as there is no collision. Hence the uniformity in ε in Lemma 6.4 and Lemma 6.5, the conservation property stated in Corollary 6.6 and the uniform coercivity of M_g allow to get the following counterpart of Corollary 2.6.

Corollary 6.7. *Let $(q_\varepsilon, p_\varepsilon)$ satisfy the assumptions of Theorem 2.9. Let $\delta > 0$. There exists $K > 0$ (depending on \mathcal{S}_0 , Ω , p_0 , γ , m_1 , \mathcal{J}_1 , δ) and $\varepsilon_0 > 0$ such that for any ε in $(0, \varepsilon_0)$, for any $t > 0$ such that as long as $(\varepsilon, q_\varepsilon(t))$ belongs to $\mathfrak{Q}_{\delta, \varepsilon_0}$, one has $|p_\varepsilon(t)|_{\mathbb{R}^3} \leq K$.*

6.3 Passage to the limit

We deduce from Corollary 6.7 two different results. The first result concerns the lifetime T_ε of the solution $(q_\varepsilon, p_\varepsilon)$, which can be only limited by a possible encounter between the solid and the boundary $\partial\Omega$.

Lemma 6.8. *There exist $\varepsilon_0 > 0$, $\underline{T} > 0$ and $\underline{\delta} > 0$, such that for any ε in $(0, \varepsilon_0)$, we have*

$$T^\varepsilon \geq \underline{T} \text{ and moreover on } [0, \underline{T}], \text{ one has } (\varepsilon, q_\varepsilon) \in \mathfrak{Q}_{\underline{\delta}, \varepsilon_0}. \quad (6.12)$$

Proof Let us introduce

$$R_0 := \max\{|x|, x \in \partial\mathcal{S}_0\}, \quad (6.13)$$

so that, whatever $t \geq 0$, $\vartheta(t)$ in \mathbb{R} and ε in $(0, 1)$, one has

$$\mathcal{S}_\varepsilon(q_\varepsilon(t)) \subset \overline{B}(h_\varepsilon(t), \varepsilon R_0). \quad (6.14)$$

We introduce $\underline{\delta} := \frac{1}{4}d(0, \partial\Omega)$ and ε_0 in $(0, 1)$ (which may be reduced later) such that $\varepsilon_0 R_0 \leq \underline{\delta}$ and

$$\forall \varepsilon \in (0, \varepsilon_0], \quad d(\overline{B}(0, \varepsilon R_0), \partial\Omega) \geq \frac{3}{4}d(0, \partial\Omega). \quad (6.15)$$

We apply Corollary 6.7 with $\underline{\delta}$ to deduce that there exists $K > 0$ such that, reducing ε_0 if necessary, we have

$$|h'_\varepsilon| \leq K, \text{ for all } t \text{ for which } d(\mathcal{S}_\varepsilon(q_\varepsilon(t)), \partial\Omega) \geq \underline{\delta}. \quad (6.16)$$

We introduce $\underline{T} := \min\left(1, \frac{d(0, \partial\Omega)}{2K}\right)$, and for ε in $(0, \varepsilon_0]$,

$$\mathcal{I}^\varepsilon = \{t \in [0, 1] \mid \forall s \in [0, t], d(\overline{B}(h_\varepsilon(s), \varepsilon R_0), \partial\Omega) \geq \underline{\delta}\}.$$

The set \mathcal{I}^ε is a closed interval containing 0, according to (6.15). Consider $\tilde{T}_\varepsilon := \max \mathcal{I}^\varepsilon$, and let us show that $\tilde{T}_\varepsilon \geq \underline{T}$. Of course, if $\tilde{T}_\varepsilon = 1$, then this is clear; let us suppose that $\tilde{T}_\varepsilon < 1$. This involves that $d(\overline{B}(h_\varepsilon(\tilde{T}_\varepsilon), \varepsilon R_0), \partial\Omega) = \underline{\delta}$. Using $\varepsilon_0 R_0 \leq \underline{\delta}$ we deduce $d(h_\varepsilon(\tilde{T}_\varepsilon), \partial\Omega) \leq 2\underline{\delta}$. With the triangle inequality and $\underline{\delta} = \frac{1}{4}d(0, \partial\Omega)$ we infer that $d(h_\varepsilon(\tilde{T}_\varepsilon), 0) \geq \frac{1}{2}d(0, \partial\Omega)$. Now the relation (6.14) implies that for all t in $[0, \tilde{T}_\varepsilon]$,

$$d(\mathcal{S}_\varepsilon(q_\varepsilon(t)), \partial\Omega) \geq d(\overline{B}(h_\varepsilon(t), \varepsilon R_0), \partial\Omega) \geq \underline{\delta}, \quad (6.17)$$

so that (6.16) is satisfied during $[0, \tilde{T}_\varepsilon]$. We deduce that $K\tilde{T}_\varepsilon \geq d(0, \partial\Omega)/2$, so $\tilde{T}_\varepsilon \geq \underline{T}$. Therefore for any t in $[0, \underline{T}]$, for any ε in $[0, \varepsilon_0]$, (6.17) holds true. This concludes the proof of Lemma 6.8. \square

The second result establishes the desired convergence on any time interval during which we have a minimal distance between $\mathcal{S}_\varepsilon(q)$ and $\partial\Omega$, uniform for small ε . Let us recall that (\bar{h}, \bar{T}) denotes the maximal solution of (2.38).

Lemma 6.9. *Let $\varepsilon_1 > 0$, $\check{\delta} > 0$ and $\check{T} > 0$ with $\check{T} < \bar{T}$, and suppose that for any ε in $(0, \varepsilon_1)$, we have*

$$d(\mathcal{S}_\varepsilon(q_\varepsilon(t)), \partial\Omega) \geq \check{\delta} \text{ on } [0, \check{T}]. \quad (6.18)$$

Then $(h_\varepsilon, \varepsilon\vartheta_\varepsilon) \rightharpoonup (\bar{h}, 0)$ in $W^{2,\infty}([0, \check{T}]; \mathbb{R}^3)$ weak- \star .

The proof of Lemma 6.9 consists in passing to the weak limit, with the help of all a priori bounds, in each term of (6.7). It is a straightforward extension of the proof of Theorem 4.7 and it is therefore omitted.

We now finish the proof of Theorem 2.9. It only remains to extend the time interval on which the above convergences are valid to any closed subinterval of $[0, \bar{T})$. Hence let $T \in (0, \bar{T})$, and let us prove that for small $\varepsilon > 0$ the time of existence T^ε is larger than T and establish the convergences on the time interval $[0, T]$. For such a T , we know that there exists $\bar{d} > 0$ such that

$$\forall t \in \left[0, \frac{T + \bar{T}}{2}\right], \quad d(\bar{h}(t), \partial\Omega) \geq \bar{d}. \quad (6.19)$$

We let $\bar{T}_\varepsilon := \max\{t > 0, d(B(h_\varepsilon(t), \varepsilon R_0), \partial\Omega) \geq \bar{d}/2\}$. Let us recall that R_0 is defined in (6.13). Using Lemma 6.8 we deduce that, reducing \bar{d} if necessary, we have that for some $\bar{\varepsilon} > 0$, $\inf_{\varepsilon \in (0, \bar{\varepsilon}]} \bar{T}_\varepsilon > 0$. Therefore $\tilde{T} := \liminf_{\varepsilon \rightarrow 0^+} \bar{T}_\varepsilon > 0$. Due to Corollary 6.7, there exists $K > 0$ and ε_0 such that for all t in $[0, T + 1]$ and ε in $(0, \varepsilon_0)$, one has the following estimate

$$|h'_\varepsilon| + |\varepsilon\vartheta'_\varepsilon| \leq K \text{ as long as } d(\mathcal{S}_\varepsilon(q_\varepsilon(t)), \partial\Omega) \geq \bar{d}/2. \quad (6.20)$$

Now we claim that

$$\tilde{T} \geq T + \frac{\bar{T} - T}{4}. \quad (6.21)$$

Suppose that this is not the case, so that we have a sequence $\varepsilon_n \rightarrow 0^+$ such that $\bar{T}_{\varepsilon_n} \rightarrow \tilde{T} < T + \frac{\bar{T} - T}{4}$. Now for any η in $(0, \tilde{T})$, on the interval $[0, \tilde{T} - \eta]$, the condition $d(\bar{B}(h_{\varepsilon_n}(t), \varepsilon_n R_0), \partial\Omega) \geq \bar{d}/2$ is satisfied for n large enough. Moreover, for such n , for all t in $[0, \tilde{T} - \eta]$, (6.14) implies that

$$d(\mathcal{S}_{\varepsilon_n}(q_{\varepsilon_n}(t)), \partial\Omega) \geq d(\bar{B}(h_{\varepsilon_n}(t), \varepsilon_n R_0), \partial\Omega) \geq \bar{d}/2.$$

Hence applying Lemma 6.9, we deduce the uniform convergence of $(h_{\varepsilon_n})_n$ to \bar{h} on $[0, \tilde{T} - \eta]$. In particular, as $n \rightarrow +\infty$,

$$d(h_{\varepsilon_n}(\tilde{T} - \eta), \partial\Omega) \rightarrow d(\bar{h}(\tilde{T} - \eta), \partial\Omega) \geq \bar{d},$$

according to (6.19). On the other hand by definition of \bar{T}_{ε_n} we have

$$d(B(h_{\varepsilon_n}(\bar{T}_{\varepsilon_n}), \varepsilon_n R_0), \partial\Omega) = \bar{d}/2.$$

Using the triangle inequality and $\bar{T}_{\varepsilon_n} \rightarrow \tilde{T}$, we get a contradiction with (6.20) for η small enough. Hence (6.21) is valid, so that, reducing $\bar{\varepsilon}$ if necessary, we have $\inf_{\varepsilon \in (0, \bar{\varepsilon})} \bar{T}_{\varepsilon} \geq T$. Now, applying again (6.14) and Lemma 6.9, we reach the conclusion. This ends the proof of Theorem 2.9.

7 Convergence to the point vortex system in Case (ii): Proof of Theorem 2.10

In this section we prove Theorem 2.10 which corresponds to Case (ii). Here again, we rely on results that are proved in subsequent sections. We work again on Equations (6.1) but here m_ε , and \mathcal{J}_ε are given by (2.43). We begin this section with a simplified case, that is the case where \mathcal{S}_0 is a homogeneous disk. Then we treat the case when \mathcal{S}_0 is not a disk, and finish with the case of a non-homogeneous disk.

7.1 The case of a homogeneous disk

When \mathcal{S}_0 is a homogeneous disk, we obtain a simplified normal form, which in particular depends only on the center of mass h and where the dynamics of the angle ϑ is trivial. Without loss of generality, we write $\mathcal{S}_0 = \bar{B}(0, 1)$.

To state this normal form, we first modify a bit Definition 6.2 as follows.

Definition 7.1. Let $\delta > 0$ and $\varepsilon_0 \in (0, 1)$ be given. We say that a vector field F_b in $L^\infty((0, \varepsilon_0) \times \Omega \times \mathbb{R}^2; \mathbb{R}^2)$ is weakly nonlinear if there exists $K > 0$ depending on $m, \mathcal{J}, \gamma, \Omega$ and δ such that for any ε in $(0, \varepsilon_0)$, any h in Ω such that $d(B(h, \varepsilon), \partial\Omega) > \delta$ and any $\ell \in \mathbb{R}^2$, one has

$$|F_b(\varepsilon, h, \ell)|_{\mathbb{R}^2} \leq K(1 + |\ell|_{\mathbb{R}^2} + \varepsilon|\ell|_{\mathbb{R}^2}^2). \quad (7.1)$$

We used the notation (2.5). Now the normal form is as follows.

Proposition 7.2. *There exists $F_{b,r}$ in $L^\infty((0,1) \times \Omega \times \mathbb{R}^2; \mathbb{R}^2)$, weakly non-linear in the sense of Definition 7.1, such that Equation (2.34) can be recast as*

$$\vartheta_\varepsilon'' = 0,$$

and

$$\begin{aligned} \left(\varepsilon^\alpha m_1 + \varepsilon^2 \pi \right) (h'_\varepsilon - u^\Omega(h_\varepsilon))' &= \gamma (h'_\varepsilon - u^\Omega(h_\varepsilon))^\perp \\ &+ \varepsilon^{\min(2,\alpha)} F_{b,r}(\varepsilon, h_\varepsilon, h'_\varepsilon). \end{aligned} \quad (7.2)$$

The coefficient π comes from the fact that $M_b^{S_0} = \pi \text{Id}_2$ in this case (where $M_b^{S_0}$ is defined in (4.23)). As a consequence, Corollary 6.7 has the following counterpart.

Corollary 7.3. *Let h_ε satisfy the assumptions of Theorem 2.8. Let $\delta > 0$. There exists $K > 0$ (depending on Ω , ℓ_0 , γ , m_1 , δ) and $\varepsilon_0 > 0$ such that for any ε in $(0, \varepsilon_0)$, for any $t > 0$, as long as $d(B(h_\varepsilon, \varepsilon), \partial\Omega) > \delta$ one has $\varepsilon^{\min(1, \frac{\alpha}{2})} |h'_\varepsilon|_{\mathbb{R}^2} \leq K$.*

Proof The proof is almost the same as for Corollary 6.7, but one has to take into account that the added mass matrix is degenerate and that dynamics of the rotation angle is trivial.

In the case of a homogeneous disk, the decomposition of the energy given in (6.8) can be described by using the following function, where $q = (\vartheta, h)$ and $p = (\omega, \ell)$:

$$\tilde{\mathcal{E}}_\varepsilon(h, \omega, \ell) = \frac{1}{2} \left(\varepsilon^\alpha m_1 \text{Id}_2 + \varepsilon^2 \tilde{M}_b(h) \right) \ell \cdot \ell + \varepsilon^{\alpha+2} \mathcal{J}_1 \omega^2 + \tilde{U}_\varepsilon(h). \quad (7.3)$$

Here we wrote $\tilde{U}_\varepsilon(h) := U_\varepsilon(q)$ for $q = (\vartheta, h)$ since actually it does not depend on the angle ϑ .

Now using the same analysis as for Corollary 6.6 and taking into account that $\vartheta'_\varepsilon(t) = \omega_0$ for all times, we deduce that the following quantity is conserved over time:

$$\begin{aligned} \frac{1}{2} \left(\varepsilon^\alpha m_1 \text{Id}_2 + \varepsilon^2 \tilde{M}_b(h) \right) h'_\varepsilon \cdot h'_\varepsilon - \gamma^2 \psi^\Omega(h_\varepsilon) \\ + \frac{1}{2} \varepsilon^4 \tilde{M}_{b,r}(\varepsilon, h_\varepsilon) h'_\varepsilon \cdot h'_\varepsilon + \varepsilon \tilde{C}_r(\varepsilon, h_\varepsilon), \end{aligned}$$

where $\tilde{C}_r(\varepsilon, h_\varepsilon) := C_r(\varepsilon, q_\varepsilon)$ since it does not depend on the angle, and $\tilde{M}_{b,r}$ is a bounded function on sets for which $d(B(h_\varepsilon, \varepsilon), \partial\Omega) > \delta$. The conclusion follows. \square

To improve the estimates on h'_ε , we obtain modulated energy estimates. The following lemma relies on straightforward computations.

Lemma 7.4. *Let h_ε satisfy the assumptions of Theorem 2.8. Then, as long as the solution exists,*

$$\frac{d}{dt} \frac{1}{2} (\varepsilon^2 \pi + \varepsilon^\alpha m_1) |h'_\varepsilon - u^\Omega(h_\varepsilon)|^2 = \varepsilon^{\min(2,\alpha)} (h'_\varepsilon - u^\Omega(h_\varepsilon)) \cdot F_{b,r}(\varepsilon, h_\varepsilon, h'_\varepsilon). \quad (7.4)$$

Moreover one has the following immediate estimate.

Lemma 7.5. *There exists $\delta > 0$, ε_0 in $(0, 1)$ and $K > 0$ such that for any $\varepsilon \in (0, \varepsilon_0)$, any $h \in \Omega$ such that $d(B(h, \varepsilon), \partial\Omega) > \delta$, one has $|u^\Omega(h)|_{\mathbb{R}^2} \leq K$.*

One concludes that Corollary 7.3 can be improved into the following.

Corollary 7.6. *Let h_ε satisfy the assumptions of Theorem 2.8. Let $\delta > 0$. There exists $K > 0$ (depending on \mathcal{S}_0 , Ω , ℓ_0 , γ , m_1 , δ) and $\varepsilon_0 > 0$ such that for any ε in $(0, \varepsilon_0)$, for any $t > 0$ such that $d(B(h_\varepsilon, \varepsilon), \partial\Omega) > \delta$ one has $|h'_\varepsilon|_{\mathbb{R}^2} \leq K$.*

Thanks to Corollary 7.6, Lemma 6.8 remains true in the case under view, that it, the time of existence T_ε of h_ε is bounded from below by some $\underline{T} > 0$. Now we have the following local convergence result of h_ε toward \bar{h} , where we recall that \bar{h} is the global solution of (2.45).

Lemma 7.7. *Let $\varepsilon_1 > 0$, $\delta > 0$ and $T > 0$, and suppose that for any ε in $(0, \varepsilon_1)$, we have*

$$d(B(h_\varepsilon(t), \varepsilon), \partial\Omega) > \delta \text{ on } [0, T]. \quad (7.5)$$

Then $h_\varepsilon \rightharpoonup \bar{h}$ in $W^{1,\infty}([0, \tilde{T}]; \mathbb{R}^2)$ weak- \star .

Proof Given $\delta > 0$, $T > 0$ and $\varepsilon_1 > 0$ as above we apply Corollary 7.6 on $[0, T]$ so that reducing $\varepsilon_1 > 0$ if necessary, one has

$$(|h'_\varepsilon|)_{\varepsilon \in (0, \varepsilon_1)} \text{ is bounded uniformly on } [0, T]. \quad (7.6)$$

Our goal is to pass to the limit in each term of (7.10). For what concerns the left hand side it is obvious that $(\varepsilon^\alpha m_1 + \varepsilon^2 \pi) (h'_\varepsilon - u^\Omega(h_\varepsilon))' \rightarrow 0$ in $W^{-1,\infty}$. Next, the term $\varepsilon^{\min(2,\alpha)} F_{b,r}(\varepsilon, h_\varepsilon, h'_\varepsilon)$ converges to 0 in L^∞ . Hence we infer that $h'_\varepsilon - \gamma u^\Omega(h_\varepsilon)$ converges weakly to 0 in $W^{-1,\infty}$. Due to the a priori estimate, this convergence occurs in L^∞ weak- \star . By (7.6), one has a subsequence of (h_ε) satisfying $h_{\varepsilon_n} \rightharpoonup h_*$ in $W^{1,\infty}$ weak- \star . In particular convergence of h_ε towards h_* is strong in L^∞ , and with the convergence of $h'_\varepsilon - \gamma u^\Omega(h_\varepsilon)$ we deduce that $h'_* = \gamma u^\Omega(h_*)$ and $h_*(0) = 0$. The uniqueness of the solution of this Cauchy problem gives $h_* = \bar{h}$ and that the whole (h_ε) converges toward \bar{h} as $\varepsilon \rightarrow 0^+$. This concludes the proof of Lemma 7.7. \square

Finally we briefly conclude the proof of Theorem 2.10 in the case of a homogeneous disk, which is the same as the one of Theorem 2.9, except that in the case (ii) under view, as mentioned below (2.45), the solution \bar{h} is global in time, and in particular that there is no collision of the point vortex with

the external boundary $\partial\Omega$. Hence here, compared to the end of the proof of Theorem 2.9 at the end of Subsection 6.3, we can pick any $T > 0$, and then we define \bar{d} so that for all $t \in [0, T+1]$, $d(\bar{h}(t), \partial\Omega) \geq \bar{d}$ rather than by (6.19) and prove $\bar{T} \geq T + \frac{1}{2}$ rather than (6.21). Of course we rely on Corollary 7.6 rather than Corollary 6.7 and on Lemma 7.7 rather than Lemma 6.9. This ends the proof of Theorem 2.10 in the case of a homogeneous disk.

7.2 Geodesic-gyroscopic normal form

In Case (ii) we will establish that (6.1) can be put into a normal form whose structure looks like (4.16) up to a refined modulation of the velocity of the center of mass. Indeed, in the same way as we defined the Kirchhoff-Routh velocity u^Ω by $u^\Omega = \nabla^\perp \psi^\Omega$ we introduce the corrector velocity u_c corresponding to the stream function ψ_c defined in (6.10) by

$$u_c(q) := \nabla_h^\perp \psi_c(q). \quad (7.7)$$

Observe that the function u_c depends on Ω , \mathcal{S}_0 , ϑ and on h , whereas u^Ω depends only on Ω and h . We will make use in a crucial way of the following second order modulation:

$$\tilde{p}_\varepsilon := (\varepsilon\vartheta'_\varepsilon, h'_\varepsilon - \gamma[u^\Omega(h_\varepsilon) + \varepsilon u_c(q_\varepsilon)])^t, \quad (7.8)$$

which can be compared to the expression (4.28) of p_ε .

Two remarks are in order. We first observe that $\gamma(u^\Omega(h) + \varepsilon u_c(q))$ is the beginning of the expansion of $-\frac{1}{\gamma} \nabla_h^\perp U_\varepsilon(q)$ where $U_\varepsilon(q)$ is the electric-type potential energy defined in (6.9), see (6.11). This modulation is therefore driven by the leading terms of the electric-type potential.

Observe also that, as long as the solid does not touch the boundary, the drift term in the velocity of the center of mass is bounded. Indeed the following refinement of Lemma 7.5 is a direct consequence of Lemma 6.1 and of the definitions of u^Ω and u_c .

Lemma 7.8. *There exists $\delta > 0$, ε_0 in $(0, 1)$ and $K > 0$ such that for any (ε, q) in $\mathcal{Q}_{\delta, \varepsilon_0}$ with $q = (\vartheta, h)$, $|u^\Omega(h) + \varepsilon u_c(q)|_{\mathbb{R}^3} \leq K$.*

Before stating our normal form, we introduce the following definition.

Definition 7.9. Let $\delta > 0$ and $\varepsilon_0 \in (0, 1)$ be given. We say that a vector field F in $C^\infty(\mathbb{R} \times \Omega; \mathbb{R}^3)$ is weakly gyroscopic if there exists $K > 0$ depending on \mathcal{S}_0 , Ω , γ and δ such that for any smooth curve $q(t) = (\vartheta(t), h(t))$ in $\mathbb{R} \times \Omega_\delta$, we have, for any $t \geq 0$ and any $\varepsilon \in (0, \varepsilon_0)$

$$\left| \int_0^t \tilde{p} \cdot F(q) \right| \leq \varepsilon K \left(1 + t + \int_0^t |\tilde{p}|_{\mathbb{R}^3}^2 \right), \quad (7.9)$$

with $\tilde{p} = (\varepsilon\vartheta', h' - \gamma[u^\Omega(h) + \varepsilon u_c(q)])^t$.

The weakly gyroscopic vector fields in the sequel have the form $F = (F_1, 0, 0)^t$. The normal form that we use in Case (ii) is as follows.

Proposition 7.10. *Let $\delta > 0$. There exists ε_0 in $(0, 1)$ and*

- F_r in $L^\infty(\Omega_{\delta, \varepsilon_0} \times \mathbb{R}^3; \mathbb{R}^3)$ depending on S_0 , γ and Ω , weakly nonlinear in the sense of Definition 6.2,
- \mathbf{E}_1^b in $C^\infty(\mathbb{R} \times \Omega; \mathbb{R}^3)$ depending on S_0 and Ω , weakly gyroscopic in the sense of Definition 7.9,

such that Equation (6.1) can be recast as

$$\begin{aligned} (\varepsilon^\alpha M_g + \varepsilon^2 M_{a, \vartheta}^{S_0}) \tilde{p}'_\varepsilon + \varepsilon \langle \Gamma_{\vartheta_\varepsilon}^{S_0}, \tilde{p}_\varepsilon, \tilde{p}_\varepsilon \rangle = F_{\vartheta_\varepsilon}^{S_0}(\tilde{p}_\varepsilon) + \varepsilon \gamma^2 \mathbf{E}_1^b(q_\varepsilon) \\ + \varepsilon^{\min(2, \alpha)} F_r(\varepsilon, q_\varepsilon, \tilde{p}_\varepsilon). \end{aligned} \quad (7.10)$$

We recall that $M_{a, \vartheta}^{S_0}$, $\Gamma_{\vartheta}^{S_0}$ and the force term $F_{\vartheta}^{S_0}$ were defined in the case without outer boundary in (4.9), (4.11) and (4.13) (see also (4.15)) respectively. Moreover the term \mathbf{E}_1^b is explicit, see (9.15).

- **Motivations.** We will use the normal form (7.10) both in order to get a uniform bound of the velocity and to pass to the limit in Case (ii). It would be actually possible to deal with the case where α is small with a less accurate normal form and still get an energy estimate. In particular in order to get a uniform bound of the velocity in Case (ii) we will perform an estimate on an energy adapted to the normal form (7.10). Observe that should the right hand side vanish the normal form (7.10) would be the geodesic equation associated with the metric $M_{\vartheta}(\varepsilon)$. On the other hand the right hand side is the sum of terms with a quite remarkable structure: the leading term $F_{\vartheta_\varepsilon}^{S_0}(\tilde{p}_\varepsilon)$ is gyroscopic in the sense of Definition 4.1, the electric-type term $\mathbf{E}_1^b(q_\varepsilon)$ is weakly gyroscopic in the sense of Definition 7.9; and the remainder F_r is weakly nonlinear.
- **Ideas of the proof of Proposition 7.10.** As for Proposition 6.3 the proof of Proposition 7.10 relies on expansions of the inertia matrix, of the Christoffel symbols and of the force terms with respect to ε . A striking and crucial phenomenon is that some subprincipal contributions (that is, of order ε) of the force terms will be gathered with the leading part of the term involving the Christoffel symbols to become a part of the second term of the left hand side of (7.10), see Lemma 9.9. The leading part of the contribution coming from the Christoffel symbols will be provided by the F^{rot} -part of the decomposition (2.27).

Remark 7.11. The normal forms above are inspired by the case without external boundary (see equation (4.16)) and by the paper [2] where the authors consider the motion of a light charged particle in a slowly varying electromagnetic field. The equation of motion for the particle is an ordinary differential equation involving a small parameter in front of the term with the highest time derivative. In order to restore some uniformity with respect to the small parameter they use a modulation, subtracting to the particle velocity the $|B|^{-2} E \times B$

drift, and a normal form, see [2, Eq. (3.5)], where the only remaining singular term appears through a Lorentz gyroscopic force. This allows to tackle the convergence of the particle motion to the so-called guiding center motion despite the fast oscillations induced by the gyroscopic force.

In the case where the solid \mathcal{S}_0 is a homogeneous disk (Subsection 7.1), $u^\Omega(h_\varepsilon)$ is the restriction to the two last components of this drift. However in the case where \mathcal{S}_0 is not a disk our drift term $(0, \gamma(u^\Omega(h) + \varepsilon u_c(q)))$ does not enter this framework. Actually the use of the $|B|^{-2} E \times B$ drift could give a modulated energy estimate only in the case $\alpha \leq 1$, and in particular not in the case of a solid with a fixed homogeneous density ($\alpha = 2$). Moreover it would not be adapted to the passage to the limit.

7.3 Modulated energy estimates

As mentioned above, in Case (i), Corollary 6.7 provides a uniform bound of h'_ε as long as the body stays at a positive distance from the external boundary. In Case (ii), the same analysis can be carried on so that Corollary 6.7 and Corollary 7.3 have the following counterpart.

Corollary 7.12. *Let $(q_\varepsilon, p_\varepsilon)$ satisfy the assumptions of Theorem 2.10. Let $\delta > 0$. There exists $K > 0$ (depending on $\mathcal{S}_0, \Omega, p_0, \gamma, m_1, \mathcal{J}_1, \delta$) and $\varepsilon_0 > 0$ such that for any ε in $(0, \varepsilon_0)$, for any $t > 0$ such that as long as $(\varepsilon, q_\varepsilon(t))$ belongs to $\mathfrak{Q}_{\delta, \varepsilon_0}$, one has $\varepsilon^{\min(1, \frac{\delta}{2})} |p_\varepsilon|_{\mathbb{R}^3} \leq K$.*

Therefore this analysis does no longer provide a uniform bound of the solid velocity. An important part of the proof consists in finding an appropriate substitute which allows a better control on the body velocity. This will be accomplished below by a modulated energy, which, roughly speaking, consists in applying the energy \mathcal{E}_ϑ (see (4.33)) to \tilde{p}_ε rather than to p_ε .

The structure established in Proposition 7.10 will allow us to obtain an estimate of the modulated energy $\mathcal{E}_\vartheta(\varepsilon, \tilde{p}_\varepsilon)$. Since the equation (7.10) looks like the equation (4.16) of the case without external boundary for which the total energy is the kinetic energy alone (as defined in Proposition 4.5), one may hope to have a good behaviour of the modulated energy $\mathcal{E}_\vartheta(\varepsilon, \tilde{p}_\varepsilon)$ when time proceeds. Indeed we have the following result (compare to Lemma 7.4).

Lemma 7.13. *Let $(q_\varepsilon, p_\varepsilon)$ satisfy the assumptions of Theorem 2.10. Then, as long as the solution exists,*

$$\frac{d}{dt} \mathcal{E}_{\vartheta_\varepsilon}(\varepsilon, \tilde{p}_\varepsilon) = \varepsilon^{\max(1-\alpha, -1)} \gamma^2 \tilde{p}_\varepsilon \cdot \mathbf{E}_1^b(q_\varepsilon) + \tilde{p}_\varepsilon \cdot F_r(\varepsilon, q, \tilde{p}_\varepsilon). \quad (7.11)$$

Proof Using the symmetry of the matrix $M_{\vartheta_\varepsilon}$ defined in (4.32), we get that the time derivative of the modulated energy is

$$\frac{d}{dt} \tilde{\mathcal{E}}_{\vartheta_\varepsilon}(\varepsilon, \tilde{p}_\varepsilon) = \tilde{p}_\varepsilon \cdot M_{\vartheta_\varepsilon}(\varepsilon) \frac{d}{dt} \tilde{p}_\varepsilon + \frac{1}{2} \tilde{p}_\varepsilon \cdot \left(\frac{d}{dt} M_{\vartheta_\varepsilon}(\varepsilon) \right) \tilde{p}_\varepsilon. \quad (7.12)$$

Combining (7.12) and (7.10) we get

$$\begin{aligned} \varepsilon^{\min(2,\alpha)} \frac{d}{dt} \tilde{\mathcal{E}}_{\vartheta_\varepsilon}(\varepsilon, \tilde{p}_\varepsilon) &= \tilde{p}_\varepsilon \cdot F_{\vartheta_\varepsilon}^{\mathcal{S}_0}(\tilde{p}_\varepsilon) \\ &+ \tilde{p}_\varepsilon \cdot \left(\frac{1}{2} \varepsilon^{\min(2,\alpha)} \left(\frac{d}{dt} M_{\vartheta_\varepsilon}(\varepsilon) \right) \tilde{p}_\varepsilon - \varepsilon \langle \Gamma_{\vartheta_\varepsilon}^{\mathcal{S}_0}, \tilde{p}_\varepsilon, \tilde{p}_\varepsilon \rangle \right) \\ &+ \varepsilon \gamma^2 \tilde{p}_\varepsilon \cdot \mathbf{E}_1^b(q_\varepsilon) + \varepsilon^{\min(2,\alpha)} \tilde{p}_\varepsilon \cdot F_r(\varepsilon, q_\varepsilon, \tilde{p}_\varepsilon). \end{aligned} \quad (7.13)$$

Using that the force term is $F_{\vartheta}^{\mathcal{S}_0}(p)$ is gyroscopic in the sense of Definition 4.1 we get that the first term of the right hand side of the equation (7.13) vanishes.

Let us now show that the second term in the right hand side of the equation (7.13) vanishes as well. Going back to the definition of $M_{\vartheta}(\varepsilon)$ we obtain that

$$\varepsilon^{\min(2,\alpha)} \frac{d}{dt} M_{\vartheta_\varepsilon}(\varepsilon) = \varepsilon^2 \frac{d}{dt} M_{a,\vartheta_\varepsilon}^{\mathcal{S}_0} = \varepsilon \frac{\partial M_{a,\vartheta}^{\mathcal{S}_0}}{\partial q}(\vartheta_\varepsilon) \cdot \tilde{p}_\varepsilon. \quad (7.14)$$

Using (4.25) then entails that the second term of the right hand side of the equation (7.13) vanishes. Therefore the equation (7.13) reduces to (7.11). \square

Now Corollary 6.7 and Lemma 7.8 already give us that $\varepsilon \tilde{p}_\varepsilon$ is bounded. Then we use that F_r in $L^\infty(\mathfrak{Q}_{\delta,\varepsilon_0} \times \mathbb{R}^3; \mathbb{R}^3)$ is weakly nonlinear in the sense of Definition 6.2, that \mathbf{E}_1^b is weakly gyroscopic in the sense of Definition 7.9 (using Lemma 6.1), Lemma 4.6 and Gronwall's lemma to get the following result (compare to Corollary 7.6).

Corollary 7.14. *Let $(q_\varepsilon, p_\varepsilon)$ satisfy the assumptions of Theorem 2.9. Let $\delta > 0$. There exists $K > 0$ (depending on \mathcal{S}_0 , Ω , p_0 , γ , m_1 , \mathcal{J}_1 , δ) and $\varepsilon_0 > 0$ such that for any ε in $(0, \varepsilon_0)$, for any $t > 0$ such that as long as $(\varepsilon, q_\varepsilon(t))$ belongs to $\mathfrak{Q}_{\delta,\varepsilon_0}$, one has $|p_\varepsilon|_{\mathbb{R}^3} \leq K$.*

Corollary 7.14 therefore provides the same estimates for Case (ii) as Corollary 6.7 for Case (i).

7.4 Passage to the limit

Here we prove the convergence of the center of mass h_ε to \bar{h} , where \bar{h} is the global solution of (2.45).

First, thanks to Corollary 7.14, Lemma 6.8 remains true in Case (ii). Moreover we have the following counterpart of Lemma 6.9 and Lemma 7.7.

Lemma 7.15. *Let $\varepsilon_1 > 0$, $\delta > 0$ and $T > 0$, and suppose that for any ε in $(0, \varepsilon_1)$, we have*

$$(\varepsilon, q_\varepsilon) \in \mathfrak{Q}_{\delta,\varepsilon_1} \text{ on } [0, T]. \quad (7.15)$$

Then $h_\varepsilon \rightharpoonup \bar{h}$ in $W^{1,\infty}([0, T]; \mathbb{R}^2)$ weak- \star .

As for Lemma 7.7, the proof of Lemma 7.15 consists in passing to the weak limit, with the help of all a priori bounds, in each term of (7.10).

Proof We consider $\delta > 0$, $T > 0$ and $\varepsilon_1 > 0$ as above and apply Corollary 7.14 on the interval $[0, T]$. Reducing $\varepsilon_1 > 0$ if necessary, one has

$$(|h'_\varepsilon| + |\varepsilon \vartheta'_\varepsilon|)_{\varepsilon \in (0, \varepsilon_1)} \text{ is bounded uniformly on } [0, T]. \quad (7.16)$$

Our goal is to pass to the limit in each term of (7.10). For what concerns the left hand side we first observe that, thanks to (7.16), $\varepsilon M_{a, \mathcal{S}_0, \vartheta_\varepsilon}$ is bounded in $W^{1, \infty}$ whereas \tilde{p}'_ε is bounded in $W^{-1, \infty}$. Since M_g is constant, it follows that $(\varepsilon^\alpha M_g + \varepsilon^2 M_{a, \vartheta_\varepsilon}^{\mathcal{S}_0}) \tilde{p}'_\varepsilon \rightarrow 0$ in $W^{-1, \infty}$. Next, the term $\varepsilon \langle \Gamma_{\vartheta_\varepsilon}^{\mathcal{S}_0}, \tilde{p}_\varepsilon, \tilde{p}_\varepsilon \rangle$ converges to 0 in L^∞ since all factors in the brackets are bounded. In the same way, the terms $\varepsilon \gamma^2 \mathbf{E}_1^b(q_\varepsilon)$ and $\varepsilon^{\min(2, \alpha)} F_r(\varepsilon, q_\varepsilon, \tilde{p}_\varepsilon)$ converge strongly to 0 in L^∞ .

Now let us consider the remaining terms in the two last lines of the equation (7.10). These are $\gamma(h'_\varepsilon - \gamma u^\Omega(h_\varepsilon) - \varepsilon \gamma u_c(q_\varepsilon))^\perp - \varepsilon \vartheta'_\varepsilon \zeta_{\vartheta_\varepsilon}$. The last term converges weakly to 0 in $W^{-1, \infty}$ as seen in Case (i), see (4.36). Hence we infer that $h'_\varepsilon - \gamma u^\Omega(h_\varepsilon)$ converges weakly to 0 in $W^{-1, \infty}$. Then one concludes exactly as for Lemma 7.7. \square

Once Lemma 7.15 is established, the conclusion of the proof of Theorem 2.10 (provided that \mathcal{S}_0 is not a disk) follows the exact same lines as in the case of a homogeneous disk. This concludes the proof of Theorem 2.10 in this case.

7.5 The case of a non-homogeneous disk

We now return to the case \mathcal{S}_0 is a disk, but this time we handle the case where it is non-homogeneous. We aim at proving Theorem 2.11. This adds some extra-difficulties in the analysis which require a separate treatment in the case $\alpha > 2$. In the case $\alpha \leq 2$, we can use the above section since the degeneracy of $M_{a, \vartheta}^{\mathcal{S}_0}$ does not prevent $M_\vartheta(\varepsilon)$ to be uniformly bounded from below, so that we can obtain Corollary 7.14 in the same way.

From now on we suppose $\alpha > 2$ and once again we assume without loss of generality $\mathcal{S}_0 = \overline{B}(0, 1)$. We will use yet another slight variant of Definition 6.2 and Definition 7.1.

Definition 7.16. Let $\delta > 0$ and $\varepsilon_0 \in (0, 1)$ be given. We say that a vector field F_b in $L^\infty((0, \varepsilon_0) \times \Omega \times \mathbb{R}^3; \mathbb{R}^2)$ is weakly nonlinear if there exists $K > 0$ depending on m , \mathcal{J} , γ , Ω and δ such that for any ε in $(0, \varepsilon_0)$, any x in Ω such that $d(B(x, \varepsilon), \partial\Omega) > \delta$ and any p in \mathbb{R}^3 , one has

$$|F_b(\varepsilon, x, p)|_{\mathbb{R}^2} \leq K(1 + |p|_{\mathbb{R}^3} + \varepsilon |p|_{\mathbb{R}^3}^2). \quad (7.17)$$

We will use the modulated variable $\tilde{p}_{b, \varepsilon}$ defined by

$$\tilde{p}_{b, \varepsilon} := h'_{c, \varepsilon} - \gamma u^\Omega(h_{c, \varepsilon}),$$

and the following normal form.

Proposition 7.17. *There exists $F_{b,r}$ in $L^\infty((0,1) \times \Omega \times \mathbb{R}^3; \mathbb{R}^2)$, weakly non-linear in the sense of Definition 7.16, such that Equation (2.34) can be recast as*

$$\varepsilon^\alpha m_1 h_\varepsilon'' + \varepsilon^2 \pi(\tilde{p}_{b,\varepsilon})' = \gamma \tilde{p}_{b,\varepsilon}^\perp + \varepsilon^2 F_{b,r}(\varepsilon, h_{c,\varepsilon}, \varepsilon \vartheta_\varepsilon', \tilde{p}_{b,\varepsilon}), \quad (7.18)$$

$$h_{c,\varepsilon} - h_\varepsilon = \varepsilon \zeta_{\vartheta_\varepsilon}, \quad (7.19)$$

$$\mathcal{J}_1 \varepsilon \vartheta_\varepsilon'' = \zeta_{\vartheta_\varepsilon}^\perp \cdot m_1 h_\varepsilon'', \quad (7.20)$$

Again the coefficient π comes from the fact that $M_b^{S_0} = \pi \text{Id}_2$ (where $M_b^{S_0}$ is defined in (4.23)). The proof of Proposition 7.17 is given in Section 9.

Once Proposition 7.17 is obtained, for $\varepsilon \in (0, \varepsilon_0)$, we multiply (7.18) by $\tilde{p}_{b,\varepsilon}$. We notice that

$$m_1 h_\varepsilon'' \cdot h_{c,\varepsilon}' = m_1 h_\varepsilon'' \cdot h_\varepsilon' + \varepsilon^2 \mathcal{J}_1 \vartheta_\varepsilon'' \vartheta_\varepsilon',$$

so we get, denoting $\mathcal{E}_b := \varepsilon^\alpha m_1 |h_\varepsilon'|^2 + \varepsilon^{\alpha+2} \mathcal{J}_1 |\vartheta_\varepsilon'|^2 + \varepsilon^2 \pi |\tilde{p}_{b,\varepsilon}|^2$,

$$\frac{d}{dt} \mathcal{E}_b(t) + \varepsilon^\alpha m_1 h_\varepsilon'' \cdot u^\Omega(h_{c,\varepsilon}) \leq C \varepsilon^2 |\tilde{p}_{b,\varepsilon}| (1 + |\tilde{p}_{b,\varepsilon}| + \varepsilon |\tilde{p}_{b,\varepsilon}|^2).$$

The standard energy estimate gives us here that $\varepsilon |h_{c,\varepsilon}'|$ is bounded, and consequently so is $\varepsilon |\tilde{p}_{b,\varepsilon}|$. Integrating over time and using an integration by parts we obtain (as long as the solid stays at positive distance from $\partial\Omega$):

$$\mathcal{E}_b(t) - \mathcal{E}_b(0) \leq C \varepsilon^2 + C \int_0^t \mathcal{E}_b(s) ds + C \varepsilon^\alpha |h_\varepsilon'(t)| + C \varepsilon^\alpha |h_\varepsilon'(0)|.$$

With Young's inequality, we finally get

$$\mathcal{E}_b(t) \leq C \int_0^t \mathcal{E}_b(s) ds + C \mathcal{E}_b(0) + C \varepsilon^2,$$

so with Gronwall's lemma we finally deduce that $\mathcal{E}_b(t) \leq C(\mathcal{E}_b(0) + \varepsilon^2)$, as long as the solid stays at positive distance from $\partial\Omega$. Since $\alpha > 2$, we deduce that $(h_{c,\varepsilon}')_{\varepsilon \in (0, \varepsilon_0)}$, $(\varepsilon^{\frac{\alpha}{2}-1} h_\varepsilon')_{\varepsilon \in (0, \varepsilon_0)}$ and $(\varepsilon^{\frac{\alpha}{2}} \vartheta_\varepsilon')_{\varepsilon \in (0, \varepsilon_0)}$ are bounded in L^∞ . (Notice in passing that using (7.19) and an interpolation argument we get that $(h_\varepsilon)_{\varepsilon \in (0, \varepsilon_0)}$ is bounded in $W^{\frac{2}{\alpha}, \infty}$.)

Then it is straightforward to adapt the previous reasoning to pass to the weak limit in (7.18). We deduce $\tilde{p}_{b,\varepsilon} \rightharpoonup 0$ in L^∞ and the conclusion follows. This ends the proof of Theorem 2.11.

8 Asymptotic development of the stream and potential functions

In this section, we establish asymptotic expansions for the circulation stream function (defined in (2.10)) and the Kirchhoff potentials (defined in (2.8)) in the domain $\mathcal{F}^\varepsilon(q)$, as ε tends to 0^+ . The asymptotic analysis of the Laplace equation when the size of an inclusion goes to 0 has been deeply studied, cf. for example [27] and [18]. However to our knowledge the results of this section are not covered by the literature.

8.1 A few reminders about single-layer potentials

In order to get the asymptotic expansions hinted above, we will look for a representation of these stream and potential functions as a superposition of single-layer integrals supported by the two connected components $\partial\mathcal{S}^\varepsilon(q)$ and $\partial\Omega$ of the boundary of the fluid domain $\mathcal{F}_\varepsilon(q)$. In this subsection, we give a few reminders about single-layer potentials which we will use in the analysis. We refer for instance to [26] and [6].

Below we consider single-layer potentials of the form:

$$SL[\mathbf{p}^C] := \int_{\mathcal{C}} \mathbf{p}^C(y) G(\cdot - y) ds(y), \quad (8.1)$$

where \mathcal{C} is a smooth Jordan curve in the plane and \mathbf{p}^C belongs to the Sobolev space $H^{-\frac{1}{2}}(\mathcal{C})$. Recall that G was defined in (2.40). We say that \mathcal{C} is the support of the single-layer potential and that \mathbf{p}^C is a density on \mathcal{C} .

- *Harmonicity and trace.* The formula (8.1) defines a function in the Sobolev space $H_{loc}^1(\mathbb{R}^2)$, harmonic in $\mathbb{R}^2 \setminus \mathcal{C}$. In particular, for any \mathbf{p}^C in $H^{-\frac{1}{2}}(\mathcal{C})$, the trace of $SL[\mathbf{p}^C]$ on \mathcal{C} is well-defined as a function of the Sobolev space $H^{\frac{1}{2}}(\mathcal{C})$.

- *Jump of the derivative and density.* The density \mathbf{p}^C is equal to the jump of the normal derivative of $SL[\mathbf{p}^C]$ across \mathcal{C} . In order to state this rigorously let us be specific on the orientation of the normal. According to Jordan's theorem, the set $\mathbb{R}^2 \setminus \mathcal{C}$ has two connected components, one bounded (the interior), say \mathcal{O}_i , and the other one unbounded (the exterior), say \mathcal{O}_e . Moreover the curve \mathcal{C} is the boundary of each component. We consider the restrictions of $SL[\mathbf{p}^C]$:

$$u_i := SL[\mathbf{p}^C]|_{\mathcal{O}_i} \quad \text{and} \quad u_e := SL[\mathbf{p}^C]|_{\mathcal{O}_e}.$$

Denote n_i (respectively n_e) the unit normal on \mathcal{C} outward to \mathcal{O}_i (resp. to \mathcal{O}_e). Then the function u_i (respectively u_e) is harmonic in \mathcal{O}_i (resp. \mathcal{O}_e) and the traces of the normal derivatives $\frac{\partial u_i}{\partial n_i}$ and $\frac{\partial u_e}{\partial n_e}$ on each side of \mathcal{C} are well-defined in $H^{-\frac{1}{2}}(\mathcal{C})$ and satisfy

$$\mathbf{p}^C = \frac{\partial u_i}{\partial n_i} + \frac{\partial u_e}{\partial n_e}. \quad (8.2)$$

In the sequel we will make use of single-layer potentials supported on the external boundary $\partial\Omega$, on the boundary $\partial\mathcal{S}_\varepsilon(q)$ of the solid body and on the boundary $\partial\mathcal{S}_0$ of the rescaled body as well. We will not use the notations n_i nor n_e but rather the notation n which always stands for the normal outward the fluid. Hence we will have to particularly take care of the signs when referring to the formula (8.2).

- *Kernel and rank.* We will use the following facts:

The operator SL is Fredholm with index zero from $L^2(\mathcal{C})$ to $H^1(\mathcal{C})$, (8.3)

If $\mathbf{p}^C \in H^{-\frac{1}{2}}(\mathcal{C})$ satisfies $\int_{\mathcal{C}} \mathbf{p}^C ds = 0$ and $SL[\mathbf{p}^C] = 0$ then $\mathbf{p}^C = 0$, (8.4)

If $\text{Cap}(\mathcal{C}) \neq 1$, then for any $\mathbf{p}^C \in H^{-\frac{1}{2}}(\mathcal{C})$, $SL[\mathbf{p}^C] = 0$ implies $\mathbf{p}^C = 0$. (8.5)

Above, with some slight abuses of notation, we omit to mention the trace operator on \mathcal{C} and we write $\int_{\mathcal{C}} \mathbf{p}^{\mathcal{C}} ds$ for the duality bracket $\langle 1, \mathbf{p}^{\mathcal{C}} \rangle_{H^{-\frac{1}{2}}(\mathcal{C}), H^{\frac{1}{2}}(\mathcal{C})}$. We refer to [26, Th. 7.17] for (8.3), to [26, Th. 8.12] for (8.4) and to [26, Th. 8.16] for (8.5).

These properties have the following consequences in our context. We will assume without loss of generality that the logarithmic capacity² $\text{Cap}(\partial\Omega)$ of $\partial\Omega$ satisfies $\text{Cap}(\partial\Omega) < 1$, using translation and dilatation of the coordinates system if necessary. Observe that the monotony property of the logarithmic capacity entails that $\text{Cap}(\partial\mathcal{S}_0) < 1$. Using this latter property, we deduce the two following results.

Proposition 8.1. *There exists a unique smooth function $\psi_{-1}^{\mathcal{S}_0}$ solution of (4.4). Moreover, it satisfies*

$$\psi_{-1}^{\mathcal{S}_0} = SL[\mathbf{p}_{-1}^{\mathcal{S}_0}], \quad \text{with } \mathbf{p}_{-1}^{\mathcal{S}_0} = \frac{\partial \psi_{-1}^{\mathcal{S}_0}}{\partial n}, \quad (8.6)$$

In potential theory $-\mathbf{p}_{-1}^{\mathcal{S}_0}$ is called the equilibrium density of $\partial\mathcal{S}_0$. The constant $C^{\mathcal{S}_0}$ in (4.4) is given by $C^{\mathcal{S}_0} = \frac{1}{2\pi} \ln(\text{Cap}(\partial\mathcal{S}_0))$. Note that the index -1 is in italic type to emphasize the fact that it is related to an asymptotic development of ψ in powers of ε and is not a coordinate.

Proposition 8.2. *Let g be a smooth function on $\partial\mathcal{S}_0$ such that*

$$\int_{\partial\mathcal{S}_0} g \mathbf{p}_{-1}^{\mathcal{S}_0} ds = 0. \quad (8.7)$$

Then there exists a unique bounded smooth function f such that

$$-\Delta f = 0 \quad \text{in } \mathbb{R}^2 \setminus \mathcal{S}_0, \quad \text{and } f = g \quad \text{on } \partial\mathcal{S}_0. \quad (8.8)$$

Moreover, there exists a unique smooth density $p^{\partial\mathcal{S}_0}$ in $C^\infty(\partial\mathcal{S}_0)$ such that $f = SL[p^{\partial\mathcal{S}_0}]$ and

$$\int_{\partial\mathcal{S}_0} p^{\partial\mathcal{S}_0} ds = 0. \quad (8.9)$$

Finally, $f = O(|x|_{\mathbb{R}^2}^{-1})$ at infinity and

$$\int_{\partial\mathcal{S}_0} \frac{\partial f}{\partial n} ds = 0. \quad (8.10)$$

Proof of Proposition 8.1 and 8.2. The uniqueness part of Proposition 8.1 and of Proposition 8.2 and the decaying at infinity in Proposition 8.2 can be established by considering holomorphy at infinity of appropriate functions, see for instance [6, Prop. 2.74. and Prop. 3.2.].

The existence part of Proposition 8.1 is given in [26, Th. 8.15]; it also follows from the properties of the single-layers potentials recalled above, in particular (8.3) and (8.5).

² also called external conformal radius or transfinite diameter in other contexts [34]

Regarding the existence part of Proposition 8.2 we proceed in two steps.

First we prove that the operator which maps $(p^{\partial\mathcal{S}_0}, C)$ in $L^2(\mathcal{S}_0) \times \mathbb{R}$ to $(SL[p^{\partial\mathcal{S}_0}] - C, \int_{\partial\mathcal{S}_0} p^{\partial\mathcal{S}_0} ds)$ in $H^1(\mathcal{S}_0) \times \mathbb{R}$ is invertible. In order to prove this, we observe that this operator is Fredholm with index zero as a consequence of (8.3). Moreover if $(p^{\partial\mathcal{S}_0}, C)$ is in the kernel of this operator, then $SL[p^{\partial\mathcal{S}_0} - \frac{C}{C^{\mathcal{S}_0}} \mathbf{p}_{-1}^{\mathcal{S}_0}] = 0$, so that according to (8.4), $p^{\partial\mathcal{S}_0} = \frac{C}{C^{\mathcal{S}_0}} \mathbf{p}_{-1}^{\mathcal{S}_0}$. Now we have

$$\int_{\partial\mathcal{S}_0} \mathbf{p}_{-1}^{\mathcal{S}_0} ds = -1, \quad (8.11)$$

as a consequence of (4.4d) and of the second identity of (8.6). Then using that $\int_{\partial\mathcal{S}_0} p^{\partial\mathcal{S}_0} ds = 0$ we deduce that $C = 0$ and therefore $p^{\partial\mathcal{S}_0} = 0$ as well.

Then $(g, 0)$ is in the image of this operator, that is there exists $(p^{\partial\mathcal{S}_0}, C)$ in $L^2(\mathcal{S}_0) \times \mathbb{R}$ such that

$$SL[p^{\partial\mathcal{S}_0}] - C = g \text{ on } \partial\mathcal{S}_0, \quad (8.12)$$

and (8.9). Observing that the trace of the operator SL on $\partial\mathcal{S}_0$ is self-adjoint we infer that

$$\int_{\partial\mathcal{S}_0} SL[p^{\partial\mathcal{S}_0}] \mathbf{p}_{-1}^{\mathcal{S}_0} ds = \int_{\partial\mathcal{S}_0} p^{\partial\mathcal{S}_0} SL[\mathbf{p}_{-1}^{\mathcal{S}_0}] ds = C^{\mathcal{S}_0} \int_{\partial\mathcal{S}_0} p^{\partial\mathcal{S}_0} ds = 0. \quad (8.13)$$

Combining (8.7), (8.11), (8.12) and (8.13) we infer that $C = 0$.

Finally the smoothness part of Proposition 8.1 and of Proposition 8.2 follows from [26, Th. 7.16] and (8.10) follows from (8.9), (8.2) and the vanishing by integration by parts of the interior contribution. \square

• *Regular integral operators.* Since we consider single-layer potentials supported on two disjoint curves and their values on both curves, we will also be led to consider regular integral operators. We recall below some straightforward results which are useful in the sequel. Given \mathcal{C} a smooth Jordan curve in $\overline{\Omega}$, we introduce, for $\delta > 0$,

$$\mathcal{C}^\delta := \{x \in \overline{\Omega} / \text{dist}(x, \mathcal{C}) < \delta\},$$

and define

$$F_\delta : C^1(\overline{\Omega} \setminus \mathcal{C}^\delta; H^{\frac{1}{2}}(\mathcal{C})) \times H^{-\frac{1}{2}}(\mathcal{C}) \rightarrow C^1(\overline{\Omega} \setminus \mathcal{C}^\delta; \mathbb{R}),$$

by setting, for any $(b, \mathbf{p}^\mathcal{C})$ in $C^1(\overline{\Omega} \setminus \mathcal{C}^\delta; H^{\frac{1}{2}}(\mathcal{C})) \times H^{-\frac{1}{2}}(\mathcal{C})$, for any x in $\overline{\Omega} \setminus \mathcal{C}^\delta$,

$$F[b, \mathbf{p}^\mathcal{C}](x) := \int_{\mathcal{C}} b(x, y) \mathbf{p}^\mathcal{C}(y) ds(y).$$

This will be applied to b defined in a larger set but singular for $x = y$; this motivates our framework for b .

Next, given another smooth Jordan curve $\tilde{\mathcal{C}}$ in $\overline{\Omega} \setminus \mathcal{C}^\delta$ and for b in $C^1(\overline{\Omega} \setminus \mathcal{C}^\delta; H^{\frac{1}{2}}(\mathcal{C}))$, we define the operator

$$F_{\delta,b} : L^2(\mathcal{C}) \rightarrow H^1(\tilde{\mathcal{C}})$$

by setting $F_{\delta,b}(\mathbf{p}^{\mathcal{C}})$ as the trace of $F_\delta[b, \mathbf{p}^{\mathcal{C}}]$ on $\tilde{\mathcal{C}}$. We will make use of the following lemma.

Lemma 8.3. *Let $\delta > 0$. The two following properties hold.*

- (i) *The operator F_δ is bilinear continuous with a norm less than 1, in other words: for any $(b, \mathbf{p}^{\mathcal{C}})$ in $C^1(\overline{\Omega} \setminus \mathcal{C}^\delta; H^{\frac{1}{2}}(\mathcal{C})) \times H^{-\frac{1}{2}}(\mathcal{C})$, one has*

$$\|F_\delta[b, \mathbf{p}^{\mathcal{C}}]\|_{C^1(\overline{\Omega} \setminus \mathcal{C}^\delta)} \leq \|b\|_{C^1(\overline{\Omega} \setminus \mathcal{C}^\delta; H^{\frac{1}{2}}(\mathcal{C}))} \|\mathbf{p}^{\mathcal{C}}\|_{H^{-\frac{1}{2}}(\mathcal{C})}.$$

- (ii) *If $\tilde{\mathcal{C}}$ is a smooth Jordan curve in $\overline{\Omega} \setminus \mathcal{C}^\delta$ and b in $C^1(\overline{\Omega} \setminus \mathcal{C}^\delta; H^{\frac{1}{2}}(\mathcal{C}))$, the operator $F_{\delta,b}$ is compact from $L^2(\mathcal{C})$ to $H^1(\tilde{\mathcal{C}})$.*

The proof of Lemma 8.3 is elementary and left to the reader.

8.2 Statements of the results

8.2.1 Circulation part

Let (ε, q) in Ω . We denote $\psi_\varepsilon(q, \cdot)$ in $\mathcal{F}_\varepsilon(q)$ the function defined as the solution to the Dirichlet boundary value problem:

$$-\Delta \psi_\varepsilon(q, \cdot) = 0 \quad \text{in } \mathcal{F}_\varepsilon(q), \quad (8.14a)$$

$$\psi_\varepsilon(q, \cdot) = C_\varepsilon(q) \quad \text{on } \partial \mathcal{S}_\varepsilon(q), \quad (8.14b)$$

$$\psi_\varepsilon(q, \cdot) = 0 \quad \text{on } \partial \Omega, \quad (8.14c)$$

where the constant $C_\varepsilon(q)$ is such that:

$$\int_{\partial \mathcal{S}_\varepsilon(q)} \frac{\partial \psi_\varepsilon}{\partial n}(q, \cdot) \, ds = -1. \quad (8.14d)$$

Here, n stands for the unit normal vector to $\partial \mathcal{S}_\varepsilon(q) \cup \partial \Omega$ directed toward the exterior of $\mathcal{F}_\varepsilon(q)$. The function ψ_ε is the counterpart, for the case where the size of the solid is of order ε , of the function ψ defined in (2.10) in the case where the size of the solid is of order 1. For any q in Ω , the existence and uniqueness of a solution $\psi_\varepsilon(q, \cdot)$ of (8.14) is classical.

In order to state a result establishing an asymptotic expansion of $C_\varepsilon(q)$ and of $\frac{\partial \psi_\varepsilon}{\partial n}(q, \cdot)$ on $\partial \mathcal{S}_\varepsilon(q)$ as $\varepsilon \rightarrow 0^+$ we introduce a few notations.

- *Definition of $\psi_0^{S_0}(q, \cdot)$ and of $P_0(q, X)$.* We denote, for any $q := (\vartheta, h)$ in $\mathbb{R} \times \Omega$, by $P_0(q, X)$ the harmonic polynomial

$$P_0(q, X) := u^\Omega(h)^\perp \cdot (R(\vartheta)X - \zeta_\vartheta). \quad (8.15)$$

Let us recall that ζ_ϑ is defined in (4.14) in terms of ζ defined in (4.12). Recalling (8.11) and the second identity of (8.6) we observe that $P_0(q, X)$ satisfies

$$\int_{\partial\mathcal{S}_0} P_0(q, \cdot) \mathbf{p}_{-I}^{S_0} ds = 0. \quad (8.16)$$

Therefore, according to Proposition 8.2 there exists a unique smooth function $\psi_0^{S_0}(q, \cdot)$ satisfying

$$-\Delta\psi_0^{S_0}(q, \cdot) = 0 \quad \text{in } \mathbb{R}^2 \setminus \mathcal{S}_0, \quad (8.17a)$$

$$\psi_0^{S_0}(q, \cdot) = P_0(q, \cdot) \quad \text{on } \partial\mathcal{S}_0, \quad (8.17b)$$

and vanishing at infinity. Moreover

$$\int_{\partial\mathcal{S}_0} \frac{\partial\psi_0^{S_0}}{\partial n} ds = 0. \quad (8.17c)$$

- *Definition of $\psi_I^\Omega(q, \cdot)$.* We also introduce the solution $\psi_I^\Omega(q, \cdot)$ of

$$-\Delta\psi_I^\Omega(q, \cdot) = 0 \quad \text{in } \Omega, \quad (8.18a)$$

$$\psi_I^\Omega(q, \cdot) = -(\nabla G)(\cdot - h) \cdot \zeta_\vartheta \quad \text{on } \partial\Omega. \quad (8.18b)$$

Above G denotes the Newtonian potential defined in (2.40). The function ψ_I^Ω can be expressed thanks to the function ψ_0^Ω defined in (2.39) according to the following formula:

$$\forall h, x \in \Omega, \quad \forall \vartheta \in \mathbb{R}, \quad D_x\psi_0^\Omega(h, x) \cdot \zeta_\vartheta = \psi_I^\Omega(\vartheta, x, h). \quad (8.19)$$

Proof of (8.19). We first recall that ψ_0^Ω is symmetric in its variables. Indeed by uniqueness of the Dirichlet problem (2.39) we have for any h in Ω , the decomposition in Ω :

$$\psi_0^\Omega(h, \cdot) = G(\cdot - h) + G^\Omega(h, \cdot), \quad (8.20)$$

where G^Ω denotes the Green function associated with the domain Ω and the homogeneous Dirichlet condition, that is

$$\Delta G^\Omega(h, \cdot) = \delta_h \text{ in } \Omega, \quad G^\Omega(h, \cdot) = 0 \text{ on } \partial\Omega.$$

Using the decomposition (8.20), that the Newtonian potential G is even and the symmetry of G^Ω we get

$$\forall h, x \in \Omega, \quad \psi_0^\Omega(h, x) = \psi_0^\Omega(x, h). \quad (8.21)$$

It follows that $(D_x\psi_0^\Omega)(x, h) \cdot \zeta_\vartheta = (D_h\psi_0^\Omega)(h, x) \cdot \zeta_\vartheta$. Next we observe that $D_h\psi_0^\Omega(h, \cdot) \cdot \zeta_\vartheta$ satisfies the same Dirichlet problem as $\psi_I^\Omega(\vartheta, h, \cdot)$, by derivation of (2.39). Formula (8.19) follows then from the uniqueness of solutions to the Dirichlet problem, after switching h and x . \square

– *Definition of $\psi_1^{S_0}(q, \cdot)$ and of $P_1(q, X)$.* Let us denote, for any $q := (\vartheta, h)$ in $\mathbb{R} \times \Omega$, by $P_1(q, X)$ the polynomial defined by

$$P_1(q, X) := -\frac{1}{2} \langle R(\vartheta)^t D_x^2 \psi_0^\Omega(h, h) R(\vartheta), T^2(\mathbf{p}_{-1}^{S_0}) + X^{\otimes 2} \rangle_{\mathbb{R}^{2 \times 2}} + R(\vartheta)^t D_x \psi_1^\Omega(q, h) \cdot (\zeta - X). \quad (8.22)$$

Above $D_x^2 \psi_0^\Omega(h, h)$ denotes the second derivative of $\psi_0^\Omega(h, \cdot)$ evaluated in h , $D_x \psi_1^\Omega(q, h)$ stands for the derivative of $\psi_1^\Omega(q, \cdot)$ evaluated in h and $X^{\otimes 2}$ stands for the 2×2 matrix $X \otimes X$. The notation $T^2(\mathbf{p}_{-1}^{S_0})$ stands for

$$T^2(\mathbf{p}_{-1}^{S_0}) := \int_{\partial \mathcal{S}_0} \frac{\partial \psi_{-1}^{S_0}}{\partial n}(X) X^{\otimes 2} ds(X). \quad (8.23)$$

This notation is justified by (8.6). Observe that $P_1(q, X)$ is harmonic (since every monomial of Taylor's expansions of harmonic functions are themselves harmonic) and satisfies

$$\int_{\partial \mathcal{S}_0} P_1(q, \cdot) \mathbf{p}_{-1}^{S_0} ds = 0. \quad (8.24)$$

Therefore, according to Proposition 8.2 there exists a unique smooth function $\psi_1^{S_0}(q, \cdot)$ satisfying

$$-\Delta \psi_1^{S_0}(q, \cdot) = 0 \quad \text{in } \mathbb{R}^2 \setminus \mathcal{S}_0, \quad (8.25a)$$

$$\psi_1^{S_0}(q, \cdot) = P_1(q, \cdot) \quad \text{on } \partial \mathcal{S}_0, \quad (8.25b)$$

and vanishing at infinity. Moreover

$$\int_{\partial \mathcal{S}_0} \frac{\partial \psi_1^{S_0}}{\partial n}(q, \cdot) ds = 0. \quad (8.25c)$$

Our main result regarding the potential ψ_ε is the following, in addition to Lemma 6.4.

Proposition 8.4. *Let $\delta > 0$. There exist ε_0 in $(0, 1)$ and a function $\mathbf{p}_r^{\partial \mathcal{S}_0}$ in $L^\infty(\Omega_{\delta, \varepsilon_0}; L^2(\partial \mathcal{S}_0; \mathbb{R}))$, depending only on \mathcal{S}_0 and Ω , such that for any (ε, q) in $\Omega_{\delta, \varepsilon_0}$ and for any X in $\partial \mathcal{S}_0$, one has*

$$\begin{aligned} \frac{\partial \psi_\varepsilon}{\partial n}(q, \varepsilon R(\vartheta)X + h) &= \frac{1}{\varepsilon} \frac{\partial \psi_{-1}^{S_0}}{\partial n}(X) + \left(\frac{\partial \psi_0^{S_0}}{\partial n}(q, X) - R(\vartheta)^t u^\Omega(h) \cdot \tau \right) \\ &+ \varepsilon \left(\frac{\partial \psi_1^{S_0}}{\partial n} - \frac{\partial P_1}{\partial n} \right)(q, X) + \varepsilon^2 \mathbf{p}_r^{\partial \mathcal{S}_0}(\varepsilon, q, X). \end{aligned} \quad (8.26)$$

We recall that the set $\Omega_{\delta, \varepsilon_0}$ was defined in (6.4).

The proofs of Lemma 6.4 and of Proposition 8.4 are gathered in Section 8.3.

8.2.2 Potential part

For any $j = 1, 2, 3$, for any q in \mathcal{Q} , we consider the functions $K_{j,\varepsilon}(q, \cdot)$ on $\partial\Omega \cup \partial\mathcal{S}_\varepsilon(q)$ given by:

$$K_{j,\varepsilon}(q, \cdot) := n \cdot \xi_j(q, \cdot) \text{ on } \partial\Omega \cup \partial\mathcal{S}_\varepsilon(q), \quad (8.27)$$

where n denotes the unit normal to $\partial\mathcal{S}_\varepsilon(q) \cup \partial\Omega$, pointing outside $\mathcal{F}^\varepsilon(q)$ and the functions $\xi_j(q, \cdot)$ are given by the formula (2.6). Then the Kirchhoff's potentials $\varphi_{j,\varepsilon}(q, \cdot)$, for $j = 1, 2, 3$, are the unique (up to an additive constant) solutions in $\mathcal{F}^\varepsilon(q)$ of the following Neumann problem:

$$\Delta\varphi_{j,\varepsilon}(q, \cdot) = 0 \quad \text{in } \mathcal{F}^\varepsilon(q), \quad (8.28a)$$

$$\frac{\partial\varphi_{j,\varepsilon}}{\partial n}(q, \cdot) = K_{j,\varepsilon}(q, \cdot) \quad \text{on } \partial\mathcal{S}_\varepsilon(q), \quad (8.28b)$$

$$\frac{\partial\varphi_{j,\varepsilon}}{\partial n}(q, \cdot) = 0 \quad \text{on } \partial\Omega. \quad (8.28c)$$

The functions $K_{j,\varepsilon}(q, \cdot)$ (respectively $\varphi_{j,\varepsilon}(q, \cdot)$) are the counterpart, for the case where the size of the solid is of order ε , of the functions defined in (2.7) (resp. in (2.8)) in the case where the size of the solid is of order 1.

We will use the vector notations:

$$\boldsymbol{\varphi}_\varepsilon = (\varphi_{1,\varepsilon}, \varphi_{2,\varepsilon}, \varphi_{3,\varepsilon})^t \text{ and } \mathbf{K}_\varepsilon = (K_{1,\varepsilon}, K_{2,\varepsilon}, K_{3,\varepsilon})^t. \quad (8.29)$$

Our result on the expansion of the Kirchhoff potentials $\varphi_{j,\varepsilon}$ is the following.

Proposition 8.5. *Let $\delta > 0$. There exists ε_0 in $(0, 1)$ and*

- (i) *there exists $\boldsymbol{\varphi}_r$ in $L^\infty(\mathfrak{Q}_{\delta,\varepsilon_0}; L^2(\partial\mathcal{S}_0; \mathbb{R}^3))$ and \check{c} in $L^\infty(\mathfrak{Q}_{\delta,\varepsilon_0}; \mathbb{R}^3)$ such that for any (ε, q) in $\mathfrak{Q}_{\delta,\varepsilon_0}$, with $q = (\vartheta, h)$, for any X in $\partial\mathcal{S}_0$,*

$$\boldsymbol{\varphi}_\varepsilon(q, \varepsilon R(\vartheta)X + h) = \varepsilon I_\varepsilon \mathcal{R}(\vartheta) \left(\boldsymbol{\varphi}^{\mathcal{S}_0}(X) + \check{c}(\varepsilon, q) + \varepsilon^2 \boldsymbol{\varphi}_r(\varepsilon, q, X) \right), \quad (8.30)$$

- (ii) *there exists $\mathbf{p}_r^{\partial\mathcal{S}_0}$ in $L^\infty(\mathfrak{Q}_{\delta,\varepsilon_0}; L^2(\partial\mathcal{S}_0; \mathbb{R}^3))$ such that for any (ε, q) in $\mathfrak{Q}_{\delta,\varepsilon_0}$, with $q = (\vartheta, h)$, for any X in $\partial\mathcal{S}_0$,*

$$\mathcal{R}(\vartheta)^t \frac{\partial\boldsymbol{\varphi}_\varepsilon}{\partial\tau}(q, \varepsilon R(\vartheta)X + h) = I_\varepsilon \left(\frac{\partial\boldsymbol{\varphi}^{\mathcal{S}_0}}{\partial\tau}(X) + \varepsilon^2 \mathbf{p}_r^{\partial\mathcal{S}_0}(\varepsilon, q, X) \right), \quad (8.31)$$

- (iii) *there exists $\mathbf{p}_r^{\partial\Omega}$ in $L^\infty(\mathfrak{Q}_{\delta,\varepsilon_0}; L^2(\partial\Omega; \mathbb{R}^3))$ such that for any (ε, q) in $\mathfrak{Q}_{\delta,\varepsilon_0}$, for any x in $\partial\Omega$,*

$$\frac{\partial\boldsymbol{\varphi}_\varepsilon}{\partial\tau}(q, x) = I_\varepsilon \varepsilon^2 \mathbf{p}_r^{\partial\Omega}(\varepsilon, q, x). \quad (8.32)$$

Moreover the remainders $\boldsymbol{\varphi}_r$, $\mathbf{p}_r^{\partial\mathcal{S}_0}$ and $\mathbf{p}_r^{\partial\Omega}$ depend only on \mathcal{S}_0 and Ω .

We recall that $\boldsymbol{\varphi}^{\mathcal{S}_0}$ was defined in (4.2). The proof of Proposition 8.5 is given in Section 8.5.

8.3 Asymptotic expansion of the circulation part: Proof of Proposition 8.4 and of Lemma 6.4

In this section we prove Proposition 8.4 and of Lemma 6.4. We proceed in four steps that we now detail. We rely on intermediate results: Lemma 8.6, Lemma 8.7, Lemma 8.9, Lemma 8.10, and Lemma 8.11 whose proofs are postponed to Subsection 8.4.

We will use the following functional space: for $-\frac{1}{2} \leq s \leq 1$, let the Hilbert space

$$F_s := H^s(\partial\mathcal{S}_0) \times H^s(\partial\Omega) \times \mathbb{R}.$$

We will mainly make use of the indices $s = 0$ and 1 and also for technical reasons of $-\frac{1}{2}$ and $\frac{1}{2}$.

First Step. Reduction to integral equations. We look for the solution $\psi_\varepsilon(q, \cdot)$ of (8.14) as a superposition of two single-layer integrals, one supported on the body's boundary and the other one supported on $\partial\Omega$. This transforms (8.14) in an integral system as follows.

We define, for any (ε, q) in \mathfrak{Q} with $q = (\vartheta, h)$ in $\mathbb{R} \times \Omega$, two operators $K^{\partial\mathcal{S}_0}(\varepsilon, q)$ and $K^{\partial\Omega}(\varepsilon, q)$ respectively from $L^2(\partial\Omega)$ to $H^1(\partial\mathcal{S}_0)$ and from $L^2(\partial\mathcal{S}_0)$ to $H^1(\partial\Omega)$, by the following formulas: given densities $\mathbf{p}^{\partial\Omega}$ and $\mathbf{p}_{\partial\mathcal{S}}$ respectively in $L^2(\partial\Omega)$ and $L^2(\partial\mathcal{S}_0)$,

$$K^{\partial\mathcal{S}_0}(\varepsilon, q)[\mathbf{p}^{\partial\Omega}](\cdot) := SL[\mathbf{p}^{\partial\Omega}](\varepsilon R(\vartheta) \cdot + h) \text{ on } \partial\mathcal{S}_0, \quad (8.33)$$

$$K^{\partial\Omega}(\varepsilon, q)[\mathbf{p}^{\partial\mathcal{S}_0}](\cdot) := \int_{\partial\mathcal{S}_0} \mathbf{p}^{\partial\mathcal{S}_0}(Y) G(\cdot - (\varepsilon R(\vartheta)Y + h)) \, ds(Y) \text{ on } \partial\Omega. \quad (8.34)$$

Thanks to Lemma 8.3 (ii), the operators $K^{\partial\mathcal{S}_0}(\varepsilon, q)$ and $K^{\partial\Omega}(\varepsilon, q)$ are compact respectively from $L^2(\partial\Omega)$ to $H^1(\partial\mathcal{S}_0)$ and from $L^2(\partial\mathcal{S}_0)$ to $H^1(\partial\Omega)$.

We also introduce for (ε, q) in \mathfrak{Q} , the operator $\mathfrak{A}(\varepsilon, q) : F_0 \rightarrow F_1$ as follows: for any $\mathbf{p} := (\mathbf{p}^{\partial\mathcal{S}_0}, \mathbf{p}^{\partial\Omega}, C)$ in F_0 ,

$$\mathfrak{A}(\varepsilon, q)[\mathbf{p}] := \left(SL[\mathbf{p}^{\partial\mathcal{S}_0}] + K^{\partial\mathcal{S}_0}(\varepsilon, q)[\mathbf{p}^{\partial\Omega}] - C, \right. \\ \left. SL[\mathbf{p}^{\partial\Omega}] + K^{\partial\Omega}(\varepsilon, q)[\mathbf{p}^{\partial\mathcal{S}_0}], \int_{\partial\mathcal{S}_0} \mathbf{p}^{\partial\mathcal{S}_0} \, ds \right). \quad (8.35)$$

In order to simplify the notations, we omitted to write the trace operators applied to the single-layers in $K^{\partial\mathcal{S}_0}(\varepsilon, q)$, $K^{\partial\Omega}(\varepsilon, q)$ and $\mathfrak{A}(\varepsilon, q)$. We also emphasize that the dependence of $\mathfrak{A}(\varepsilon, q)$ on (ε, q) occurs only through the compact operators $K^{\partial\mathcal{S}_0}(\varepsilon, q)$ and $K^{\partial\Omega}(\varepsilon, q)$.

Now the equation (8.14) is transformed into an integral system thanks to the following lemma.

Lemma 8.6. *For any (ε, q) in \mathfrak{Q} , let $\mathbf{p}_\varepsilon(q, \cdot) = (p_\varepsilon^{\partial\mathcal{S}_0}(q, \cdot), p_\varepsilon^{\partial\Omega}(q, \cdot), C_\varepsilon(q))$ in F_0 such that*

$$\mathfrak{A}(\varepsilon, q)[\mathbf{p}_\varepsilon(q, \cdot) + (0, 0, G(\varepsilon))] = (0, 0, -1). \quad (8.36)$$

Then the function in $\mathcal{F}^\varepsilon(q)$

$$\psi_\varepsilon(q, \cdot) := SL[p_\varepsilon^{\partial\mathcal{S}_\varepsilon(q)}(q, \cdot)] + SL[p_\varepsilon^{\partial\Omega}(q, \cdot)], \quad (8.37)$$

where the density $p_\varepsilon^{\partial\mathcal{S}_\varepsilon(q)}(q, \cdot)$ on $\partial\mathcal{S}_\varepsilon(q)$ is defined through the relation:

$$\text{for } X \in \partial\mathcal{S}_0, \quad p_\varepsilon^{\partial\mathcal{S}_0}(q, X) := \varepsilon p_\varepsilon^{\partial\mathcal{S}_\varepsilon(q)}(q, \varepsilon R(\vartheta)X + h), \quad (8.38)$$

is the solution of (8.14). Moreover the normal derivative $\frac{\partial\psi_\varepsilon}{\partial n}(q, \cdot)$ on $\partial\mathcal{S}_\varepsilon(q)$ is given by:

$$\text{for } X \in \partial\mathcal{S}_0, \quad \frac{\partial\psi_\varepsilon}{\partial n}(q, \varepsilon R(\vartheta)X + h) = \frac{1}{\varepsilon} p_\varepsilon^{\partial\mathcal{S}_0}(q, X). \quad (8.39)$$

The proof of Lemma 8.6 is postponed to Section 8.4.1.

Second Step. Construction of an approximate solution. In this step we describe an approximation \mathbf{p}_{app} up to order $O(\varepsilon^3)$ of the solution \mathbf{p}_ε of (8.36) and reformulate the equation (8.36) in terms of the rest $\mathbf{p}_\varepsilon - \mathbf{p}_{\text{app}}$.

We introduce the various terms involved in the approximation.

- *Densities on $\partial\mathcal{S}_0$.* Recall that the functions $\psi_{-1}^{\mathcal{S}_0}$, $\psi_0^{\mathcal{S}_0}(q, \cdot)$ and $\psi_1^{\mathcal{S}_0}(q, \cdot)$, defined respectively in (4.4), (8.17) and (8.25), are harmonic in $\mathbb{R}^2 \setminus \mathcal{S}_0$.

Let $\mathbf{p}_{-1}^{\mathcal{S}_0}$, $\mathbf{p}_0^{\mathcal{S}_0}(q, \cdot)$ and $\mathbf{p}_1^{\mathcal{S}_0}(q, \cdot)$ be the densities on $\partial\mathcal{S}_0$ associated respectively with $\psi_{-1}^{\mathcal{S}_0}$, $\psi_0^{\mathcal{S}_0}(q, \cdot)$ and $\psi_1^{\mathcal{S}_0}(q, \cdot)$ as explained in Propositions 8.1 and 8.2. Thus we have, in $\mathbb{R}^2 \setminus \mathcal{S}_0$:

$$\psi_{-1}^{\mathcal{S}_0} = SL[\mathbf{p}_{-1}^{\mathcal{S}_0}] \quad \text{and} \quad \psi_j^{\mathcal{S}_0}(q, \cdot) = SL[\mathbf{p}_j^{\mathcal{S}_0}(q, \cdot)] \quad (j = 0, 1). \quad (8.40)$$

The single layer potentials, as being supported by $\partial\mathcal{S}_0$, are actually defined in \mathbb{R}^2 and harmonic in $\mathbb{R}^2 \setminus \partial\mathcal{S}_0$. The identities above can therefore be used to extend the functions $\psi_{-1}^{\mathcal{S}_0}$, $\psi_0^{\mathcal{S}_0}(q, \cdot)$ and $\psi_1^{\mathcal{S}_0}(q, \cdot)$ in \mathcal{S}_0 . Recalling the definitions (8.15) and (8.22) of the harmonic polynomials P_0 and P_1 , this leads explicitly to:

- $\psi_{-1}^{\mathcal{S}_0}$ is extended by $C^{\mathcal{S}_0}$ in \mathcal{S}_0 ,
- $\psi_0^{\mathcal{S}_0}(q, \cdot)$ is extended by $P_0(q, \cdot)$ in \mathcal{S}_0 and
- $\psi_1^{\mathcal{S}_0}(q, \cdot)$ is extended by $P_1(q, \cdot)$ in \mathcal{S}_0 .

- *Densities on $\partial\Omega$.* We follow the same ideas to extend in \mathbb{R}^2 the functions $\psi_0^\Omega(h, \cdot)$, $\psi_1^\Omega(q, \cdot)$ and $\psi_2^\Omega(q, \cdot)$, so far defined and harmonic in Ω (see (2.39), (8.18) and (8.44)). We consider their densities $\mathbf{p}_0^\Omega(q, \cdot)$, $\mathbf{p}_1^\Omega(q, \cdot)$ and $\mathbf{p}_2^\Omega(q, \cdot)$, supported in $\partial\Omega$ and such that, in Ω :

$$\psi_0^\Omega(h, \cdot) = SL[\mathbf{p}_0^\Omega(h, \cdot)], \quad \psi_j^\Omega(q, \cdot) = SL[\mathbf{p}_j^\Omega(q, \cdot)] \quad (j = 1, 2). \quad (8.41)$$

The single layer potentials being defined in \mathbb{R}^2 and harmonic in $\mathbb{R}^2 \setminus \partial\Omega$, these identities are used to extend the functions $\psi_0^\Omega(h, \cdot)$, $\psi_1^\Omega(q, \cdot)$ and $\psi_2^\Omega(q, \cdot)$ in \mathbb{R}^2 . Explicitly, this leads to:

- $\psi_0^\Omega(h, \cdot)$ is extended by $G(\cdot - h)$ in $\mathbb{R}^2 \setminus \Omega$,
- $\psi_1^\Omega(q, \cdot)$ is extended by $-(\nabla G)(\cdot - h) \cdot \zeta_\vartheta$ in $\mathbb{R}^2 \setminus \Omega$, and
- $\psi_2^\Omega(q, \cdot)$ is extended by $Q_2(q, \cdot)$, defined in (8.42), in $\mathbb{R}^2 \setminus \Omega$.

- *Definition of $\psi_2^\Omega(q, \cdot)$ and of $Q_2(q, \cdot)$.* In order to define $\psi_2^\Omega(q, \cdot)$, we introduce the harmonic function in $\mathbb{R}^2 \setminus \Omega$:

$$Q_2(q, x) := \frac{1}{2} \langle R(\vartheta)^t D^2 G(x - h) R(\vartheta), T^2(\mathbf{p}_{-1}^{S_0}) \rangle_{\mathbb{R}^{2 \times 2}} + R(\vartheta)^t \nabla G(x - h) \cdot T^1(\mathbf{p}_0^{S_0}(q, \cdot)), \quad (8.42)$$

where

$$T^1(\mathbf{p}_0^{S_0}(q, \cdot)) := \int_{\partial S_0} Y \mathbf{p}_0^{S_0}(q, Y) ds(Y). \quad (8.43)$$

Then we consider $\psi_2^\Omega(q, \cdot)$ as the solution of

$$-\Delta \psi_2^\Omega(q, \cdot) = 0 \quad \text{in } \Omega, \quad (8.44a)$$

$$\psi_2^\Omega(q, \cdot) = Q_2(q, \cdot) \quad \text{on } \partial\Omega, \quad (8.44b)$$

extended by $Q_2(q, x)$ for x in $\mathbb{R}^2 \setminus \Omega$. These functions $\psi_2^\Omega(q, \cdot)$ and $Q_2(q, \cdot)$ do not appear in the claim of Proposition 8.4 and of Lemma 6.4 but will be useful later.

With the choices above we expect to construct a solution of (8.36) with $p_\varepsilon^{\partial S_0}(\varepsilon, q, \cdot)$ and $p_\varepsilon^{\partial\Omega}(\varepsilon, q, \cdot)$ close respectively to

$$\mathbf{p}_{\text{app}}^{\partial S_0}(\varepsilon, q, \cdot) := \mathbf{p}_{-1}^{S_0} + \varepsilon \mathbf{p}_0^{S_0}(q, \cdot) + \varepsilon^2 \mathbf{p}_1^{S_0}(q, \cdot), \quad (8.45)$$

$$\mathbf{p}_{\text{app}}^{\partial\Omega}(\varepsilon, q, \cdot) := \mathbf{p}_0^\Omega(h, \cdot) + \varepsilon \mathbf{p}_1^\Omega(q, \cdot) + \varepsilon^2 \mathbf{p}_2^\Omega(q, \cdot).$$

The corresponding approximation $C_{\text{app}}(\varepsilon, q)$ of $C_\varepsilon(q)$ is chosen as:

$$C_{\text{app}}(\varepsilon, q) := -G(\varepsilon) + C_0(h) + \varepsilon C_1(q) + \varepsilon^2 C_2(q), \quad (8.46)$$

with

$$\begin{aligned}
C_0(h) &:= C^{\mathcal{S}_0} + \psi_0^\Omega(h, h), \\
C_I(q) &:= 2D_x \psi_0^\Omega(h, h) \cdot \zeta_\vartheta, \\
C_2(q) &:= \psi_2^\Omega(q, h) + D_x \psi_1^\Omega(q, h) \cdot \zeta_\vartheta \\
&\quad + \frac{1}{2} \langle R(\vartheta)^t D_x^2 \psi_0^\Omega(h, h) R(\vartheta), T^2(\mathbf{p}_{-I}^{\mathcal{S}_0}) \rangle_{\mathbb{R}^2 \times 2}. \tag{8.47}
\end{aligned}$$

Using that ψ_0^Ω is symmetric with respect to its two arguments (see (8.21)), and using (2.41), we see that the first terms of the expansion above are the same as those claimed in Lemma 6.4, that is

$$C_0(h) := C^{\mathcal{S}_0} + 2\psi^\Omega(h) \quad \text{and} \quad C_I(q) := 2\psi_c(q). \tag{8.48}$$

We finally denote

$$\mathbf{p}_{\text{app}}(\varepsilon, q, \cdot) := (\mathbf{p}_{\text{app}}^{\partial \mathcal{S}_0}(\varepsilon, q, \cdot), \mathbf{p}_{\text{app}}^{\partial \Omega}(\varepsilon, q, \cdot), C_{\text{app}}(\varepsilon, q)). \tag{8.49}$$

Now the equation (8.36) translates as follows. Let us introduce $g^{\partial \mathcal{S}_0}(\varepsilon, q, \cdot)$ and $g^{\partial \Omega}(\varepsilon, q, \cdot)$ two functions respectively defined on $\partial \mathcal{S}_0$ and $\partial \Omega$, for $q = (\vartheta, h)$, by

$$-g^{\partial \mathcal{S}_0}(\varepsilon, q, \cdot) := \sum_{j=0}^2 \int_{\partial \Omega} \mathbf{p}_j^\Omega(q, y) \eta_{\mathcal{S}-j}(\varepsilon, q, \cdot, y) \, ds(y), \tag{8.50a}$$

$$\begin{aligned}
-g^{\partial \Omega}(\varepsilon, q, x) &:= \int_{\partial \mathcal{S}_0} \mathbf{p}_{-1}^{\mathcal{S}_0}(y) \eta_{\mathcal{S}}(\varepsilon, (\vartheta, x), -y, h) \, ds(y) \\
&\quad + \sum_{j=0}^1 \int_{\partial \mathcal{S}_0} \mathbf{p}_j^{\mathcal{S}_0}(q, y) \eta_{2-j}(\varepsilon, (\vartheta, x), -y, h) \, ds(y), \tag{8.50b}
\end{aligned}$$

where we have denoted, for $N \geq 1$,

$$\eta_N(\varepsilon, q, \cdot, y) := \int_0^1 \frac{(1-\sigma)^{N-1}}{(N-1)!} D^N G(\sigma \varepsilon R(\vartheta) \cdot + h - y) \cdot (R(\vartheta) \cdot)^{\otimes N} \, d\sigma. \tag{8.51}$$

Let

$$\mathbf{g}(\varepsilon, q, \cdot) := (g^{\partial \mathcal{S}_0}(\varepsilon, q, \cdot), g^{\partial \Omega}(\varepsilon, q, \cdot), 0). \tag{8.52}$$

We can deduce from the definitions of the densities \mathbf{p}_j^Ω for $j = 0, 1, 2$, $\mathbf{p}_{-1}^{\mathcal{S}_0}$ and $\mathbf{p}_j^{\mathcal{S}_0}$ for $j = 0$ and 1 , and from Lemma 8.3, (ii) that $g^{\partial \mathcal{S}_0}(\varepsilon, q, \cdot)$ and $g^{\partial \Omega}(\varepsilon, q, \cdot)$ belong respectively to $H^1(\partial \mathcal{S}_0)$ and to $H^1(\partial \Omega)$. Actually we even have

$$\mathbf{g} \in L^\infty(\Omega^\delta; F_1). \tag{8.53}$$

We can now state the result of this second step.

Lemma 8.7. *For any (ε, q) in \mathfrak{Q} , let $\mathbf{p}_r(\varepsilon, q, \cdot)$ in F_0 satisfy:*

$$\mathfrak{A}(\varepsilon, q)[\mathbf{p}_r(\varepsilon, q, \cdot)] = \mathfrak{g}(\varepsilon, q, \cdot). \quad (8.54)$$

Then

$$\mathbf{p}_\varepsilon(q, \cdot) := \mathbf{p}_{app}(\varepsilon, q, \cdot) + \varepsilon^3 \mathbf{p}_r(\varepsilon, q, \cdot), \quad (8.55)$$

is solution of (8.36).

The proof of Lemma 8.7 is postponed to Section 8.4.2. Let us stress in particular that the third argument of the left hand side of (8.54) does not contain the singular term $G(\varepsilon)$ anymore (compare with (8.36)) and that the third component of the right hand side of (8.54) is now 0.

Third Step. Existence and estimate of the remainders. In this third step we prove, for (ε, q) in $\mathfrak{Q}_{\delta, \varepsilon_0}$ with δ and ε_0 positive and small enough, the existence of $\mathbf{p}_r(\varepsilon, q, \cdot)$ in F_0 satisfying (8.54) and provide an estimate in F_0 , uniform over (ε, q) in $\mathfrak{Q}_{\delta, \varepsilon_0}$.

We will make use of the fact that the the third argument of the right hand side of (8.54) vanishes. Accordingly, we denote

$$\tilde{F}_1 := H^1(\partial\mathcal{S}_0) \times H^1(\partial\Omega) \times \{0\}, \quad (8.56)$$

which is a closed subspace of F_1 and prove the following result.

Lemma 8.8. *Let $\delta > 0$. There exists ε_0 in $(0, 1)$, such that for any \mathfrak{g} in $L^\infty(\mathfrak{Q}_{\delta, \varepsilon_0}; \tilde{F}_1)$, there exists \mathbf{p}_r in $L^\infty(\mathfrak{Q}_{\delta, \varepsilon_0}; F_0)$ such that $\mathbf{p}_r(\varepsilon, q, \cdot)$ solves (8.54) for any (ε, q) in $\mathfrak{Q}_{\delta, \varepsilon_0}$.*

Proof of Lemma 8.8. In order to prove Lemma 8.8 let us start with stating a perturbative result. We will use the notation that given X and Y two Banach spaces, $\mathcal{L}(X; Y)$ is the space of bounded linear operators from X to Y . Now the framework is as follows. Let $\delta > 0$. Recall that \mathfrak{Q}^δ and Ω_δ were defined in (6.3) and (6.5). We introduce the following families of operators.

- First we consider a family of operators in $\mathcal{L}(L^2(\partial\Omega); H^1(\partial\mathcal{S}_0))$:

$$\begin{aligned} \tilde{K}^{\partial\mathcal{S}_0} \in \text{Lip}(\overline{\Omega_\delta}; \mathcal{L}(L^2(\partial\Omega); H^1(\partial\mathcal{S}_0))) \text{ such that for all } h \text{ in } \overline{\Omega_\delta}, \\ \tilde{K}^{\partial\mathcal{S}_0}(h) \text{ is compact from } L^2(\partial\Omega) \text{ to } H^1(\partial\mathcal{S}_0). \end{aligned} \quad (8.57)$$

- Next we consider two families of operators: one in $\mathcal{L}(L^2(\partial\Omega); H^1(\partial\mathcal{S}_0))$ and the other one in $\mathcal{L}(L^2(\partial\mathcal{S}_0); H^1(\partial\Omega))$:

$$(T^{\partial\mathcal{S}_0}(\varepsilon, q))_{(\varepsilon, q) \in \mathfrak{Q}^\delta} \text{ bounded in } \mathcal{L}(L^2(\partial\Omega); H^1(\partial\mathcal{S}_0)), \quad (8.58a)$$

$$(T^{\partial\Omega}(\varepsilon, q))_{(\varepsilon, q) \in \mathfrak{Q}^\delta} \text{ bounded in } \mathcal{L}(L^2(\partial\mathcal{S}_0); H^1(\partial\Omega)). \quad (8.58b)$$

Given these operators we can construct the following one. For (ε, q) in \mathfrak{Q}^δ , let $A(\varepsilon, q) : F_0 \rightarrow F_1$ given by the following formula: for any $\mathbf{p} := (\mathbf{p}^{\partial\mathcal{S}_0}, \mathbf{p}^{\partial\Omega}, C)$ in F_0 ,

$$A(\varepsilon, q)[\mathbf{p}] := (A(\varepsilon, q)[\mathbf{p}]_i)_{1 \leq i \leq 3} \in F_1, \quad (8.59)$$

with

$$A(\varepsilon, q)[\mathbf{p}]_1 := SL[\mathbf{p}^{\partial\mathcal{S}_0}] + \tilde{K}^{\partial\mathcal{S}_0}(h)[\mathbf{p}^{\partial\Omega}] + \varepsilon T^{\partial\mathcal{S}_0}(\varepsilon, q)[\mathbf{p}^{\partial\Omega}] - C, \quad (8.60a)$$

$$A(\varepsilon, q)[\mathbf{p}]_2 := SL[\mathbf{p}^{\partial\Omega}] + \varepsilon T^{\partial\Omega}(\varepsilon, q)[\mathbf{p}^{\partial\mathcal{S}_0}], \quad (8.60b)$$

$$A(\varepsilon, q)[\mathbf{p}]_3 := \int_{\partial\mathcal{S}_0} \mathbf{p}^{\partial\mathcal{S}_0} ds. \quad (8.60c)$$

Our perturbative result is the following lemma whose proof is postponed to Section 8.4.3.

Lemma 8.9. *Let $\delta > 0$ and for (ε, q) in \mathfrak{Q}^δ , $A(\varepsilon, q)$ given as above, with assumptions (8.57) and (8.58). Then there exists ε_0 in $(0, 1)$ such that for any (ε, q) in $\mathfrak{Q}_{\delta, \varepsilon_0}$, $A(\varepsilon, q)$ is an isomorphism from F_0 to F_1 and*

$$\sup_{(\varepsilon, q) \in \mathfrak{Q}_{\delta, \varepsilon_0}} \|A(\varepsilon, q)^{-1}\|_{\mathcal{L}(F_1; F_0)} < \infty. \quad (8.61)$$

In our case, Lemma 8.9 is applied as follows. Recalling (8.33)-(8.34) we define, for any (ε, q) in \mathfrak{Q} , with $q = (\vartheta, h)$,

– for any density $\mathbf{p}^{\partial\Omega}$ in $L^2(\partial\Omega)$,

$$\tilde{K}^{\partial\mathcal{S}_0}(h)[\mathbf{p}^{\partial\Omega}] = K^{\partial\mathcal{S}_0}(0, 0, h)[\mathbf{p}^{\partial\Omega}] = SL[\mathbf{p}^{\partial\Omega}](h) \quad (8.62)$$

as a constant function on $\partial\mathcal{S}_0$, and

$$T^{\partial\mathcal{S}_0}(\varepsilon, q)[\mathbf{p}^{\partial\Omega}] := \int_{\partial\Omega} \mathbf{p}^{\partial\Omega}(y) \eta_I(\varepsilon, q, \cdot, y) ds(y) \quad \text{on } \partial\mathcal{S}_0, \quad (8.63)$$

– for any density $\mathbf{p}^{\partial\mathcal{S}_0}$ in $L^2(\partial\mathcal{S}_0)$,

$$T^{\partial\Omega}(\varepsilon, q)[\mathbf{p}^{\partial\mathcal{S}_0}] := \int_{\partial\mathcal{S}_0} \mathbf{p}^{\partial\mathcal{S}_0}(Y) \eta_I(\varepsilon, \vartheta, \cdot, -Y, h) ds(y) \quad \text{on } \partial\Omega. \quad (8.64)$$

The following lemma whose proof is postponed to Section 8.4.4 entails that the hypotheses of Lemma 8.9 are satisfied.

Lemma 8.10. *Let $\delta > 0$. With the definitions above, (8.57) and (8.58) hold true.*

Then we consider the operator $A(\varepsilon, q)$ associated with these operators $\tilde{K}^{\partial\mathcal{S}_0}(h)$, $T^{\partial\mathcal{S}_0}(\varepsilon, q)$ and $T^{\partial\Omega}(\varepsilon, q)$ as given by (8.59)-(8.60). The next lemma shows that this operator $A(\varepsilon, q)$ provides the existence of a solution to (8.54) with uniform estimates.

Lemma 8.11. *Let $\delta > 0$. There exists ε_0 in $(0, 1)$ such that for any (ε, q) in $\mathfrak{Q}_{\delta, \varepsilon_0}$,*

$$\mathbf{p}_r(\varepsilon, q, \cdot) := A(\varepsilon, q)^{-1} \mathbf{g}(\varepsilon, q, \cdot) \quad (8.65)$$

belongs to F_0 and solves (8.54). Moreover \mathbf{p}_r is in $L^\infty(\mathfrak{Q}_{\delta, \varepsilon_0}; F_0)$.

The proof of Lemma 8.11 is postponed to Section 8.4.5. Now once assumed Lemma 8.11, Lemma 8.8 follows in a straightforward manner. \square

Fourth Step. Conclusion.

End of proof of Lemma 6.4 We apply Lemma 8.8 to (8.52). Thanks to (8.53) the assumption is satisfied. Regarding $C_\varepsilon(q)$ this yields an expansion actually better than the one stated in Lemma 6.4, that is, according to (8.46) and (8.48) and what precedes, there exists C_r in $L^\infty(\mathfrak{Q}_{\delta, \varepsilon_0}; \mathbb{R})$ such that

$$\begin{aligned} C_\varepsilon(q) = & -G(\varepsilon) + C^{S_0} + \psi_0^\Omega(h, h) + 2\varepsilon D_x \psi_0^\Omega(h, h) \cdot \zeta_\vartheta + \varepsilon^2 C_2(q) \\ & + \varepsilon^3 C_r(\varepsilon, q), \end{aligned} \quad (8.66)$$

where $C_2(q)$ is given by (8.47). In order to prove Lemma 6.4 it is therefore sufficient to observe that $C_2(q)$ is bounded uniformly in \mathbb{R} for (ε, q) in \mathfrak{Q}^δ and to redefine $C_r(\varepsilon, q)$ such that $\varepsilon^2 C_r(\varepsilon, q)$ is equal to the sum of the last two terms in (8.66). \square

End of proof of Proposition 8.4 Combining (8.39), (8.45) and (8.55), we get that on $\partial\mathcal{S}_0$

$$\frac{\partial \psi_\varepsilon}{\partial n}(q, \varepsilon R(\vartheta) \cdot + h) = \mathbf{p}_{-1}^{S_0}(\cdot) + \varepsilon \mathbf{p}_0^{S_0}(q, \cdot) + \varepsilon^2 \mathbf{p}_1^{S_0}(q, \cdot) + \varepsilon^3 \mathbf{p}_r^{\partial S_0}(\varepsilon, q, \cdot),$$

with $\mathbf{p}_r^{\partial S_0}$ in $L^\infty(\mathfrak{Q}_{\delta, \varepsilon_0}; L^2(\partial\mathcal{S}_0; \mathbb{R}))$. Moreover using (8.2) we have that

$$\mathbf{p}_{-1}^{S_0} = \frac{\partial \psi_{-1}^{S_0}}{\partial n}, \quad \mathbf{p}_0^{S_0} = \frac{\partial \psi_0^{S_0}}{\partial n} - \frac{\partial P_0}{\partial n} \quad \text{and} \quad \mathbf{p}_1^{S_0} = \frac{\partial \psi_1^{S_0}}{\partial n} - \frac{\partial P_1}{\partial n}. \quad (8.67)$$

Referring to the definition of P_0 in (8.15) we obtain $\frac{\partial P_0}{\partial n}(q, X) = -R(\vartheta)^t u^\Omega(h) \cdot \tau$, for X on $\partial\mathcal{S}_0$, which concludes the proof of Proposition 8.4. \square

8.4 Proof of the intermediate lemmas

In this subsection, we establish the intermediate lemmas used in Subsection 8.3.

8.4.1 Proof of Lemma 8.6

First observe that for any densities $p_\varepsilon^{\partial\mathcal{S}_\varepsilon(q)}(q, \cdot)$ in $H^{-\frac{1}{2}}(\partial\mathcal{S}_\varepsilon(q))$ and $p_\varepsilon^{\partial\Omega}(q, \cdot)$ in $H^{-\frac{1}{2}}(\partial\Omega)$, the right hand side of (8.37) is in $H_{loc}^1(\mathbb{R}^2)$ and harmonic in $\mathcal{F}_\varepsilon(q)$ and in $\mathbb{R}^2 \setminus \mathcal{F}_\varepsilon(q)$. In particular the equation (8.14a) is satisfied when $\psi^\varepsilon(q, \cdot)$ is given by (8.37) without further assumptions about $p_\varepsilon^{\partial\mathcal{S}_\varepsilon(q)}(q, \cdot)$ or $p_\varepsilon^{\partial\Omega}(q, \cdot)$.

Next we write (8.36) explicitly in the form:

$$-G(\varepsilon) + SL[p_\varepsilon^{\partial\mathcal{S}_0}(q, \cdot)] + K^{\partial\mathcal{S}_0}(\varepsilon, q)[p_\varepsilon^{\partial\Omega}(q, \cdot)] = C_\varepsilon(q) \quad \text{on } \partial\mathcal{S}_0, \quad (8.68a)$$

$$SL[p_\varepsilon^{\partial\Omega}(q, \cdot)] + K^{\partial\Omega}(\varepsilon, q)[p_\varepsilon^{\partial\mathcal{S}_0}(q, \cdot)] = 0 \quad \text{on } \partial\Omega, \quad (8.68b)$$

$$\int_{\partial\mathcal{S}_0} p_\varepsilon^{\partial\mathcal{S}_0}(q, \cdot) \, ds = -1. \quad (8.68c)$$

Thanks to a change of variable, and using $G(\varepsilon(x-y)) = G(\varepsilon) + G(x-y)$, (8.33), (8.34) and (8.38), this can be recast as

$$SL[p_\varepsilon^{\partial\mathcal{S}_\varepsilon(q)}(q, \cdot)] + SL[p_\varepsilon^{\partial\Omega}(q, \cdot)] = C_\varepsilon(q) \quad \text{on } \partial\mathcal{S}_\varepsilon(q), \quad (8.69a)$$

$$SL[p_\varepsilon^{\partial\mathcal{S}_\varepsilon(q)}(q, \cdot)] + SL[p_\varepsilon^{\partial\Omega}(q, \cdot)] = 0 \quad \text{on } \partial\Omega, \quad (8.69b)$$

$$\int_{\partial\mathcal{S}_\varepsilon(q)} p_\varepsilon^{\partial\mathcal{S}_\varepsilon(q)}(q, \cdot) \, ds = -1. \quad (8.69c)$$

In particular we infer from (8.69a) and (8.69b) that, when $\psi_\varepsilon(q, \cdot)$ is given by (8.37) with $\mathbf{p}_\varepsilon(q, \cdot) = (p_\varepsilon^{\partial\mathcal{S}_0}(q, \cdot), p_\varepsilon^{\partial\Omega}(q, \cdot), C_\varepsilon(q))$ solution of (8.36), the boundary conditions (8.14b) and (8.14c) are satisfied. Moreover, by uniqueness of the solutions to the Poisson problem:

$$\Delta\Psi = 0 \text{ in } \mathcal{S}_\varepsilon(q), \quad \Psi = C_\varepsilon(q) \text{ on } \partial\mathcal{S}_\varepsilon(q),$$

the right hand side of (8.37) is equal to $C_\varepsilon(q)$ in $\mathcal{S}_\varepsilon(q)$.

The single-layer potential $SL[p_\varepsilon^{\partial\Omega}(q, \cdot)]$ is smooth in a neighborhood of $\partial\mathcal{S}_\varepsilon(q)$. Hence, according to (8.2), when $\psi^\varepsilon(q, \cdot)$ is given by (8.37), the density $p_\varepsilon^{\partial\mathcal{S}_\varepsilon(q)}(q, \cdot)$ is equal to the jump across $\partial\mathcal{S}_\varepsilon(q)$ of the normal derivatives of the function equal to $\psi_\varepsilon(q, \cdot)$ in $\mathcal{F}_\varepsilon(q)$ and to $C_\varepsilon(q)$ in $\mathcal{S}_\varepsilon(q)$, that is

$$p_\varepsilon^{\partial\mathcal{S}_\varepsilon(q)}(q, \cdot) = \frac{\partial\psi_\varepsilon}{\partial n}(q, \cdot) \text{ on } \partial\mathcal{S}_\varepsilon(q).$$

Hence we obtain (8.39) by using (8.38) and the condition (8.14d) by using (8.69c). This concludes the proof of Lemma 8.6.

8.4.2 Proof of Lemma 8.7

Let (ε, q) in \mathfrak{Q} and $\mathbf{p}_r(\varepsilon, q, \cdot) := (\mathbf{p}_r^{\partial\mathcal{S}_0}(\varepsilon, q, \cdot), p_r^{\partial\Omega}(\varepsilon, q, \cdot), C_r(\varepsilon, q))$ in F_0 satisfying (8.54) (where \mathfrak{A} is defined in (8.35)), that is

$$SL[\mathbf{p}_r^{\partial\mathcal{S}_0}(\varepsilon, q, \cdot)] + K^{\partial\mathcal{S}_0}(\varepsilon, q)[\mathbf{p}_r^{\partial\Omega}(\varepsilon, q, \cdot)] - C_r(\varepsilon, q) = g^{\partial\mathcal{S}_0}(\varepsilon, q, \cdot) \text{ on } \partial\mathcal{S}_0, \quad (8.70)$$

$$SL[\mathbf{p}_r^{\partial\Omega}(\varepsilon, q, \cdot)] + K^{\partial\Omega}(\varepsilon, q)[\mathbf{p}_r^{\partial\mathcal{S}_0}(\varepsilon, q, \cdot)] = g^{\partial\Omega}(\varepsilon, q, \cdot) \text{ on } \partial\Omega, \quad (8.71)$$

$$\int_{\partial\mathcal{S}_0} \mathbf{p}_r^{\partial\mathcal{S}_0}(\varepsilon, q, \cdot) \, ds = 0. \quad (8.72)$$

Let $\mathbf{p}_\varepsilon(q, \cdot) = (p_\varepsilon^{\partial\mathcal{S}_0}(q, \cdot), p_\varepsilon^{\partial\Omega}(q, \cdot), C_\varepsilon(q))$ in F_0 given by (8.55). In order to prove (8.36) we now verify the three parts of (8.68).

Verification of (8.68c). Using (8.67), the fact that P_0 and P_1 are harmonic and the conditions (4.4d), (8.17c) and (8.25c) we get that

$$\int_{\partial\mathcal{S}_0} \mathbf{p}_{-1}^{\mathcal{S}_0} \, ds = -1 \quad \text{and} \quad \int_{\partial\mathcal{S}_0} \mathbf{p}_0^{\mathcal{S}_0}(q, \cdot) \, ds = \int_{\partial\mathcal{S}_0} \mathbf{p}_1^{\mathcal{S}_0}(q, \cdot) \, ds = 0. \quad (8.73)$$

As a consequence of (8.73) and of (8.72) we get that the condition (8.68c) is fulfilled.

Verification of (8.68a). Using (8.49), (8.55), (8.6), (4.4b), (8.40), we obtain, on $\partial\mathcal{S}_0$,

$$SL[p_\varepsilon^{\partial\mathcal{S}_0}(q, \cdot)] = C^{\mathcal{S}_0} + \varepsilon\psi_0^{\mathcal{S}_0}(q, \cdot) + \varepsilon^2\psi_1^{\mathcal{S}_0}(q, \cdot) + \varepsilon^3 SL[p_{\varepsilon,r}^{\partial\mathcal{S}_0}(\varepsilon, q, \cdot)]. \quad (8.74)$$

Using (8.33), (8.49), (8.55), (8.41) and then Taylor's formula and (8.19) we get on $\partial\mathcal{S}_0$

$$\begin{aligned} K^{\partial\mathcal{S}_0}(\varepsilon, q)[p_\varepsilon^{\partial\Omega}(q, \cdot)](X) &= \psi_0^\Omega(h, h) + \varepsilon D_x \psi_0^\Omega(h, h) \cdot (R(\vartheta)X + \zeta_\vartheta) \\ &+ \varepsilon^2 \left(\psi_2^\Omega(q, h) + D_x \psi_1^\Omega(q, h) \cdot R(\vartheta)X + \frac{1}{2} D_x^2 \psi_0^\Omega(h, h) \cdot (R(\vartheta)X, R(\vartheta)X) \right) \\ &+ \varepsilon^3 \left(K^{\partial\mathcal{S}_0}(\varepsilon, q)[\mathbf{p}_r^{\partial\Omega}(\varepsilon, q, \cdot)](X) - g^{\partial\mathcal{S}_0}(\varepsilon, q, X) \right). \end{aligned} \quad (8.75)$$

We recall that the function $g^{\partial\mathcal{S}_0}(\varepsilon, q, \cdot)$ is defined in (8.50a).

Gathering (8.46), (8.55), (8.74) and (8.75) we obtain, on $\partial\mathcal{S}_0$,

$$\begin{aligned} &-G(\varepsilon) + SL[p_\varepsilon^{\partial\mathcal{S}_0}(q, \cdot)] + K^{\partial\mathcal{S}_0}(\varepsilon, q)[p_\varepsilon^{\partial\Omega}(q, \cdot)] - C_\varepsilon(q) \\ &= \varepsilon(\psi_0^{\mathcal{S}_0}(q, \cdot) - P_0(q, \cdot)) + \varepsilon^2(\psi_1^{\mathcal{S}_0}(q, \cdot) - P_1(q, \cdot)) \\ &+ \varepsilon^3 \left(SL[p_{\varepsilon,r}^{\partial\mathcal{S}_0}(\varepsilon, q, \cdot)] + K^{\partial\mathcal{S}_0}(\varepsilon, q)[\mathbf{p}_r^{\partial\Omega}(\varepsilon, q, \cdot)] - C_r(\varepsilon, q) - g^{\partial\mathcal{S}_0}(\varepsilon, q, \cdot) \right), \end{aligned}$$

where $P_0(q, \cdot)$ and $P_1(q, \cdot)$ are the harmonic polynomials defined respectively in (8.15) and in (8.22). Now taking into account the boundary conditions (8.17b) and (8.44b), and (8.70) we get that (8.68a) holds true.

Remark 8.12. Let us explain a bit how the ansatz of $C_\varepsilon(q)$ given by (8.46) and (8.55) was guessed. Taking into account (8.74) and (8.75) we multiply $-G(\varepsilon) + SL[p_\varepsilon^{\partial S_0}(q, \cdot)] + K^{\partial S_0}(\varepsilon, q)[p_\varepsilon^{\partial \Omega}(q, \cdot)]$ by $-\mathbf{p}_{-1}^{S_0}$, and integrate over ∂S_0 , combining with (8.68a). We then simplify the resulting equation with the following observations:

- using (4.14), (4.12) and the first equality in (8.67), we have:

$$\int_{\partial S_0} \mathbf{p}_{-1}^{S_0}(X) R(\vartheta) X \, ds(X) = -\zeta_\vartheta, \quad (8.76)$$

so that

$$\begin{aligned} \int_{\partial S_0} \mathbf{p}_{-1}^{S_0}(X) D_x \psi_0^\Omega(h, h) \cdot R(\vartheta) X \, ds(X) &= -D_x \psi_0^\Omega(h, h) \cdot \zeta_\vartheta, \\ \int_{\partial S_0} \mathbf{p}_{-1}^{S_0}(X) D_x \psi_1^\Omega(q, h) \cdot R(\vartheta) X \, ds(X) &= -D_x \psi_1^\Omega(q, h) \cdot \zeta_\vartheta; \end{aligned}$$

- using again the first equality in (8.67) and (8.23), we get:

$$\int_{\partial S_0} \mathbf{p}_{-1}^{S_0}(X) X^{\otimes 2} \, ds(X) = T^2(\mathbf{p}_{-1}^{S_0}), \quad (8.77)$$

so that

$$\begin{aligned} \int_{\partial S_0} \mathbf{p}_{-1}^{S_0}(X) D_x^2 \psi_0^\Omega(h, h) \cdot (R(\vartheta) X, R(\vartheta) X) \, ds(X) &= \\ \langle R(\vartheta)^t D_x^2 \psi_0^\Omega(h, h) R(\vartheta), T^2(\mathbf{p}_{-1}^{S_0}) \rangle_{\mathbb{R}^{2 \times 2}}; \end{aligned}$$

- using (8.16), (8.17b), (8.24), (8.25b).

Then we deduce that if (8.68c) and (8.68a) hold true then $C_\varepsilon(q)$ should be given by (8.46) and (8.55) with

$$\begin{aligned} C_r(\varepsilon, q) &:= \\ &- \int_{\partial S_0} \mathbf{p}_{-1}^{S_0} (SL[p_\varepsilon^{\partial S, r}(\varepsilon, q, \cdot)] + K^{\partial S_0}(\varepsilon, q)[\mathbf{p}_r^{\partial \Omega}(\varepsilon, q, \cdot)] - g^{\partial S_0}(\varepsilon, q, \cdot)) \, ds. \end{aligned}$$

We see that the ansatz (8.46)-(8.55) leads to a remainder $C_r(\varepsilon, q)$ of order $O(1)$, which encourages the try of (8.46) as an approximate solution.

Verification of (8.68b). First, using (8.49), (8.55) and (8.41), we obtain, on $\partial \Omega$,

$$SL[p_\varepsilon^{\partial \Omega}(q, \cdot)] = \psi_0^\Omega(h, \cdot) + \varepsilon \psi_1^\Omega(q, \cdot) + \varepsilon^2 \psi_2^\Omega(q, \cdot) + \varepsilon^3 SL[\mathbf{p}_r^{\partial \Omega}(\varepsilon, q, \cdot)]. \quad (8.78)$$

On the other hand, using (8.34), (8.55), Taylor's formula, (8.50b), (8.43), (8.72), (8.73), (8.76), (8.77), we get, for x in $\partial\Omega$,

$$\begin{aligned} K^{\partial\Omega}(\varepsilon, q)[p_\varepsilon^{\partial S_0}(q, \cdot)](x) &= -G(x-h) + \varepsilon DG(x-h) \cdot \zeta_\vartheta \\ &+ \varepsilon^2 \left(-R(\vartheta)^t DG(x-h) \cdot T^1(\mathbf{p}_0^{S_0}(q, \cdot)) \right. \\ &\quad \left. - \frac{1}{2} \langle R(\vartheta)^t D_x^2 G(x-h) R(\vartheta), T^2(\mathbf{p}_{-1}^{S_0}) \rangle_{\mathbb{R}^2 \times \mathbb{R}^2} \right) \\ &+ \varepsilon^3 \left(K^{\partial\Omega}(\varepsilon, q)[p_r^{\partial S_0}(\varepsilon, q, \cdot)](x) - g^{\partial\Omega}(\varepsilon, q, x) \right). \end{aligned} \quad (8.79)$$

Gathering (8.78) and (8.79), the equation (8.68b) now reads, for x in $\partial\Omega$,

$$\begin{aligned} SL[p_\varepsilon^{\partial\Omega}(q, \cdot)](x) + K^{\partial\Omega}(\varepsilon, q)[p_\varepsilon^{\partial S_0}(q, \cdot)](x) &= \psi_0^\Omega(h, x) - G(x-h) \\ &+ \varepsilon \left(\psi_1^\Omega(q, x) + DG(x-h) \cdot \zeta_\vartheta \right) + \varepsilon^2 \left(\psi_2^\Omega(q, x) - Q_2(q, x) \right) \\ &+ \varepsilon^3 \left(SL[p_r^{\partial\Omega}(\varepsilon, q, \cdot)] + K^{\partial\Omega}(\varepsilon, q)[p_{\varepsilon, r}^{\partial S_0}(\varepsilon, q, \cdot)](x) - g^{\partial\Omega}(\varepsilon, q, x) \right), \end{aligned} \quad (8.80)$$

where $Q_2(q, x)$ denotes the harmonic polynomial defined in (8.42).

Taking now the boundary conditions (2.39), (8.18b), (8.44b) and (8.71) into account, we deduce from the equation (8.80) that (8.68b) holds true. This concludes the proof of Lemma 8.7.

8.4.3 Proof of Lemma 8.9

It is straightforward to see that for any (ε, q) in \mathfrak{Q} , $A(\varepsilon, q)$ is linear continuous. Let (ε, q) in \mathfrak{Q} , with $q = (\vartheta, h)$ in $\mathbb{R} \times \Omega$. Let us introduce, for any $\mathbf{p} := (\mathbf{p}^{\partial S_0}, \mathbf{p}^{\partial\Omega}, C)$ in $F_{-\frac{1}{2}}$,

$$\begin{aligned} L[\mathbf{p}] &:= (SL[\mathbf{p}^{\partial S_0}], SL[\mathbf{p}^{\partial\Omega}], C), \\ K(h)[\mathbf{p}] &:= (\tilde{K}^{\partial S_0}(h)[\mathbf{p}^{\partial\Omega}] - C, 0, \int_{\partial S_0} \mathbf{p}^{\partial S_0} ds - C), \\ T(\varepsilon, q)[\mathbf{p}] &:= (T^{\partial S}(\varepsilon, q)[\mathbf{p}^{\partial\Omega}], T^{\partial\Omega}(\varepsilon, q)[\mathbf{p}^{\partial S_0}], 0), \end{aligned}$$

so that we can write A in the following form: on F_0 ,

$$A(\varepsilon, q) = L + K(h) + \varepsilon T(\varepsilon, q). \quad (8.81)$$

We first consider the operator $L + K(h)$. According to (8.3), the operator L is Fredholm with index zero and since for each h in $\overline{\Omega}_\delta$, $K(h)$ is compact, we deduce that $L + K(h)$ is Fredholm with index zero. It follows that in order to prove that $L + K(h)$ is an isomorphism, it is sufficient to prove that its kernel is trivial.

Consider $\mathbf{p} := (\mathbf{p}^{\partial S_0}, \mathbf{p}^{\partial\Omega}, C)$ in $F_{-\frac{1}{2}}$ such that $(L + K(h))[\mathbf{p}] = 0$. Since the logarithmic capacity $\text{Cap}(\partial\Omega)$ of $\partial\Omega$ satisfies $\text{Cap}(\partial\Omega) \neq 1$, according to (8.5), the second equation $SL[\mathbf{p}^{\partial\Omega}] = 0$ implies $\mathbf{p}^{\partial\Omega} = 0$. Then reporting

in the first equation, we get $SL[\mathbf{p}^{S_0}] = C$, whereas the third equation reads $\int_{\partial S_0} \mathbf{p}^{\partial S_0} ds = 0$. Thus according to (8.4), we obtain $\mathbf{p}^{\partial S_0} = 0$ and thus $C = 0$. This proves that the kernel of $L + K(h)$ is trivial, and consequently that for any h in Ω_δ , $L + K(h)$ is an isomorphism.

Now using that the dependence of K on h is Lipschitz, we deduce that $L + K(h)$ has locally a bounded inverse. By compactness of $\overline{\Omega_\delta}$, it follows that $L + K(h)$ has a bounded inverse for h running over $\overline{\Omega_\delta}$.

Since the operators $(T_\varepsilon)_{\varepsilon \in (0,1)}$ are bounded in the space of bounded operators from F_0 to F_1 we can then easily deduce the result from (8.81). This concludes the proof of Lemma 8.9.

8.4.4 Proof of Lemma 8.10

We use Lemma 8.3 with $\mathcal{C} = \partial\Omega$, $b = G$ and $\mathbf{p}^{\mathcal{C}} = \mathbf{p}^{\partial\Omega}$ to obtain that $\tilde{K}^{\partial S_0}$ satisfies (8.57). Next we apply Lemma 8.3, (i) for any (ε, q) in $\Omega_{\delta, \varepsilon_0}$, with $\mathcal{C} = \partial\Omega$, $b(x, y) = \eta_I(\varepsilon, q, x, y)$ and $\mathbf{p}^{\mathcal{C}} = \mathbf{p}^{\partial\Omega}$ and with $\mathcal{C} = \partial S_0$, $b(x, y) = \eta_I(\varepsilon, \vartheta, x, -y, h)$ and $\mathbf{p}^{\mathcal{C}} = \mathbf{p}_{S_0}$ to get that $T^{\partial S_0}(\varepsilon, q)$ and $T^{\partial\Omega}(\varepsilon, q)$ satisfy (8.58). This ends the proof of Lemma 8.10

8.4.5 Proof of Lemma 8.11

Let $\delta > 0$. Let us first observe that for any (ε, q) in Ω_δ , for any $\mathbf{p} := (\mathbf{p}^{\partial S_0}, \mathbf{p}^{\partial\Omega}, C)$ in F_0 satisfying the condition

$$\int_{\partial S_0} \mathbf{p}^{\partial S_0} ds = 0, \quad (8.82)$$

one has

$$\mathfrak{A}(\varepsilon, q)[\mathbf{p}] = A(\varepsilon, q)[\mathbf{p}].$$

Indeed (8.62) and first order Taylor expansions yield with (8.63) and (8.64) that

$$\begin{aligned} K^{\partial S_0}(\varepsilon, q)[\mathbf{p}^{\partial\Omega}] - \tilde{K}^{\partial S_0}(h)[\mathbf{p}^{\partial\Omega}] &= \varepsilon T^{\partial S_0}(\varepsilon, q)[\mathbf{p}^{\partial\Omega}], \\ K^{\partial\Omega}(\varepsilon, q)[\mathbf{p}^{\partial S_0}] &= \varepsilon T^{\partial\Omega}(\varepsilon, q)[\mathbf{p}^{\partial S_0}]. \end{aligned}$$

We emphasize in particular that the last equality relies on the condition (8.82).

Now, consider ε_0 in $(0, 1)$ obtained by applying Lemma 8.9. For any (ε, q) in $\Omega_{\delta, \varepsilon_0}$, consider $\mathbf{p}_r(\varepsilon, q, \cdot) = (\mathbf{p}_r^{\partial S_0}(\varepsilon, q, \cdot), p_r^{\partial\Omega}(\varepsilon, q, \cdot), C_r(\varepsilon, q))$ given by (8.65). It belongs to F_0 and satisfies (8.72) and consequently

$$\mathfrak{A}(\varepsilon, q)[p_r(\varepsilon, q, \cdot)] = A(\varepsilon, q)[p_r(\varepsilon, q, \cdot)] = \mathfrak{g}(\varepsilon, q, \cdot).$$

Moreover we have the estimate

$$\|p_r(\varepsilon, q, \cdot)\|_{F_0} \leq \|A(\varepsilon, q)^{-1}\|_{\mathcal{L}(F_1; F_0)} \|\mathfrak{g}(\varepsilon, q, \cdot)\|_{F_1}.$$

The estimates (8.53) and (8.61) entail that \mathbf{p}_r is in $L^\infty(Q_{\delta, \varepsilon_0}; F_0)$, which concludes the proof of Lemma 8.11.

8.5 Asymptotic expansion of the potential part: Proof of Proposition 8.5

The proof of Proposition 8.5 is very close to the one of Proposition 8.4. We will only explain how to transform the (Neumann) problem defining the Kirchhoff potentials into a Dirichlet one, so that the proof of Proposition 8.5 follows from a tedious adaptation of the steps of the proof of Proposition 8.4 detailed in Subsection 8.3.

We emphasize that the indices below correspond to coordinates in \mathbb{R}^3 and are in normal font type (while indices related to the order in an asymptotic development in powers of ε are written in italic type).

We consider the functions $\bar{\varphi}_{j,\varepsilon}(q, \cdot)$, for $j = 1, 2, 3$, as the solution to the following Dirichlet boundary value problem in $\mathcal{F}_\varepsilon(q)$:

$$-\Delta \bar{\varphi}_{j,\varepsilon}(q, \cdot) = 0 \quad \text{in } \mathcal{F}_\varepsilon(q), \quad (8.83a)$$

$$\bar{\varphi}_{j,\varepsilon}(q, \cdot) = \bar{K}_j(q, \cdot) + c_{j,\varepsilon}(q) \quad \text{on } \partial\mathcal{S}_\varepsilon(q), \quad (8.83b)$$

$$\bar{\varphi}_{j,\varepsilon}(q, \cdot) = 0 \quad \text{on } \partial\Omega, \quad (8.83c)$$

where the functions $\bar{K}_j(q, \cdot)$ are given by

$$\bar{K}_j(q, \cdot) := \begin{cases} \frac{1}{2}|x-h|^2 & \text{if } j = 1, \\ -R(\vartheta)^t(x-h) \cdot e_2 & \text{if } j = 2, \\ R(\vartheta)^t(x-h) \cdot e_1 & \text{if } j = 3, \end{cases}$$

where e_1 and e_2 are the unit vectors of the canonical basis, and the constants $c_{j,\varepsilon}(q)$ are such that:

$$\int_{\partial\mathcal{S}_\varepsilon(q)} \frac{\partial \bar{\varphi}_{j,\varepsilon}}{\partial n}(q, \cdot) \, ds = 0. \quad (8.83d)$$

Precisely, the constants $c_{j,\varepsilon}(q)$ are given by

$$c_{j,\varepsilon}(q) = C_\varepsilon(q) \int_{\partial\mathcal{S}_\varepsilon(q)} \frac{\partial \bar{\varphi}_{j,\varepsilon}}{\partial n}(q, \cdot) \, ds, \quad (8.84)$$

where $\bar{\phi}_{j,\varepsilon}$, $j = 1, 2, 3$ are the solutions of

$$-\Delta \bar{\phi}_{j,\varepsilon}(q, \cdot) = 0 \quad \text{in } \mathcal{F}_\varepsilon(q), \quad (8.85)$$

$$\bar{\phi}_{j,\varepsilon}(q, \cdot) = \bar{K}_j(q, \cdot) \quad \text{on } \partial\mathcal{S}_\varepsilon(q), \quad (8.86)$$

$$\bar{\phi}_{j,\varepsilon}(q, \cdot) = 0 \quad \text{on } \partial\Omega. \quad (8.87)$$

We will use the vector notation:

$$\bar{\varphi}_\varepsilon := (\bar{\varphi}_{1,\varepsilon}, \bar{\varphi}_{2,\varepsilon}, \bar{\varphi}_{3,\varepsilon})^t. \quad (8.88)$$

Up to a rotation, the functions $\bar{\varphi}_{j,\varepsilon}(q, \cdot)$ are harmonically conjugated to the Kirchhoff's potentials $\varphi_{j,\varepsilon}(q, \cdot)$ (see (8.28)), as shown in the following result.

Lemma 8.13. For any (ε, q) in \mathfrak{Q} , with $q = (\vartheta, h)$, there holds in $\mathcal{F}_\varepsilon(q)$,

$$\nabla \varphi_{j,\varepsilon}(q, \cdot) = \nabla^\perp \check{\varphi}_{j,\varepsilon}(q, \cdot), \quad (8.89)$$

where

$$\left(\check{\varphi}_{1,\varepsilon}(q, \cdot), \check{\varphi}_{2,\varepsilon}(q, \cdot), \check{\varphi}_{3,\varepsilon}(q, \cdot) \right) := \mathcal{R}(\vartheta) \left(\bar{\varphi}_{1,\varepsilon}(q, \cdot), \bar{\varphi}_{2,\varepsilon}(q, \cdot), \bar{\varphi}_{3,\varepsilon}(q, \cdot) \right). \quad (8.90)$$

Proof We recall that for any (ε, q) in \mathfrak{Q} , the system

$$\operatorname{div} u = 0 \quad \text{in } \mathcal{F}_\varepsilon(q), \quad (8.91a)$$

$$\operatorname{curl} u = 0 \quad \text{in } \mathcal{F}_\varepsilon(q), \quad (8.91b)$$

$$u \cdot n = 0 \quad \text{on } \partial\Omega, \quad (8.91c)$$

$$u \cdot n = K_j(q, \cdot) \quad \text{on } \partial\mathcal{S}_\varepsilon(q), \quad (8.91d)$$

$$\int_{\partial\mathcal{S}_\varepsilon(q)} u \cdot \tau \, ds = 0, \quad (8.91e)$$

has a unique solution u , say in $H^1(\mathcal{F}_\varepsilon(q))$. Then one observes that both $\nabla \varphi_{j,\varepsilon}(q, \cdot)$ and $\nabla^\perp \check{\varphi}_{j,\varepsilon}(q, \cdot)$ solve (8.91). In particular let us emphasize that, on $\partial\mathcal{S}_\varepsilon(q)$,

$$\left(n \cdot \nabla^\perp \bar{\varphi}_{j,\varepsilon}(q, \cdot) \right)_{j=1,2,3} = \left(\frac{\partial \bar{K}_j}{\partial \tau}(q, \cdot) \right)_{j=1,2,3} = \mathcal{R}(\vartheta)^t \mathbf{K}(q, \cdot),$$

so that, for $j = 1, 2, 3$, $\nabla^\perp \check{\varphi}_{j,\varepsilon}(q, \cdot)$ satisfies (8.91d), and the condition (8.83d) ensures that (8.91e) is satisfied. \square

In the case without exterior boundary we consider in the same way $\bar{\varphi}_j^{\mathcal{S}_0}$ as the solution of

$$-\Delta \bar{\varphi}_j^{\mathcal{S}_0} = 0 \quad \text{in } \mathbb{R}^2 \setminus \mathcal{S}_0, \quad (8.92a)$$

$$\bar{\varphi}_j^{\mathcal{S}_0}(\cdot) = \bar{K}_j(0, \cdot) + c_j \quad \text{on } \partial\mathcal{S}_0, \quad (8.92b)$$

$$\bar{\varphi}_j^{\mathcal{S}_0}(x) \rightarrow 0 \quad \text{as } |x| \rightarrow +\infty, \quad (8.92c)$$

where the constant $c_{\mathcal{S}_0,j}$ is such that

$$\int_{\partial\mathcal{S}_0} \frac{\partial \bar{\varphi}_j^{\mathcal{S}_0}}{\partial n} \, ds = 0. \quad (8.92d)$$

The existence and uniqueness of such a constant $c_{\mathcal{S}_0,j}$ is provided by a similar argument as for (8.84)-(8.85). Proceeding as in the proof of Lemma 8.13 we get

$$\nabla \varphi_j^{\mathcal{S}_0} = \nabla^\perp \bar{\varphi}_j^{\mathcal{S}_0}, \quad (8.93)$$

where the functions $\varphi_j^{\mathcal{S}_0}$, for $j = 1, 2, 3$, are the Kirchhoff's potentials in $\mathbb{R}^2 \setminus \mathcal{S}_0$ defined in (4.2). As before we introduce the vector notation for the functions $\bar{\varphi}_j^{\mathcal{S}_0}$:

$$\bar{\varphi}^{\mathcal{S}_0} := (\bar{\varphi}_1^{\mathcal{S}_0}, \bar{\varphi}_2^{\mathcal{S}_0}, \bar{\varphi}_3^{\mathcal{S}_0}). \quad (8.94)$$

Then, following the strategy of Proposition 8.4 we obtain the following result.

Proposition 8.14. *Let $\delta > 0$. There exists ε_0 in $(0, 1)$ and*

(i) *there exists $\mathbf{p}_r^{\partial S_0}$ in $L^\infty(\mathfrak{Q}_{\delta, \varepsilon_0}; L^2(\partial S_0; \mathbb{R}^3))$ such that for any (ε, q) in $\mathfrak{Q}_{\delta, \varepsilon_0}$, with $q = (\vartheta, h)$, for any X in ∂S_0 ,*

$$\frac{\partial \bar{\varphi}_\varepsilon}{\partial n}(q, \varepsilon R(\vartheta)X + h) = I_\varepsilon \left(\frac{\partial \bar{\varphi}^{S_0}}{\partial n}(X) + \varepsilon^2 \mathbf{p}_r^{\partial S_0}(\varepsilon, q, X) \right), \quad (8.95)$$

(ii) *there exists $\mathbf{p}_r^{\partial \Omega}$ in $L^\infty(\mathfrak{Q}_{\delta, \varepsilon_0}; L^2(\partial \Omega; \mathbb{R}^3))$ such that for any (ε, q) in $\mathfrak{Q}_{\delta, \varepsilon_0}$, for any x in $\partial \Omega$,*

$$\frac{\partial \bar{\varphi}_\varepsilon}{\partial n}(q, x) = I_\varepsilon \varepsilon^2 \mathbf{p}_r^{\partial \Omega}(\varepsilon, q, x). \quad (8.96)$$

Once Proposition 8.14 is obtained, Proposition 8.5 follows by using additionally Lemma 8.13 and (8.93). We omit the details.

9 Proof of the normal forms

This section is devoted to the proof of the normal forms in Proposition 6.3, Proposition 7.10 and Proposition 7.17, as well as the expansion of the added inertia in Proposition 9.1 that was used to establish Lemma 6.5. The proof of the normal forms (6.7), (7.10) and (7.18) consists first in expanding the functions $M_\varepsilon(q)$, $\langle \Gamma_\varepsilon(q), p, p \rangle$ and $F_\varepsilon(q, p)$ with respect to ε thanks to the expansions of the previous sections and to Lamb's lemma (Lemma 4.3), and then in plugging these expansions into (6.1). Next, further modifications are needed in order to reach the exact forms (6.7), (7.10) and (7.18).

9.1 Asymptotic expansion of the added inertia and the Christoffel symbols

In this subsection, we use the asymptotic developments of Section 8 to deduce expansions for the added inertia matrix and for the Christoffel symbols.

We begin by giving the expansions in terms of ε of the inertia matrix $M_{a, \varepsilon}(q)$ which is the counterpart for the body of size ε of the added mass $M_a(q)$ defined in (2.15b). Precisely, it is defined for (ε, q) in \mathfrak{Q} by

$$M_{a, \varepsilon}(q) := \int_{\partial S_\varepsilon(q)} \varphi_\varepsilon(q, \cdot) \otimes \frac{\partial \varphi_\varepsilon}{\partial n}(q, \cdot) \, ds = \int_{\partial S_\varepsilon(q)} \varphi_\varepsilon(q, \cdot) \otimes \mathbf{K}_\varepsilon(q, \cdot) \, ds.$$

The function φ_ε mentioned above is defined in (8.28), (8.29). Let us also recall that the matrix $M_{a, \vartheta}^{S_0}$ is defined in (4.8) and (4.9), and I_ε is defined in (4.27). The expansion is as follows.

Proposition 9.1. *Let $\delta > 0$. There exists ε_0 in $(0, 1)$ and a function M_r in $L^\infty(\mathfrak{Q}_{\delta, \varepsilon_0}; \mathbb{R}^{3 \times 3})$ depending on S_0 and Ω , such that, for all (ε, q) in $\mathfrak{Q}_{\delta, \varepsilon_0}$, with $q = (\vartheta, h)$,*

$$M_{a, \varepsilon}(q) = \varepsilon^2 I_\varepsilon \left(M_{a, \vartheta}^{S_0} + \varepsilon^2 M_r(\varepsilon, q) \right) I_\varepsilon. \quad (9.1)$$

Proof Using a change of variable, (2.9) and (8.29) we get that for (ε, q) in \mathfrak{Q} ,

$$\mathbf{K}_\varepsilon(q, \varepsilon R(\vartheta) \cdot + h) = I_\varepsilon \mathcal{R}(\vartheta) \mathbf{K}(0, \cdot) \text{ on } \partial \mathcal{S}_0. \quad (9.2)$$

It follows that

$$M_{a,\varepsilon}(q) = \varepsilon \int_{\partial \mathcal{S}_0} \varphi_\varepsilon(q, \varepsilon R(\vartheta) \cdot + h) \otimes I_\varepsilon \mathcal{R}(\vartheta) \mathbf{K}(0, \cdot) \, ds.$$

We now apply Proposition 8.5, (i) to get

$$\begin{aligned} M_{a,\varepsilon}(q) &= \varepsilon^2 I_\varepsilon \left(\int_{\partial \mathcal{S}_0} \mathcal{R}(\vartheta) \left(\varphi^{S_0} + \check{c}(\varepsilon, q) + \varepsilon^2 \varphi_r(\varepsilon, q, \cdot) \right) \otimes \mathcal{R}(\vartheta) \mathbf{K}(0, \cdot) \, ds \right) I_\varepsilon \\ &= \varepsilon^2 I_\varepsilon \left(M_{a,\vartheta}^{S_0} + \varepsilon^2 \mathcal{R}(\vartheta) \int_{\partial \mathcal{S}_0} \varphi_r(\varepsilon, q, \cdot) \otimes \mathbf{K}(0, \cdot) \, ds \mathcal{R}(\vartheta)^t \right) I_\varepsilon, \end{aligned}$$

since

$$\int_{\partial \mathcal{S}_0} \check{c}(\varepsilon, q) \otimes \mathbf{K}(0, \cdot) \, ds = \check{c}(\varepsilon, q) \otimes \int_{\partial \mathcal{S}_0} \mathbf{K}(0, \cdot) \, ds = 0,$$

and

$$M_{a,\vartheta}^{S_0} = \mathcal{R}(\vartheta) \int_{\partial \mathcal{S}_0} \varphi^{S_0} \otimes \mathbf{K}(0, \cdot) \, ds \mathcal{R}(\vartheta)^t,$$

thanks to (4.2b), (4.8) and (4.9). Above ε_0 belongs to $(0, 1)$ and φ_r is in the space $L^\infty(\mathfrak{Q}_{\delta,\varepsilon_0}; L^2(\partial \mathcal{S}_0; \mathbb{R}^3))$. Then we set

$$M_r(\varepsilon, q) := \mathcal{R}(\vartheta) \int_{\partial \mathcal{S}_0} \varphi_r(\varepsilon, q, \cdot) \otimes \mathbf{K}(0, \cdot) \, ds \mathcal{R}(\vartheta)^t,$$

and we observe that M_r is in $L^\infty(\mathfrak{Q}_{\delta,\varepsilon_0}; \mathbb{R})$ and depends only on \mathcal{S}_0 and Ω , which concludes the proof of Proposition 9.1. \square

We now consider the Christoffel symbols $\Gamma_\varepsilon^{\text{rot}}$ given for (ε, q) in \mathfrak{Q} and $p = (\omega, \ell) \in \mathbb{R}^3$, by

$$\langle \Gamma_\varepsilon^{\text{rot}}(q), p, p \rangle := - \begin{pmatrix} 0 \\ P_a^\varepsilon \end{pmatrix} \times p - \omega M_{a,\varepsilon}(q) \begin{pmatrix} 0 \\ \ell^\perp \end{pmatrix} \in \mathbb{R}^3, \quad (9.3)$$

where P_a^ε denotes the two last lines of $M_{a,\varepsilon}(q)p$. The formula (9.3) is the counterpart for a body of size ε of the Christoffel symbols given by (2.25) when $\varepsilon = 1$.

The next result proves that the leading term of $\Gamma_\varepsilon^{\text{rot}}$ is given, up to an appropriate scaling, by the Christoffel symbols $\langle \Gamma_\vartheta^{S_0}, p, p \rangle$ of the solid as if it was immersed in a fluid filling the plane. We recall that $\langle \Gamma_\vartheta^{S_0}, p, p \rangle$ is defined in (4.11). Precisely, we have the following result.

Proposition 9.2. *Let $\delta > 0$ be fixed. There exists ε_0 in $(0, 1)$ and Γ_r^{rot} in $L^\infty(\mathfrak{Q}_{\delta,\varepsilon_0}; \mathcal{BL}(\mathbb{R}^3 \times \mathbb{R}^3; \mathbb{R}^3))$ depending on \mathcal{S}_0 , γ and Ω , such that for any (ε, q) in $\mathfrak{Q}_{\delta,\varepsilon_0}$, with $q = (\vartheta, h)$, for any $p = (\omega, \ell)$ in \mathbb{R}^3 ,*

$$\langle \Gamma_\varepsilon^{\text{rot}}(q), p, p \rangle = \varepsilon I_\varepsilon (\langle \Gamma_\vartheta^{S_0}, I_\varepsilon p, I_\varepsilon p \rangle) + \varepsilon^2 \langle \Gamma_r^{\text{rot}}(\varepsilon, q), I_\varepsilon p, I_\varepsilon p \rangle. \quad (9.4)$$

Proposition 9.2 follows from Proposition 9.1 by straightforward computations.

Finally we study the Christoffel symbols $\Gamma_\varepsilon^{\partial\Omega}$ given for (ε, q) in \mathfrak{Q} and $p = (p_1, p_2, p_3) \in \mathbb{R}^3$, by

$$\langle \Gamma_\varepsilon^{\partial\Omega}(q), p, p \rangle := \left(\sum_{1 \leq k, l \leq 3} (\Gamma_\varepsilon^{\partial\Omega})_{k,l}^j(q) p_k p_l \right)_{1 \leq j \leq 3} \in \mathbb{R}^3, \quad (9.5)$$

where for every $j, k, l \in \{1, 2, 3\}$, we set

$$(\Gamma_\varepsilon^{\partial\Omega})_{k,l}^j(q) := \frac{1}{2} [A_{\varepsilon, kj}^l(q) + A_{\varepsilon, jl}^k(q) - A_{\varepsilon, kl}^j(q)],$$

with

$$A_{\varepsilon, kj}^l(q) := \int_{\partial\Omega} \left(\frac{\partial \varphi_{j,\varepsilon}}{\partial \tau} \frac{\partial \varphi_{k,\varepsilon}}{\partial \tau} K_l \right) (q, \cdot) ds.$$

The Christoffel symbols $\Gamma_\varepsilon^{\partial\Omega}$ are the counterpart $\Gamma^{\partial\Omega}$ defined in (2.26). They are expanded as follows.

Proposition 9.3. *Let $\delta > 0$ be fixed. There exists ε_0 in $(0, 1)$ and $\Gamma_r^{\partial\Omega}$ in $L^\infty(\mathfrak{Q}_{\delta, \varepsilon_0}; \mathcal{B}\mathcal{L}(\mathbb{R}^3 \times \mathbb{R}^3; \mathbb{R}^3))$ depending on \mathcal{S}_0, γ and Ω , such that for any (ε, q) in $\mathfrak{Q}_{\delta, \varepsilon_0}$, with $q := (\vartheta, h)$, for any $p := (\omega, \ell)$ in \mathbb{R}^3 ,*

$$\langle \Gamma_\varepsilon^{\partial\Omega}(q), p, p \rangle = \varepsilon^3 I_\varepsilon \langle \Gamma_r^{\partial\Omega}(\varepsilon, q), I_\varepsilon p, I_\varepsilon p \rangle. \quad (9.6)$$

Proof Proposition 9.3 follows from Proposition 8.5, (iii). Indeed, (9.5) can be rewritten as:

$$\langle \Gamma_\varepsilon^{\partial\Omega}(q), p, p \rangle = \int_{\partial\Omega} \left[\frac{\partial \varphi_\varepsilon}{\partial \tau} (\mathbf{K}_\varepsilon \cdot p) \left(\frac{\partial \varphi_\varepsilon}{\partial \tau} \cdot p \right) - \frac{1}{2} \mathbf{K}_\varepsilon \left(\frac{\partial \varphi_\varepsilon}{\partial \tau} \cdot p \right)^2 \right] (q, \cdot) ds.$$

Observe that \mathbf{K}_ε is actually independent of ε on $\partial\Omega$. Let us denote $\hat{\mathbf{K}}_\varepsilon := \varepsilon I_\varepsilon^{-1} \mathbf{K}_\varepsilon$. According to (8.32), we obtain:

$$\langle \Gamma_\varepsilon^{\partial\Omega}(q), p, p \rangle = \varepsilon^3 I_\varepsilon \int_{\partial\Omega} \left[\mathbf{p}_r^{\partial\Omega}(\hat{\mathbf{K}}_\varepsilon \cdot I_\varepsilon p) (\mathbf{p}_r^{\partial\Omega} \cdot I_\varepsilon p) - \frac{1}{2} \hat{\mathbf{K}}_\varepsilon (\mathbf{p}_r^{\partial\Omega} \cdot I_\varepsilon p)^2 \right] (q, \cdot) ds,$$

which gives the expected result. \square

9.2 Asymptotic expansion of E_ε

We now consider the expansion of E_ε which is given, for (ε, q) in \mathfrak{Q} , by

$$E_\varepsilon(q) := -\frac{1}{2} \int_{\partial\mathcal{S}_\varepsilon(q)} \left| \frac{\partial \psi_\varepsilon}{\partial n}(q, \cdot) \right|^2 \mathbf{K}_\varepsilon(q, \cdot) ds.$$

This formula is the counterpart of (2.17b) for a body of size ε . We recall that the function $\psi_\varepsilon(q, \cdot)$ is defined in (8.14) and the vector field $\mathbf{K}_\varepsilon(q, \cdot)$ in (8.27)-(8.29).

The two first orders of the asymptotic expansion will be given respectively thanks to two terms $\mathbf{E}_\theta(q)$ and $\mathbf{E}_I(q)$ which are now defined.

First we set

$$\mathbf{E}_\theta(q) := - \begin{pmatrix} u^\Omega(h) \cdot \zeta_\vartheta \\ u^\Omega(h)^\perp \end{pmatrix}, \text{ where } q = (\vartheta, h).$$

We recall that u^Ω and ζ_ϑ were defined in (2.42) and (4.14)-(4.12) respectively.

Next we define $\mathbf{E}_I(q)$ as

$$\mathbf{E}_I(q) := \mathbf{E}_I^a(q) + \mathbf{E}_I^b(q) + \mathbf{E}_I^c(q), \quad (9.7)$$

where the three addends are given by the definitions below.

The inertial subprincipal term \mathbf{E}_I^a . The definition of the term \mathbf{E}_I^a will use some functions of the entries of the matrix $M_a^{\mathcal{S}_0}$ defined in (4.8). Let us first recall that we decomposed $M_a^{\mathcal{S}_0}$ in (4.23). We also define the real traceless symmetric 2×2 matrix M^\dagger defined by

$$M^\dagger = (M_{i,j}^\dagger)_{1 \leq i,j \leq 2} := \frac{1}{2} \left(M_b^{\mathcal{S}_0}(\perp) + (M_b^{\mathcal{S}_0}(\perp))^t \right) = \frac{1}{2} \left(M_b^{\mathcal{S}_0}(\perp) - (\perp) M_b^{\mathcal{S}_0} \right), \quad (9.8)$$

where (\perp) denotes the 2×2 matrix

$$(\perp) := \begin{pmatrix} 0 & -1 \\ 1 & 0 \end{pmatrix}. \quad (9.9)$$

The matrix M^\dagger depends only on \mathcal{S}_0 . Its coefficients can be described as follows:

$$M_{1,1}^\dagger = -M_{2,2}^\dagger = \int_{\partial \mathcal{S}_0} \frac{\partial \varphi_3^{\mathcal{S}_0}}{\partial n} \varphi_2^{\mathcal{S}_0} ds, \quad \text{and} \quad (9.10a)$$

$$M_{1,2}^\dagger = M_{2,1}^\dagger = \frac{1}{2} \int_{\partial \mathcal{S}_0} \left(\frac{\partial \varphi_3^{\mathcal{S}_0}}{\partial n} \varphi_3^{\mathcal{S}_0} - \frac{\partial \varphi_2^{\mathcal{S}_0}}{\partial n} \varphi_2^{\mathcal{S}_0} \right) ds, \quad (9.10b)$$

where the functions $\varphi_j^{\mathcal{S}_0}$, for $j = 1, 2, 3$, are the Kirchhoff's potentials in $\mathbb{R}^2 \setminus \mathcal{S}_0$ defined in (4.2). Recalling (4.23), we also consider

$$M_{b,\vartheta}^{\mathcal{S}_0} := R(\vartheta) M_b^{\mathcal{S}_0} R(\vartheta)^t, \quad \mu_\vartheta^{\mathcal{S}_0} := R(\vartheta) \mu^{\mathcal{S}_0}, \quad M_\vartheta^\dagger := R(\vartheta) M^\dagger R(\vartheta)^t. \quad (9.11)$$

Then we define:

$$\mathbf{E}_I^a(q) := \begin{pmatrix} u^\Omega(h)^\perp M_\vartheta^\dagger u^\Omega(h)^\perp \\ 0 \\ 0 \end{pmatrix}, \text{ where } q = (\vartheta, h). \quad (9.12)$$

The weakly gyroscopic subprincipal term \mathbf{E}_1^b . Let us introduce the geometrical constant 2×2 matrix

$$\begin{aligned}\sigma &:= \int_{\partial\mathcal{S}_0} \frac{\partial\psi_{-1}^{\mathcal{S}_0}}{\partial n}(X) X \otimes X^\perp ds(X) + \zeta \otimes \zeta^\perp \\ &= \int_{\partial\mathcal{S}_0} \frac{\partial\psi_{-1}^{\mathcal{S}_0}}{\partial n}(X) (X \otimes X^\perp - \zeta \otimes \zeta^\perp) ds(X),\end{aligned}\quad (9.13)$$

which only depends on \mathcal{S}_0 . Next we introduce its symmetric part

$$\sigma^s := \frac{1}{2}(\sigma + \sigma^t), \quad (9.14)$$

and the associated field force $\mathbf{E}_1^b(q)$ defined, for $q = (\vartheta, h)$ in $\mathbb{R} \times \Omega$, by

$$\mathbf{E}_1^b(q) := \begin{pmatrix} -\langle D_x^2 \psi_0^\Omega(h, h), R(-2\vartheta)\sigma^s \rangle_{\mathbb{R}^{2 \times 2}} \\ 0 \\ 0 \end{pmatrix}. \quad (9.15)$$

Let us recall that the function ψ_0^Ω is defined in (2.39). The main property of \mathbf{E}_1^b is the following.

Lemma 9.4. *The vector field \mathbf{E}_1^b in $C^\infty(\mathbb{R} \times \Omega; \mathbb{R}^3)$ defined by (9.15) is weakly gyroscopic in the sense of Definition 7.9.*

Proof Multiply (9.15) by \tilde{p} and integrate. The conclusion follows from an integration by parts, crude bounds, Lemma 6.1, the smoothness of the function ψ_0^Ω and Lemma 7.8. \square

The drift subprincipal term \mathbf{E}_1^c . Let us introduce the force field $\mathbf{E}_1^c(q)$ defined, for $q = (\vartheta, h)$ in $\mathbb{R} \times \Omega$, by

$$\mathbf{E}_1^c(q) := - \begin{pmatrix} \zeta_\vartheta \cdot u_c(q) \\ (u_c(q))^\perp \end{pmatrix}. \quad (9.16)$$

Above $u_c(q)$ denotes the corrector velocity defined in (7.7).

Now the goal of this subsection is to establish the following result.

Proposition 9.5. *Let $\delta > 0$. There exists ε_0 in $(0, 1)$ and a function E_r in $L^\infty(\mathfrak{Q}_{\delta, \varepsilon_0}; \mathbb{R}^3)$ depending on \mathcal{S}_0 and Ω , such that for any (ε, q) in $\mathfrak{Q}_{\delta, \varepsilon_0}$,*

$$E_\varepsilon(q) = I_\varepsilon \left(\mathbf{E}_0(q) + \varepsilon \mathbf{E}_1(q) + \varepsilon^2 E_r(\varepsilon, q) \right). \quad (9.17)$$

Proof We proceed in three steps: first we use a change of variable in order to recast $E_\varepsilon(q)$ as an integral on the fixed boundary $\partial\mathcal{S}_0$. Then we plug the expansion of ψ^ε into this integral. Finally we use several times Lamb's lemma in order to compute the terms of the resulting expansion.

First, thanks to a change of variable, using (9.2), we get

$$E_\varepsilon(q) = -\frac{\varepsilon}{2} I_\varepsilon \mathcal{R}(\vartheta) \int_{\partial \mathcal{S}_0} \left| \frac{\partial \psi_\varepsilon}{\partial n}(q, \varepsilon R(\vartheta) \cdot + h) \right|^2 \mathbf{K}(0, \cdot) ds,$$

where $\mathbf{K}(q, \cdot)$ is the vector field defined in (2.9). Now let $\delta > 0$. Using Proposition 8.4 we get that there exists ε_0 in $(0, 1)$ such that for any (ε, q) in $\mathfrak{Q}_{\delta, \varepsilon_0}$,

$$E_\varepsilon(q) = I_\varepsilon \mathcal{R}(\vartheta) \left(\frac{1}{\varepsilon} \underline{\mathbf{E}}_{-1} + \underline{\mathbf{E}}_0(q) + \varepsilon \underline{\mathbf{E}}_1(q) + \varepsilon^2 \underline{\mathbf{E}}_r(\varepsilon, q) \right), \quad (9.18)$$

with

$$\begin{aligned} \underline{\mathbf{E}}_{-1} &:= -\frac{1}{2} \int_{\partial \mathcal{S}_0} \left| \frac{\partial \psi_{-1}^{\mathcal{S}_0}}{\partial n} \right|^2 \mathbf{K}(0, \cdot) ds, \\ \underline{\mathbf{E}}_0(q) &:= -\int_{\partial \mathcal{S}_0} \frac{\partial \psi_{-1}^{\mathcal{S}_0}}{\partial n} \left(\frac{\partial \psi_0^{\mathcal{S}_0}}{\partial n}(q, \cdot) - R(\vartheta)^t u^\Omega(h) \cdot \tau \right) \mathbf{K}(0, \cdot) ds, \end{aligned} \quad (9.19)$$

$$\underline{\mathbf{E}}_1(q) := \underline{\mathbf{E}}_1^a(q) + \underline{\mathbf{E}}_1^b(q), \quad (9.20)$$

where

$$\underline{\mathbf{E}}_1^a(q) := -\frac{1}{2} \int_{\partial \mathcal{S}_0} \left| \frac{\partial \psi_0^{\mathcal{S}_0}}{\partial n}(q, \cdot) - R(\vartheta)^t u^\Omega(h) \cdot \tau \right|^2 \mathbf{K}(0, \cdot) ds, \quad (9.21)$$

$$\underline{\mathbf{E}}_1^b(q) := -\int_{\partial \mathcal{S}_0} \frac{\partial \psi_{-1}^{\mathcal{S}_0}}{\partial n} \left(\frac{\partial \psi_1^{\mathcal{S}_0}}{\partial n}(q, \cdot) - \frac{\partial P_1}{\partial n}(q, \cdot) \right) \mathbf{K}(0, \cdot) ds, \quad (9.22)$$

and $\underline{\mathbf{E}}_r$ in $L^\infty(\mathfrak{Q}_{\delta, \varepsilon_0}; \mathbb{R}^3)$ depending only on \mathcal{S}_0 and Ω .

We now compute each term thanks to Lamb's lemma. More precisely we establish the following equalities:

$$\underline{\mathbf{E}}_{-1} = 0, \quad (9.23)$$

$$\mathcal{R}(\vartheta) \underline{\mathbf{E}}_0(q) = \mathbf{E}_0(q), \quad (9.24)$$

$$\mathcal{R}(\vartheta) \underline{\mathbf{E}}_1^a(q) = \mathbf{E}_1^a(q), \quad (9.25)$$

$$\mathcal{R}(\vartheta) \underline{\mathbf{E}}_1^b(q) = \mathbf{E}_1^b(q) + \mathbf{E}_1^c(q). \quad (9.26)$$

The proof of Proposition 9.5 is then concluded after observing that $\mathcal{R}(\vartheta) \underline{\mathbf{E}}_r$ is also in $L^\infty(\mathfrak{Q}_{\delta, \varepsilon_0}; \mathbb{R}^3)$ and depends only on \mathcal{S}_0 and Ω .

In order to simplify the notations we omit to write the dependence on q except if this dependence reduces on ϑ or h . Similarly we omit to write that the function \mathbf{K} , its coordinates K_j and the vector fields ξ_j , which appear thanks to Lamb's lemma, are evaluated at $q = 0$.

Proof of (9.23). Computation of $\underline{\mathbf{E}}_{-1}$. We use Lemma 4.3 with $u = v = \nabla^\perp \psi_{-1}^{\mathcal{S}_0}$ and observe that $\nabla^\perp \psi_{-1}^{\mathcal{S}_0}$ is tangent to \mathcal{S}_0 to obtain (9.23).

Proof of (9.24). Computation of $\underline{\mathbb{E}}_0$. We observe that

$$\nabla^\perp \psi_{-1}^{S_0} = -\frac{\partial \psi_{-1}^{S_0}}{\partial n} \tau \text{ on } \partial \mathcal{S}_0, \quad (9.27)$$

$$\tau \cdot \nabla^\perp \psi_0^{S_0} = -\frac{\partial \psi_0^{S_0}}{\partial n} \text{ on } \partial \mathcal{S}_0, \quad (9.28)$$

so that, for $j = 1, 2, 3$,

$$-\int_{\partial \mathcal{S}_0} \frac{\partial \psi_{-1}^{S_0}}{\partial n} \cdot \frac{\partial \psi_0^{S_0}}{\partial n} K_j \, ds = -\int_{\partial \mathcal{S}_0} \nabla \psi_{-1}^{S_0} \cdot \nabla \psi_0^{S_0} K_j \, ds$$

and we use Lemma 4.3 with $(u, v) = (\nabla^\perp \psi_{-1}^{S_0}, \nabla^\perp \psi_0^{S_0})$ to obtain

$$-\int_{\partial \mathcal{S}_0} \frac{\partial \psi_{-1}^{S_0}}{\partial n} \cdot \frac{\partial \psi_0^{S_0}}{\partial n} K_j \, ds = -\int_{\partial \mathcal{S}_0} (\xi_j \cdot \nabla^\perp \psi_{-1}^{S_0})(n \cdot \nabla^\perp \psi_0^{S_0}) \, ds.$$

Then we use again (9.27) and observe that applying the tangential derivative to (8.17b), taking (8.15) into account, yields

$$n \cdot \nabla^\perp \psi_0^{S_0} = \frac{\partial \psi_0^{S_0}}{\partial \tau} = R(\vartheta)^t u^\Omega(h)^\perp \cdot \tau \text{ on } \partial \mathcal{S}_0. \quad (9.29)$$

Thus

$$-\int_{\partial \mathcal{S}_0} \frac{\partial \psi_{-1}^{S_0}}{\partial n} \cdot \frac{\partial \psi_0^{S_0}}{\partial n} K_j \, ds = \int_{\partial \mathcal{S}_0} \frac{\partial \psi_{-1}^{S_0}}{\partial n} (\xi_j \cdot \tau) (R(\vartheta)^t u^\Omega(h)^\perp \cdot \tau) \, ds. \quad (9.30)$$

Then using (4.4d), (4.12) and (4.14), we get

$$\mathcal{R}(\vartheta) \underline{\mathbb{E}}_0 = \mathcal{R}(\vartheta) \left(\int_{\partial \mathcal{S}_0} \frac{\partial \psi_{-1}^{S_0}}{\partial n} (\xi_j \cdot R(\vartheta)^t u^\Omega(h)^\perp) \, ds \right)_j = - \begin{pmatrix} u^\Omega(h) \cdot \zeta_\vartheta \\ u^\Omega(h)^\perp \end{pmatrix} = \underline{\mathbb{E}}_0.$$

Proof of (9.25). Computation of $\underline{\mathbb{E}}_1^a$. We start with expanding the square in (9.21), to get

$$\underline{\mathbb{E}}_1^a = \underline{\mathbb{E}}_1^{a,1} + \int_{\partial \mathcal{S}_0} \frac{\partial \psi_0^{S_0}}{\partial n} (R(\vartheta)^t u^\Omega(h) \cdot \tau) \mathbf{K} \, ds - \frac{1}{2} \int_{\partial \mathcal{S}_0} |R(\vartheta)^t u^\Omega(h) \cdot \tau|^2 \mathbf{K} \, ds, \quad (9.31)$$

with

$$\underline{\mathbb{E}}_1^{a,1} = -\frac{1}{2} \int_{\partial \mathcal{S}_0} \left| \frac{\partial \psi_0^{S_0}}{\partial n} \right|^2 \mathbf{K} \, ds = -\frac{1}{2} \int_{\partial \mathcal{S}_0} |\nabla \psi_0^{S_0}|^2 \mathbf{K} \, ds + \frac{1}{2} \int_{\partial \mathcal{S}_0} \left| \frac{\partial \psi_0^{S_0}}{\partial \tau} \right|^2 \mathbf{K} \, ds.$$

We apply Lemma 4.3 with $u = v = \nabla^\perp \psi_0^{S_0}$ to get

$$\frac{1}{2} \int_{\partial \mathcal{S}_0} |\nabla \psi_0^{S_0}|^2 \mathbf{K} \, ds = \left(\int_{\partial \mathcal{S}_0} (\nabla^\perp \psi_0^{S_0} \cdot n) (\nabla^\perp \psi_0^{S_0} \cdot \xi_j) \, ds \right)_{j=1,2,3}$$

Let us denote by $\underline{\Xi}_{1,j}^{a,1}$, $j = 1, 2, 3$, the coordinates of the vector $\underline{\Xi}_1^{a,1}$. We use (9.29) to get

$$\underline{\Xi}_{1,j}^{a,1} = - \int_{\partial S_0} (\xi_j \cdot \nabla^\perp \psi_0^{S_0}) (R(\vartheta)^t u^\Omega(h)^\perp \cdot \tau) ds + \frac{1}{2} \int_{\partial S_0} (R(\vartheta)^t u^\Omega(h)^\perp \cdot \tau)^2 K_j ds.$$

Then we decompose $\xi_j \cdot \nabla^\perp \psi_0^{S_0}$ in normal and tangential parts and use again (9.29) to obtain:

$$\underline{\Xi}_{1,j}^{a,1} = \int_{\partial S_0} \frac{\partial \psi_0^{S_0}}{\partial n} (R(\vartheta)^t u^\Omega(h)^\perp \cdot \tau) (\xi_j \cdot \tau) ds - \frac{1}{2} \int_{\partial S_0} (R(\vartheta)^t u^\Omega(h)^\perp \cdot \tau)^2 K_j ds.$$

Now we plug this expression of $\underline{\Xi}_1^{a,1}$ into (9.31) to get

$$\begin{aligned} \underline{\Xi}_{1,j}^a &= -\frac{1}{2} \int_{\partial S_0} (R(\vartheta)^t u^\Omega(h)^\perp)^2 K_j ds + \int_{\partial S_0} \frac{\partial \psi_0^{S_0}}{\partial n} (R(\vartheta)^t u^\Omega(h) \cdot \tau) K_j ds \\ &\quad + \int_{\partial S_0} \frac{\partial \psi_0^{S_0}}{\partial n} (R(\vartheta)^t u^\Omega(h)^\perp \cdot \tau) (\xi_j \cdot \tau) ds. \end{aligned}$$

We observe that the first term in the right hand side vanishes and we combine the two other ones to get

$$\underline{\Xi}_{1,j}^a = \int_{\partial S_0} \frac{\partial \psi_0^{S_0}}{\partial n} (R(\vartheta)^t u^\Omega(h)^\perp) \cdot \xi_j ds.$$

Using (8.17c) we infer

$$\underline{\Xi}_{1,j}^a = 0 \text{ for } j = 2, 3. \quad (9.32)$$

Now for $j = 1$, we start with observing that

$$\underline{\Xi}_{1,1}^a = (R(\vartheta)^t u^\Omega(h)^\perp) \cdot \int_{\partial S_0} \frac{\partial \psi_0^{S_0}}{\partial n} x^\perp ds. \quad (9.33)$$

To compute the right hand side we introduce the matrix

$$\overline{M} := \int_{\partial S_0} \begin{pmatrix} \overline{\varphi}_2^{S_0} \\ \overline{\varphi}_3^{S_0} \end{pmatrix} \otimes \begin{pmatrix} \frac{\partial \overline{\varphi}_3^{S_0}}{\partial n} \\ -\frac{\partial \overline{\varphi}_2^{S_0}}{\partial n} \end{pmatrix} ds, \quad (9.34)$$

where the functions $\overline{\varphi}_j^{S_0}$, for $j = 1, 2, 3$, defined in (8.92), are harmonic conjugates to the functions $\varphi_j^{S_0}$.

Lemma 9.6. *For any $q := (\vartheta, h)$ in $\mathbb{R} \times \Omega$,*

$$\int_{\partial S_0} \frac{\partial \psi_0^{S_0}}{\partial n} (q, x) x^\perp ds(x) = \overline{M} R(\vartheta)^t u^\Omega(h)^\perp. \quad (9.35)$$

Proof First it follows from (8.92b) that, on $\partial\mathcal{S}_0$,

$$x^\perp = \begin{pmatrix} \overline{\varphi}_2^{\mathcal{S}_0} - c_2^{\mathcal{S}_0} \\ \overline{\varphi}_3^{\mathcal{S}_0} - c_3^{\mathcal{S}_0} \end{pmatrix}. \quad (9.36)$$

We now express the stream function $\psi_0^{\mathcal{S}_0}(q, \cdot)$ thanks to the functions $\overline{\varphi}_3^{\mathcal{S}_0}$ and $\overline{\varphi}_2^{\mathcal{S}_0}$. Let $q := (\vartheta, h)$ in $\mathbb{R} \times \Omega$. On $\partial\mathcal{S}_0$, it follows from (8.15), (8.17b), (8.92) and Proposition 8.1 that there exists c in \mathbb{R} such that, on $\overline{\mathbb{R}^2} \setminus \overline{\mathcal{S}_0}$,

$$\psi_0^{\mathcal{S}_0}(q, \cdot) = R(\vartheta)^t u^\Omega(h)^\perp \cdot \begin{pmatrix} \overline{\varphi}_3^{\mathcal{S}_0} \\ -\overline{\varphi}_2^{\mathcal{S}_0} \end{pmatrix} + c \psi_{-1}^{\mathcal{S}_0}.$$

Then using (4.4d), (8.17c) and (8.92d) we obtain $c = 0$, which proves that for any $q := (\vartheta, h)$ in $\mathbb{R} \times \Omega$, on $\overline{\mathbb{R}^2} \setminus \overline{\mathcal{S}_0}$,

$$\psi_0^{\mathcal{S}_0}(q, \cdot) = R(\vartheta)^t u^\Omega(h)^\perp \cdot \begin{pmatrix} \overline{\varphi}_3^{\mathcal{S}_0} \\ -\overline{\varphi}_2^{\mathcal{S}_0} \end{pmatrix} \quad (9.37)$$

Plugging (9.36) and (9.37) into the left hand side of (9.35) and using again (8.17c) establishes Lemma 9.6. \square

Then, combining (9.33) and (9.35), we obtain:

$$\underline{\Xi}_{1,1}^a = (R(\vartheta)^t u^\Omega(h)^\perp) \cdot \overline{M} R(\vartheta)^t u^\Omega(h)^\perp.$$

Let us now connect the matrices M^\dagger defined in (9.8) and \overline{M} defined in (9.34). Using integrations by parts and (8.93), we get, for any $i, j = 2, 3$,

$$\int_{\partial\mathcal{S}_0} \frac{\partial \overline{\varphi}_i^{\mathcal{S}_0}}{\partial n} \overline{\varphi}_j^{\mathcal{S}_0} ds = \int_{\partial\mathcal{S}_0} \frac{\partial \varphi_i^{\mathcal{S}_0}}{\partial n} \varphi_j^{\mathcal{S}_0} dx.$$

Combining this with (9.10) yields

$$M^\dagger = \frac{1}{2}(\overline{M} + \overline{M}^t). \quad (9.38)$$

Recalling the definition of M_ϑ^\dagger in (9.11), we deduce that

$$\underline{\Xi}_{1,1}^a = u^\Omega(h)^\perp M_\vartheta^\dagger u^\Omega(h)^\perp. \quad (9.39)$$

Gathering (9.12), (9.32) and (9.39) we obtain (9.25).

Proof of (9.26). Computation of $\underline{\Xi}_1^b$. We start with splitting $\underline{\Xi}_1^b$ into two parts as follows:

$$\underline{\Xi}_1^b = - \int_{\partial\mathcal{S}_0} \frac{\partial \psi_{-1}^{\mathcal{S}_0}}{\partial n} \frac{\partial \psi_1^{\mathcal{S}_0}}{\partial n} \mathbf{K} ds + \int_{\partial\mathcal{S}_0} \frac{\partial \psi_{-1}^{\mathcal{S}_0}}{\partial n} \frac{\partial P_1}{\partial n} \mathbf{K} ds.$$

Using (9.27) and (9.28), we see that the first term of the right hand side above is equal to

$$- \int_{\partial\mathcal{S}_0} \nabla^\perp \psi_{-1}^{\mathcal{S}_0} \cdot \nabla^\perp \psi_1^{\mathcal{S}_0} \mathbf{K} ds.$$

We denote $\underline{\mathbf{E}}_{1,j}^b$, $j = 1, 2, 3$, the coordinates of $\underline{\mathbf{E}}_1^b$. We apply Lemma 4.3 with $u = \nabla^\perp \psi_{-1}^{S_0}$ and $v = \nabla^\perp \psi_1^{S_0}$ for any $j = 1, 2, 3$, to get

$$\underline{\mathbf{E}}_{1,j}^b = - \int_{\partial S_0} \frac{\partial \psi_1^{S_0}}{\partial \tau} \xi_j \cdot \nabla^\perp \psi_{-1}^{S_0} ds + \int_{\partial S_0} \frac{\partial \psi_{-1}^{S_0}}{\partial n} \frac{\partial P_1}{\partial n} K_j ds.$$

We now use that, on ∂S_0 ,

$$\xi_j \cdot \nabla^\perp \psi_{-1}^{S_0} = - \frac{\partial \psi_{-1}^{S_0}}{\partial n} \xi_j \cdot \tau \quad \text{and} \quad \frac{\partial \psi_1^{S_0}}{\partial \tau} = \frac{\partial P_1}{\partial \tau},$$

the last identity being a consequence of (8.25b), to deduce

$$\underline{\mathbf{E}}_{1,j}^b = \int_{\partial S_0} \frac{\partial \psi_{-1}^{S_0}}{\partial n} \xi_j \cdot \nabla P_1 ds. \quad (9.40)$$

Using the expression of P_1 in (8.22), we obtain

$$\underline{\mathbf{E}}_{1,j}^b = - \langle D_x^2 \psi_0^\Omega(h, h), R(\vartheta) A_j^1 R(\vartheta)^t \rangle_{\mathbb{R}^{2 \times 2}} - D_x \psi_1^\Omega(q, h) \cdot R(\vartheta) A_j^2, \quad (9.41)$$

where

$$A_j^1 := \int_{\partial S_0} \frac{\partial \psi_{-1}^{S_0}}{\partial n} x \otimes \xi_j ds \quad \text{and} \quad A_j^2 := \int_{\partial S_0} \frac{\partial \psi_{-1}^{S_0}}{\partial n} \xi_j ds.$$

• We start with the case $j = 1$. Consider the first term in the right hand side of (9.41). Using (9.13) we see that $A_1^1 = \sigma - \zeta \otimes \zeta^\perp$ and we observe that, since $D_x^2 \psi_0^\Omega(h, h)$ is symmetric,

$$\langle D_x^2 \psi_0^\Omega(h, h), R(\vartheta) \sigma R(\vartheta)^t \rangle_{\mathbb{R}^{2 \times 2}} = \langle D_x^2 \psi_0^\Omega(h, h), R(\vartheta) \sigma^s R(\vartheta)^t \rangle,$$

where σ^s is the symmetric part of σ defined in (9.14). Then using that σ^s is a traceless symmetric 2×2 matrix we get

$$\langle D_x^2 \psi_0^\Omega(h, h), R(\vartheta) \sigma R(\vartheta)^t \rangle_{\mathbb{R}^{2 \times 2}} = \langle D_x^2 \psi_0^\Omega(h, h), R(-2\vartheta) \sigma^s \rangle = -\mathbf{E}_{1,1}^b(q),$$

where $\mathbf{E}_{1,1}^b(q)$ denotes the first coordinate of the vector field $\mathbf{E}_1^b(q)$ defined in (9.15). Therefore we obtain for $j = 1$,

$$- \langle D_x^2 \psi_0^\Omega(h, h), R(\vartheta) A_1^1 R(\vartheta)^t \rangle_{\mathbb{R}^{2 \times 2}} = \mathbf{E}_{1,1}^b(q) + \langle D_x^2 \psi_0^\Omega(h, h), \zeta_\vartheta \otimes \zeta_\vartheta^\perp \rangle_{\mathbb{R}^{2 \times 2}}.$$

Concerning the second term in the right hand side of (9.41), we use $A_1^2 = -\zeta^\perp$ (see (4.12)) to get that for $j = 1$,

$$-D_x \psi_1^\Omega(q, h) \cdot R(\vartheta) A_1^2 = D_x \psi_1^\Omega(q, h) \cdot \zeta_\vartheta^\perp.$$

Thus

$$\underline{\mathbf{E}}_{1,1}^b = \mathbf{E}_{1,1}^b(q) + \langle D_x^2 \psi_0^\Omega(h, h), \zeta_\vartheta \otimes \zeta_\vartheta^\perp \rangle_{\mathbb{R}^{2 \times 2}} + D_x \psi_1^\Omega(q, h) \cdot \zeta_\vartheta^\perp.$$

The last two terms in the right hand side can be expressed in terms of the corrector velocity $u_c(q)$ defined in (7.7), as follows from the following statement.

Lemma 9.7. For any $q = (\vartheta, h)$ in $\Omega \times \mathbb{R}$,

$$u_c(q) = \left(D_x^2 \psi_0^\Omega(h, h) \cdot \zeta_\vartheta + D_x \psi_1^\Omega(q, h) \right)^\perp. \quad (9.42)$$

Proof From the definition of ψ_c in (6.10) and the one of $u_c(q)$ in (7.7) we deduce that for any $q = (\vartheta, h)$ in $\Omega \times \mathbb{R}$,

$$u_c(q) = \left(D_x^2 \psi_0^\Omega(h, h) \cdot \zeta_\vartheta + D_{xh}^2 \psi_0^\Omega(h, h) \cdot \zeta_\vartheta \right)^\perp,$$

which yields (9.42) thanks to (8.19). \square

Hence we finally obtain

$$\mathbb{E}_{1,1}^b = \mathbb{E}_{1,1}^b(q) - \zeta_\vartheta \cdot u_c(q). \quad (9.43)$$

• On the other hand, for $j = 2$ or 3 , we have $A_j^1 = -\zeta \otimes \xi_j$ and $A_j^2 = -\xi_j$, and therefore

$$\begin{aligned} \underline{\mathbb{E}}_{1,j} &= \langle D_x^2 \psi_0^\Omega(h, h), R(\vartheta)(\zeta \otimes \xi_j) R(\vartheta)^t \rangle_{\mathbb{R}^{2 \times 2}} + D \psi_1^\Omega(q, h) \cdot R(\vartheta) \xi_j \\ &= \left(D_x^2 \psi_0^\Omega(h, h) \cdot \zeta_\vartheta + D \psi_1^\Omega(q, h) \right) \cdot R(\vartheta) \xi_j \\ &= -R(\vartheta)^t u_c(q)^\perp \cdot \xi_j. \end{aligned}$$

Thus

$$R(\vartheta)(\underline{\mathbb{E}}_{1,1}^b)_{j=2,3} = -u_c(q)^\perp. \quad (9.44)$$

Gathering (9.15), (9.16), (9.43) and (9.44) we obtain (9.26). This ends the proof of Proposition 9.5 \square

9.3 Asymptotic expansion of B_ε

We now tackle the expansion of B_ε which is given, for (ε, q) in \mathfrak{Q} , by

$$B_\varepsilon(q) := \int_{\partial \mathcal{S}_\varepsilon(q)} \frac{\partial \psi_\varepsilon}{\partial n}(q, \cdot) \left(\mathbf{K}_\varepsilon(q, \cdot) \times \frac{\partial \varphi_\varepsilon}{\partial \tau}(q, \cdot) \right) ds.$$

This formula is the counterpart of (2.17a) for a body of size ε . Let us recall that the Kirchhoff potentials φ_ε are defined in (8.28)-(8.29).

The expansion that we obtain for $B_\varepsilon(q)$ is given in the following statement where $B_\vartheta^{S_0}$ is defined in (4.15), M_ϑ^\dagger in (9.11), I_ε in (4.27) and where:

$$\mathbf{B}_1(q) := \begin{pmatrix} 0 \\ -2M_\vartheta^\dagger u^\Omega(h)^\perp \end{pmatrix} \text{ for } q = (\vartheta, h). \quad (9.45)$$

Proposition 9.8. Let $\delta > 0$. There exists ε_0 in $(0, 1)$ and a function B_r in $L^\infty(\mathfrak{Q}_{\delta, \varepsilon_0}; \mathbb{R}^3)$ depending only on \mathcal{S}_0 and Ω , such that for any (ε, q) in $\mathfrak{Q}_{\delta, \varepsilon_0}$, where $q = (\vartheta, h)$,

$$B_\varepsilon(q) = \varepsilon I_\varepsilon^{-1} \left(B_\vartheta^{S_0} + \varepsilon \mathbf{B}_1(q) + \varepsilon^2 B_r(\varepsilon, q) \right). \quad (9.46)$$

Proof We proceed as in the proof of Proposition 9.5. Let us state the following formula which is useful several times in the sequel:

$$\forall p_a := (\omega_a, \ell_a) \text{ and } p_b := (\omega_b, \ell_b) \text{ in } \mathbb{R} \times \mathbb{R}^2, \quad \varepsilon p_a \times p_b = I_\varepsilon \left((I_\varepsilon p_a) \times (I_\varepsilon p_b) \right). \quad (9.47)$$

By a change of variable, using (9.47) and (9.2), we get

$$B_\varepsilon(q) = \varepsilon I_\varepsilon^{-1} \mathcal{R}(\vartheta) \int_{\partial \mathcal{S}_0} \frac{\partial \psi_\varepsilon}{\partial n}(q, \varepsilon R(\vartheta) \cdot + h) \left(\mathbf{K}(0, \cdot) \times \mathcal{R}(\vartheta)^t \frac{\partial \varphi_\varepsilon}{\partial \tau}(q, \varepsilon R(\vartheta) \cdot + h) \right) ds.$$

Now let $\delta > 0$. We use Proposition 8.4 and Proposition 8.5 to obtain that there exists ε_0 in $(0, 1)$ such that for any (ε, q) in $\mathfrak{Q}_{\delta, \varepsilon_0}$,

$$B_\varepsilon(q) = \varepsilon I_\varepsilon^{-1} \mathcal{R}(\vartheta) \left(\underline{\mathbf{B}}_0 + \varepsilon \underline{\mathbf{B}}_1(q) + \varepsilon^2 B_r(\varepsilon, q) \right),$$

with

$$\begin{aligned} \underline{\mathbf{B}}_0 &:= \int_{\partial \mathcal{S}_0} \frac{\partial \psi_{-1}^{S_0}}{\partial n} \left(\mathbf{K}(0, \cdot) \times \frac{\partial \varphi^{S_0}}{\partial \tau} \right) ds, \\ \underline{\mathbf{B}}_1(q) &:= \int_{\partial \mathcal{S}_0} \left(\frac{\partial \psi_0^{S_0}}{\partial n} - R(\vartheta)^t u^\Omega(h) \cdot \tau \right) \left(\mathbf{K}(0, \cdot) \times \frac{\partial \varphi^{S_0}}{\partial \tau} \right) ds, \end{aligned}$$

and B_r in $L^\infty(\mathfrak{Q}_{\delta, \varepsilon_0}; \mathbb{R}^3)$ depending only on \mathcal{S}_0 and Ω .

We now compute each term thanks to Lamb's lemma. More precisely we will prove the following equalities:

$$\mathcal{R}(\vartheta) \underline{\mathbf{B}}_0 = B_\vartheta^{S_0}, \quad (9.48)$$

$$\mathcal{R}(\vartheta) \underline{\mathbf{B}}_1 = \mathbf{B}_1. \quad (9.49)$$

As in the proof of Proposition 9.8 we will omit to write the dependence on q , except if this dependence reduces to a dependence on ϑ or h , and it will be understood that the functions \mathbf{K} , its coordinates K_j and the vector fields ξ_j are evaluated at $q = 0$.

Proof of (9.48). Computation of $\underline{\mathbf{B}}^0$. For $j = 1, 2, 3$, we denote by $\underline{\mathbf{B}}_{0,j}$, the coordinates of $\underline{\mathbf{B}}_0$. We have

$$\begin{aligned} \underline{\mathbf{B}}_{0,1} &= \int_{\partial \mathcal{S}_0} \frac{\partial \psi_{-1}^{S_0}}{\partial n} \frac{\partial \varphi_3^{S_0}}{\partial \tau} K_2 ds - \int_{\partial \mathcal{S}_0} \frac{\partial \psi_{-1}^{S_0}}{\partial n} \frac{\partial \varphi_2^{S_0}}{\partial \tau} K_3 ds \\ &= - \int_{\partial \mathcal{S}_0} \nabla^\perp \psi_{-1}^{S_0} \cdot \nabla \varphi_3^{S_0} K_2 ds + \int_{\partial \mathcal{S}_0} \nabla^\perp \psi_{-1}^{S_0} \cdot \nabla \varphi_2^{S_0} K_3 ds, \end{aligned}$$

using (9.27). Then we use Lemma 4.3 with $(u, v) = (\nabla^\perp \psi_{-1}^{S_0}, \nabla \varphi_2^{S_0})$ and $(u, v) = (\nabla^\perp \psi_{-1}^{S_0}, \nabla \varphi_3^{S_0})$, (4.4d) and (9.27) to obtain

$$\underline{\mathbf{B}}_{0,1} = \int_{\partial \mathcal{S}_0} \frac{\partial \psi_{-1}^{S_0}}{\partial n} \left((\tau \cdot \xi_2)(n \cdot \xi_3) - (\tau \cdot \xi_3)(n \cdot \xi_2) \right) ds = -1.$$

Proceeding in the same way and using (4.12) we get

$$\underline{\mathbf{B}}_{0,2} = \int_{\partial S_0} \frac{\partial \psi_{-1}^{S_0}}{\partial n} \left((\tau \cdot \xi_3)(n \cdot \xi_1) - (\tau \cdot \xi_1)(n \cdot \xi_3) \right) ds = \zeta^\perp \cdot \xi_2,$$

and $\underline{\mathbf{B}}_{0,3} = \zeta^\perp \cdot \xi_3$. This gives (9.48).

Proof of (9.49). Computation of $\underline{\mathbf{B}}_1$. Let us start with the first coordinate $\underline{\mathbf{B}}_{1,1}$ of $\underline{\mathbf{B}}_1$, that is:

$$\begin{aligned} \underline{\mathbf{B}}_{1,1} &= - \int_{\partial S_0} \left(\frac{\partial \psi_0^{S_0}}{\partial n} - R(\vartheta)^t u^\Omega(h) \cdot \tau \right) \left(\frac{\partial \varphi_2^{S_0}}{\partial \tau} K_3 - \frac{\partial \varphi_3^{S_0}}{\partial \tau} K_2 \right) ds \\ &= \underline{\mathbf{B}}_{1,1}^a + \underline{\mathbf{B}}_{1,1}^b + \underline{\mathbf{B}}_{1,1}^c, \end{aligned}$$

with

$$\begin{aligned} \underline{\mathbf{B}}_{1,1}^a &:= - \int_{\partial S_0} \frac{\partial \psi_0^{S_0}}{\partial n} \frac{\partial \varphi_2^{S_0}}{\partial \tau} K_3 ds, \\ \underline{\mathbf{B}}_{1,1}^b &:= \int_{\partial S_0} \frac{\partial \psi_0^{S_0}}{\partial n} \frac{\partial \varphi_3^{S_0}}{\partial \tau} K_2 ds, \\ \underline{\mathbf{B}}_{1,1}^c &:= \int_{\partial S_0} R(\vartheta)^t u^\Omega(h) \cdot \left(\frac{\partial \varphi_2^{S_0}}{\partial \tau} K_3 - \frac{\partial \varphi_3^{S_0}}{\partial \tau} K_2 \right) \tau ds. \end{aligned}$$

We start with

$$\underline{\mathbf{B}}_{1,1}^a = \int_{\partial S_0} \nabla^\perp \psi_0^{S_0} \cdot \nabla \varphi_2^{S_0} K_3 ds - \int_{\partial S_0} \frac{\partial \psi_0^{S_0}}{\partial \tau} K_2 K_3 ds.$$

We use Lemma 4.3 with $u = \nabla^\perp \psi_0^{S_0}$ and $v = \nabla \varphi_2^{S_0}$ to obtain

$$\begin{aligned} \int_{\partial S_0} \nabla^\perp \psi_0^{S_0} \cdot \nabla \varphi_2^{S_0} K_3 ds &= \int_{\partial S_0} (\nabla^\perp \psi_0^{S_0} \cdot \xi_3) (\nabla \varphi_2^{S_0} \cdot n) ds \\ &\quad + \int_{\partial S_0} (\nabla^\perp \psi_0^{S_0} \cdot n) (\nabla \varphi_2^{S_0} \cdot \xi_3) ds \\ &= \int_{\partial S_0} \left(\frac{\partial \psi_0^{S_0}}{\partial \tau} K_3 - \frac{\partial \psi_0^{S_0}}{\partial n} \xi_3 \cdot \tau \right) K_2 ds + \int_{\partial S_0} \frac{\partial \psi_0^{S_0}}{\partial \tau} K_2 K_3 ds \\ &\quad + \int_{\partial S_0} \frac{\partial \psi_0^{S_0}}{\partial \tau} \frac{\partial \varphi_2^{S_0}}{\partial \tau} (\xi_3 \cdot \tau) ds. \end{aligned}$$

Therefore

$$\underline{\mathbf{B}}_{1,1}^a = \int_{\partial S_0} \left(\frac{\partial \psi_0^{S_0}}{\partial \tau} K_3 - \frac{\partial \psi_0^{S_0}}{\partial n} \xi_3 \cdot \tau \right) K_2 ds + \int_{\partial S_0} \frac{\partial \psi_0^{S_0}}{\partial \tau} \frac{\partial \varphi_2^{S_0}}{\partial \tau} (\xi_3 \cdot \tau) ds.$$

By switching the indexes 2 and 3 we obtain

$$\underline{\mathbf{B}}_{1,1}^b = - \int_{\partial\mathcal{S}_0} \left(\frac{\partial\psi_0^{S_0}}{\partial\tau} K_2 - \frac{\partial\psi_0^{S_0}}{\partial n} \xi_2 \cdot \tau \right) K_3 \, ds - \int_{\partial\mathcal{S}_0} \frac{\partial\psi_0^{S_0}}{\partial\tau} \frac{\partial\varphi_3^{S_0}}{\partial\tau} (\xi_2 \cdot \tau) \, ds.$$

We sum these two terms, observe that $-(\xi_3 \cdot \tau)K_2 + (\xi_2 \cdot \tau)K_3 = K_2^2 + K_3^2 = 1$ and use (8.17c) to get

$$\int_{\partial\mathcal{S}_0} (K_2^2 + K_3^2) \frac{\partial\psi_0^{S_0}}{\partial n} \, ds = \int_{\partial\mathcal{S}_0} \frac{\partial\psi_0^{S_0}}{\partial n} \, ds = 0.$$

We deduce

$$\underline{\mathbf{B}}_{1,1}^a + \underline{\mathbf{B}}_{1,1}^b = \int_{\partial\mathcal{S}_0} \frac{\partial\psi_0^{S_0}}{\partial\tau} \left(\frac{\partial\varphi_2^{S_0}}{\partial\tau} (\xi_3 \cdot \tau) - \frac{\partial\varphi_3^{S_0}}{\partial\tau} (\xi_2 \cdot \tau) \right) \, ds.$$

Now, using (9.29), we obtain

$$\begin{aligned} \underline{\mathbf{B}}_{1,1}^a + \underline{\mathbf{B}}_{1,1}^b &= - \int_{\partial\mathcal{S}_0} R(\vartheta)^t u^\Omega(h) \cdot \left(\frac{\partial\varphi_2^{S_0}}{\partial\tau} (\xi_3 \cdot \tau) - \frac{\partial\varphi_3^{S_0}}{\partial\tau} (\xi_2 \cdot \tau) \right) n \, ds \\ &= R(\vartheta)^t u^\Omega(h) \cdot \int_{\partial\mathcal{S}_0} \left(\frac{\partial\varphi_2^{S_0}}{\partial\tau} (\xi_2 \cdot n) + \frac{\partial\varphi_3^{S_0}}{\partial\tau} (\xi_3 \cdot n) \right) n \, ds. \end{aligned} \quad (9.50)$$

On the other hand we observe that

$$\underline{\mathbf{B}}_{1,1}^c = R(\vartheta)^t u^\Omega(h) \cdot \int_{\partial\mathcal{S}_0} \left(\frac{\partial\varphi_2^{S_0}}{\partial\tau} (\xi_2 \cdot \tau) + \frac{\partial\varphi_3^{S_0}}{\partial\tau} (\xi_3 \cdot \tau) \right) \tau \, ds. \quad (9.51)$$

As a consequence, gathering (9.50) and (9.51) we get

$$\underline{\mathbf{B}}_{1,1} = R(\vartheta)^t u^\Omega(h) \cdot \int_{\partial\mathcal{S}_0} \left[\frac{\partial\varphi_2^{S_0}}{\partial\tau} \xi_2 + \frac{\partial\varphi_3^{S_0}}{\partial\tau} \xi_3 \right] \, ds = 0,$$

as follows by integration by parts.

Let us now consider the second coordinate $\underline{\mathbf{B}}_{1,2}$ of $\underline{\mathbf{B}}_1$, that is:

$$\begin{aligned} \underline{\mathbf{B}}_{1,2} &= \int_{\partial\mathcal{S}_0} \left(\frac{\partial\psi_0^{S_0}}{\partial n} - R(\vartheta)^t u^\Omega(h) \cdot \tau \right) \left(\frac{\partial\varphi_1^{S_0}}{\partial\tau} K_3 - \frac{\partial\varphi_3^{S_0}}{\partial\tau} K_1 \right) \, ds \\ &= \underline{\mathbf{B}}_{1,2}^a + \underline{\mathbf{B}}_{1,2}^b + \underline{\mathbf{B}}_{1,2}^c, \end{aligned}$$

with

$$\begin{aligned} \underline{\mathbf{B}}_{1,2}^a &:= - \int_{\partial\mathcal{S}_0} \frac{\partial\psi_0^{S_0}}{\partial n} \frac{\partial\varphi_3^{S_0}}{\partial\tau} K_1 \, ds, \\ \underline{\mathbf{B}}_{1,2}^b &:= \int_{\partial\mathcal{S}_0} \frac{\partial\psi_0^{S_0}}{\partial n} \frac{\partial\varphi_1^{S_0}}{\partial\tau} K_3 \, ds, \\ \underline{\mathbf{B}}_{1,2}^c &:= \int_{\partial\mathcal{S}_0} R(\vartheta)^t u^\Omega(h) \cdot \left(\frac{\partial\varphi_3^{S_0}}{\partial\tau} K_1 - \frac{\partial\varphi_1^{S_0}}{\partial\tau} K_3 \right) \tau \, ds. \end{aligned}$$

Proceeding as above with $\underline{\mathbf{B}}_{1,1}^a$ and $\underline{\mathbf{B}}_{1,1}^b$ we get

$$\begin{aligned}\underline{\mathbf{B}}_{1,2}^a &= \int_{\partial\mathcal{S}_0} \left(\frac{\partial\psi_0^{\mathcal{S}_0}}{\partial\tau} K_1 - \frac{\partial\psi_0^{\mathcal{S}_0}}{\partial n} \xi_1 \cdot \tau \right) K_3 \, ds + \int_{\partial\mathcal{S}_0} \frac{\partial\psi_0^{\mathcal{S}_0}}{\partial\tau} \frac{\partial\varphi_3^{\mathcal{S}_0}}{\partial\tau} (\xi_1 \cdot \tau) \, ds, \\ \underline{\mathbf{B}}_{1,2}^b &= - \int_{\partial\mathcal{S}_0} \left(\frac{\partial\psi_0^{\mathcal{S}_0}}{\partial\tau} K_3 - \frac{\partial\psi_0^{\mathcal{S}_0}}{\partial n} \xi_3 \cdot \tau \right) K_1 \, ds - \int_{\partial\mathcal{S}_0} \frac{\partial\psi_0^{\mathcal{S}_0}}{\partial\tau} \frac{\partial\varphi_1^{\mathcal{S}_0}}{\partial\tau} (\xi_3 \cdot \tau) \, ds.\end{aligned}$$

We sum these two terms, observing that $(\xi_1 \cdot \tau)K_3 - (\xi_3 \cdot \tau)K_1 = x^\perp \cdot \xi_2$, to get

$$\begin{aligned}\underline{\mathbf{B}}_{1,2}^a + \underline{\mathbf{B}}_{1,2}^b &= - \int_{\partial\mathcal{S}_0} (x^\perp \cdot \xi_2) \frac{\partial\psi_0^{\mathcal{S}_0}}{\partial n} \, ds \\ &\quad + \int_{\partial\mathcal{S}_0} \frac{\partial\psi_0^{\mathcal{S}_0}}{\partial\tau} \left(\frac{\partial\varphi_3^{\mathcal{S}_0}}{\partial\tau} (\xi_1 \cdot \tau) - \frac{\partial\varphi_1^{\mathcal{S}_0}}{\partial\tau} (\xi_3 \cdot \tau) \right) \, ds. \quad (9.52)\end{aligned}$$

Using (9.35) we obtain

$$\int_{\partial\mathcal{S}_0} (x^\perp \cdot \xi_2) \frac{\partial\psi_0^{\mathcal{S}_0}}{\partial n} \, ds = \overline{MR}(\vartheta)^t u^\Omega(h)^\perp \cdot \xi_2, \quad (9.53)$$

with \overline{M} given by (9.34). We also use (9.29) to modify the second term in the right hand side of (9.52) and then get

$$\begin{aligned}\underline{\mathbf{B}}_{1,2}^a + \underline{\mathbf{B}}_{1,2}^b &= -\overline{MR}(\vartheta)^t u^\Omega(h)^\perp \cdot \xi_2 \\ &\quad - \int_{\partial\mathcal{S}_0} R(\vartheta)^t u^\Omega(h) \cdot \left(\frac{\partial\varphi_3^{\mathcal{S}_0}}{\partial\tau} (\xi_1 \cdot \tau) - \frac{\partial\varphi_1^{\mathcal{S}_0}}{\partial\tau} (\xi_3 \cdot \tau) \right) n \, ds.\end{aligned}$$

Adding $\underline{\mathbf{B}}_{1,2}^c$ we get

$$\underline{\mathbf{B}}_{1,2} = -\overline{MR}(\vartheta)^t u^\Omega(h)^\perp \cdot \xi_2 - R(\vartheta)^t u^\Omega(h) \cdot \underline{\mathbf{B}}_{1,2}^d, \quad (9.54)$$

with

$$\begin{aligned}\underline{\mathbf{B}}_{1,2}^d &:= \int_{\partial\mathcal{S}_0} \frac{\partial\varphi_3^{\mathcal{S}_0}}{\partial\tau} \left((\xi_1^\perp \cdot n)n + (\xi_1^\perp \cdot \tau)\tau \right) \, ds \\ &\quad + \int_{\partial\mathcal{S}_0} \frac{\partial\varphi_1^{\mathcal{S}_0}}{\partial\tau} \left(-(\xi_3^\perp \cdot n)n - (\xi_3^\perp \cdot \tau)\tau \right) \, ds \\ &= \int_{\partial\mathcal{S}_0} \frac{\partial\varphi_3^{\mathcal{S}_0}}{\partial\tau} \xi_1^\perp \, ds - \int_{\partial\mathcal{S}_0} \frac{\partial\varphi_1^{\mathcal{S}_0}}{\partial\tau} \xi_3^\perp \, ds.\end{aligned}$$

Using an integration by parts we see that the second term of the right hand side above vanishes and the first term gives

$$\int_{\partial\mathcal{S}_0} \frac{\partial\varphi_3^{\mathcal{S}_0}}{\partial\tau}(x) x^\perp \, ds(x) = - \int_{\partial\mathcal{S}_0} \varphi_3^{\mathcal{S}_0} n \, ds = -\overline{M}\xi_2(0, \cdot).$$

Therefore

$$R(\vartheta)^t u^\Omega(h) \cdot \underline{\mathbf{B}}_{1,2}^d = \overline{M}^t R(\vartheta)^t u^\Omega(h)^\perp \cdot \xi_2. \quad (9.55)$$

Gathering (9.54) and (9.55) and using (9.38) we get

$$\underline{\mathbf{B}}_{1,2} = -2M^\dagger R(\vartheta)^t u^\Omega(h)^\perp \cdot \xi_2. \quad (9.56)$$

Proceeding in the same way for the third coordinate, using

$$(\xi_1 \cdot \tau)K_2 - (\xi_2 \cdot \tau)K_1 = -x^\perp \cdot \xi_3 \text{ and } \int_{\partial \mathcal{S}_0} \frac{\partial \overline{\varphi}_2^{S_0}}{\partial \tau}(x) x^\perp ds(x) = \overline{M} \xi_3(0, \cdot),$$

we get

$$\underline{\mathbf{B}}_{1,3} = -2M^\dagger R(\vartheta)^t u^\Omega(h)^\perp \cdot \xi_3. \quad (9.57)$$

Combining (9.56) and (9.57) and recalling the definition of M_ϑ^\dagger in (9.11), we get (9.49). This ends the proof of Proposition 9.8. \square

9.4 End of the proof of the normal forms

In order to prove Proposition 6.3 and Proposition 7.10 we have to put Equation (2.20) under the normal forms (6.7) and (7.10). We focus on the more delicate Case (ii). Case (i) can be proved with the same strategy with some simplifications, since the proof of the normal form (7.10) corresponding to Case (ii) actually requires to perform additional manipulations. At the end of this section, we add a few words about Proposition 7.17.

Using (2.17c) the equation (2.20) reads

$$M_\varepsilon(q_\varepsilon)q_\varepsilon'' = \gamma^2 E_\varepsilon(q_\varepsilon) + \gamma q_\varepsilon' \times B_\varepsilon(q_\varepsilon) - \langle \Gamma_\varepsilon(q_\varepsilon), q_\varepsilon', q_\varepsilon' \rangle. \quad (9.58)$$

The proof now consists in plugging the previous expansions of Γ_ε , E_ε and B_ε into the right hand side of (9.58) and to rely on some crucial cancellations. Let $\delta > 0$. Using the decomposition (2.27) of Γ , the definition (4.28) of p_ε , the expansions (9.2) and (9.6) for the Christoffel symbols, (9.17) for the electric field, (9.46) for the magnetic field and the relation (9.47) we get, for ε_0 in $(0, 1)$ small enough, as long as $(\varepsilon, q_\varepsilon)$ belongs to $\mathfrak{Q}_{\delta, \varepsilon_0}$:

$$\begin{aligned} \gamma^2 E_\varepsilon(q_\varepsilon) + \gamma q_\varepsilon' \times B_\varepsilon(q_\varepsilon) - \langle \Gamma_\varepsilon(q_\varepsilon), q_\varepsilon', q_\varepsilon' \rangle &= I_\varepsilon \left[\left(\gamma^2 \mathbf{E}_0(q_\varepsilon) + \gamma p_\varepsilon \times B_{\vartheta_\varepsilon}^{S_0} \right) \right. \\ &\quad \left. + \varepsilon \left(\gamma^2 \mathbf{E}_1(q_\varepsilon) + \gamma p_\varepsilon \times \mathbf{B}_1(q_\varepsilon) - \langle \Gamma_{\vartheta_\varepsilon}^{S_0}, p_\varepsilon, p_\varepsilon \rangle \right) + \varepsilon^2 \check{F}_r(\varepsilon, q_\varepsilon, p_\varepsilon) \right], \end{aligned} \quad (9.59)$$

where

$$\begin{aligned} \check{F}_r(\varepsilon, q_\varepsilon, p_\varepsilon) &= \gamma^2 E_r(\varepsilon, q_\varepsilon) + \gamma p_\varepsilon \times B_r(\varepsilon, q_\varepsilon) - \varepsilon \langle \Gamma_r^{\text{rot}}(\varepsilon, q_\varepsilon), p_\varepsilon, p_\varepsilon \rangle \\ &\quad - \varepsilon \langle \Gamma_r^{\partial \Omega}(\varepsilon, q_\varepsilon), p_\varepsilon, p_\varepsilon \rangle. \end{aligned}$$

Then recalling that $F_{\vartheta_\varepsilon}^{\mathcal{S}_0}$ is defined in (4.13), we observe that the zero order term in the right hand side of (9.59) (in terms of powers of ε) can be recast as follows:

$$\gamma^2 \mathbf{E}_0(q_\varepsilon) + \gamma p_\varepsilon \times B_{\vartheta_\varepsilon}^{\mathcal{S}_0}(q_\varepsilon) = F_{\vartheta_\varepsilon}^{\mathcal{S}_0}((\varepsilon \vartheta'_\varepsilon, h'_\varepsilon - \gamma u^\Omega(h_\varepsilon))). \quad (9.60)$$

Now, in order to deal with the subprincipal term of the right hand side of (9.59), let us state the following crucial lemma, where we consider only the part $\mathbf{E}_1^a(q_\varepsilon)$ defined in (9.12) of the decomposition (9.7) of the term $\mathbf{E}_1(q_\varepsilon)$.

Lemma 9.9. *The following holds:*

$$\gamma^2 \mathbf{E}_1^a(q_\varepsilon) + \gamma p_\varepsilon \times \mathbf{B}_1(q_\varepsilon) - \langle \Gamma_{\vartheta}^{\mathcal{S}_0}, p_\varepsilon, p_\varepsilon \rangle = -\langle \Gamma_{\vartheta}^{\mathcal{S}_0}, \hat{p}_\varepsilon, \hat{p}_\varepsilon \rangle. \quad (9.61)$$

where

$$\hat{p}_\varepsilon := (\varepsilon \vartheta'_\varepsilon, h'_\varepsilon - \gamma u^\Omega(h_\varepsilon))^t, \quad (9.62)$$

Remark 9.10. As for (9.60), this relation is algebraic, in the sense that it does not rely on $p_\varepsilon = I_\varepsilon q'_\varepsilon$ or on the fact that q_ε satisfies (2.20).

Proof of Lemma 9.9. We will recast the second and third terms of the left hand side in terms of the matrix M^\dagger defined in (9.8). Let us start with the Christoffel term. Using the definition of $M_{a,\vartheta}^{\mathcal{S}_0}$ in (4.9) and the decomposition of $M_a^{\mathcal{S}_0}$ in (4.23) we get

$$M_{a,\vartheta}^{\mathcal{S}_0} = \begin{pmatrix} m^\# & \mu_{\vartheta}^t \\ \mu_{\vartheta} & M_{b,\vartheta}^{\mathcal{S}_0} \end{pmatrix}, \quad (9.63)$$

with $M_{b,\vartheta}^{\mathcal{S}_0}$ and μ_{ϑ} as in (9.11). In particular we infer from (4.11) that

$$\langle \Gamma_{\vartheta}^{\mathcal{S}_0}, p_\varepsilon, p_\varepsilon \rangle = \left(\begin{array}{c} -(M_{b,\vartheta}^{\mathcal{S}_0} h'_\varepsilon)^\perp \cdot h'_\varepsilon \\ (\varepsilon \vartheta'_\varepsilon)^2 \mu_{\vartheta}^\perp + \varepsilon \vartheta'_\varepsilon ((M_{b,\vartheta}^{\mathcal{S}_0} h'_\varepsilon)^\perp - M_{b,\vartheta}^{\mathcal{S}_0} (h'_\varepsilon)^\perp) \end{array} \right).$$

It remains to recast this expression thanks to the matrix M^\dagger defined in (9.8). This is done thanks to the following elementary identities: recalling (9.9) one has for any ϑ in \mathbb{R} and X in \mathbb{R}^2 ,

$$(M_{b,\vartheta}^{\mathcal{S}_0} X)^\perp \cdot X = X^\perp \cdot M_{\vartheta}^\dagger X^\perp, \quad (9.64)$$

$$(M_{b,\vartheta}^{\mathcal{S}_0} X)^\perp - M_{b,\vartheta}^{\mathcal{S}_0} X^\perp = -2M_{\vartheta}^\dagger X, \quad (9.65)$$

$$(\perp) M_{\vartheta}^\dagger (\perp) = M_{\vartheta}^\dagger. \quad (9.66)$$

Therefore we have:

$$\langle \Gamma_{\vartheta}^{\mathcal{S}_0}, p_\varepsilon, p_\varepsilon \rangle = \left(\begin{array}{c} -(h'_\varepsilon)^\perp \cdot M_{\vartheta}^\dagger (h'_\varepsilon)^\perp \\ (\varepsilon \vartheta'_\varepsilon)^2 \mu_{\vartheta}^\perp - 2\varepsilon \vartheta'_\varepsilon M_{\vartheta}^\dagger h'_\varepsilon \end{array} \right). \quad (9.67)$$

Now using (9.45), (9.66) and the fact that for any ϑ in \mathbb{R} , M_ϑ^\dagger is symmetric, we obtain

$$p_\varepsilon \times \mathbf{B}_I(q_\varepsilon) = \begin{pmatrix} -2(h'_\varepsilon)^\perp \cdot M_{\vartheta_\varepsilon}^\dagger u^\Omega(h_\varepsilon)^\perp \\ -2(\varepsilon\vartheta'_\varepsilon)^2 (M_{\vartheta_\varepsilon}^\dagger u^\Omega(h_\varepsilon)^\perp)^\perp \end{pmatrix} = \begin{pmatrix} -2(h'_\varepsilon)^\perp \cdot M_{\vartheta_\varepsilon}^\dagger u^\Omega(h_\varepsilon)^\perp \\ -2(\varepsilon\vartheta'_\varepsilon)^2 M_{\vartheta_\varepsilon}^\dagger u^\Omega(h_\varepsilon) \end{pmatrix}. \quad (9.68)$$

Now it suffices to combine (9.68), (9.67) and (9.12) to deduce (9.61). \square

As a consequence, combining (9.59), (9.60) and (9.61) we get,

$$\begin{aligned} & \gamma^2 E_\varepsilon(q_\varepsilon) + \gamma q'_\varepsilon \times B_\varepsilon(q_\varepsilon) - \langle \Gamma_\varepsilon(q_\varepsilon), q'_\varepsilon, q'_\varepsilon \rangle \\ &= I_\varepsilon \left[F_{\vartheta_\varepsilon}^{\mathcal{S}_0}(\hat{p}_\varepsilon) + \varepsilon \{ -\langle \Gamma_{\vartheta_\varepsilon}^{\mathcal{S}_0}, \hat{p}_\varepsilon, \hat{p}_\varepsilon \rangle + \gamma^2 (\mathbf{E}_I^b(q_\varepsilon) + \mathbf{E}_I^c(q_\varepsilon)) \} + \varepsilon^2 F_r(\varepsilon, q_\varepsilon, \hat{p}_\varepsilon) \right]. \end{aligned} \quad (9.69)$$

Moreover F_r belongs to $L^\infty(\mathfrak{Q}_{\delta, \varepsilon_0} \times \mathbb{R}^3; \mathbb{R}^3)$, depends on \mathcal{S}_0 , γ and Ω and is weakly nonlinear in the sense of Definition 6.2. Next the part \mathbf{E}_I^c of the subprincipal term in (9.69) can be absorbed by the principal term up to a modification of size ε of the arguments (that is, thanks to the second order modulation). More precisely, considering (4.15) and (9.16), we have

$$F_{\mathcal{S}_0, \vartheta_\varepsilon}(\tilde{p}_\varepsilon) = F_{\mathcal{S}_0, \vartheta_\varepsilon}(\hat{p}_\varepsilon) + \varepsilon \gamma^2 \mathbf{E}_I^c(q_\varepsilon), \quad (9.70)$$

where \tilde{p}_ε is given by (7.8). Thus we deduce from (9.69) and (9.70) that

$$\begin{aligned} & \gamma^2 E_\varepsilon(q_\varepsilon) + \gamma q'_\varepsilon \times B_\varepsilon(q_\varepsilon) - \langle \Gamma_\varepsilon(q_\varepsilon), q'_\varepsilon, q'_\varepsilon \rangle \\ &= I_\varepsilon \left[F_{\vartheta_\varepsilon}^{\mathcal{S}_0}(\tilde{p}_\varepsilon) - \varepsilon \langle \Gamma_{\vartheta_\varepsilon}^{\mathcal{S}_0}, \tilde{p}_\varepsilon, \tilde{p}_\varepsilon \rangle + \varepsilon \gamma^2 \mathbf{E}_I^b(q_\varepsilon) + \varepsilon^2 \hat{F}_r(\varepsilon, q_\varepsilon, \tilde{p}_\varepsilon) \right], \end{aligned} \quad (9.71)$$

where the term \hat{F}_r is defined by

$$\begin{aligned} \hat{F}_r(\varepsilon, q_\varepsilon, \tilde{p}_\varepsilon) &:= \check{F}_r(\varepsilon, q_\varepsilon, \tilde{p}_\varepsilon + \varepsilon \gamma p_c(q_\varepsilon)) - 2\gamma \langle \Gamma_{\vartheta_\varepsilon}^{\mathcal{S}_0}, \tilde{p}_\varepsilon, p_c(q_\varepsilon) \rangle \\ &\quad - \varepsilon \gamma^2 \langle \Gamma_{\vartheta_\varepsilon}^{\mathcal{S}_0}, p_c(q_\varepsilon), p_c(q_\varepsilon) \rangle, \end{aligned}$$

where $p_c(q_\varepsilon) := (0, u_c(q_\varepsilon))$. One can easily check that \hat{F}_r is still weakly nonlinear.

Using Proposition 9.1 and (9.71), and recalling the notation (4.32), the equation (9.58) can now be recast as follows:

$$\begin{aligned} \varepsilon^{\min(2, \alpha)} \left(M_{\vartheta_\varepsilon}(\varepsilon) + \varepsilon^{4-\min(2, \alpha)} M_r(\varepsilon, q_\varepsilon) \right) p'_\varepsilon &= F_{\vartheta_\varepsilon}^{\mathcal{S}_0}(\tilde{p}_\varepsilon) - \varepsilon \langle \Gamma_{\vartheta_\varepsilon}^{\mathcal{S}_0}, \tilde{p}_\varepsilon, \tilde{p}_\varepsilon \rangle \\ &\quad + \varepsilon \gamma^2 \mathbf{E}_I^b(q_\varepsilon) + \varepsilon^2 \hat{F}_r(\varepsilon, q_\varepsilon, \tilde{p}_\varepsilon). \end{aligned} \quad (9.72)$$

We need to perform further modifications on this equation in order to achieve the normal forms (6.7)-(7.10) exactly, due to the fact that the mass matrix in (9.72) contains some extra lower-order terms, and that the time derivative

is applied to p_ε rather than to \tilde{p}_ε . In order to deal with the first discrepancy, reducing ε_0 in $(0, 1)$ if necessary, we simply multiply (9.72) by the matrix

$$M_\vartheta(\varepsilon) \left(M_\vartheta(\varepsilon) + \varepsilon^{4-\min(2,\alpha)} M_r(\varepsilon, q) \right)^{-1}.$$

On the other hand, for the second discrepancy, denoting $\tilde{\ell}_\varepsilon := h' - \gamma(u^\Omega(h_\varepsilon) + \varepsilon u_c(q_\varepsilon))$ we compute

$$\begin{aligned} \tilde{\ell}'_\varepsilon &= h''_\varepsilon - \gamma \tilde{\ell}_\varepsilon \cdot (\nabla u^\Omega)(h_\varepsilon) - \gamma^2 (u^\Omega(h_\varepsilon) + \varepsilon u_c(q_\varepsilon)) \cdot \nabla u^\Omega(h_\varepsilon) \\ &\quad - \gamma D_\vartheta u_c(q_\varepsilon) \varepsilon \vartheta'_\varepsilon - \varepsilon \gamma D_h u_c(q_\varepsilon) \cdot (\tilde{\ell}_\varepsilon + \gamma(u^\Omega(h_\varepsilon) + \varepsilon u_c(q_\varepsilon))). \end{aligned}$$

Thus we obtain (7.10) with F_r in $L^\infty(\mathfrak{Q}_{\delta, \varepsilon_0} \times \mathbb{R}^3; \mathbb{R}^3)$ weakly nonlinear in the sense of Definition 6.2. This ends the proof of Proposition 7.10.

Starting from (2.34a), the proof of Proposition 7.17 is similar to the one of Proposition 7.10, with some simplifications, since in this case $\tilde{M}_b^{S_0} = \pi \text{Id}_2$ and consequently $M^\dagger = 0$. It follows that $\mathbf{E}_I^a = 0$ and $\mathbf{B}_I = 0$.

Now we expand $\tilde{M}_{b,\varepsilon}$, $\Gamma_{b,\varepsilon}$, $E_{b,\varepsilon}$ and $\tilde{B}_{1,\varepsilon}$ (which depend merely on h_c and ε) in terms of ε . Note that the two last coordinates of \mathbf{E}_I^b (which are the only ones to be relevant here, recall (2.32)) are zero and that \mathbf{E}_I^c gives the term $-u_c(q)^\perp$. Noting that $u^\Omega(h) + \varepsilon u_c(q) = u^\Omega(h_c) + O(\varepsilon^2)$ and recalling (2.33) we infer

$$E_{b,\varepsilon} = -u^\Omega(h_c)^\perp + \varepsilon^2 E_{b,r} \text{ and } \tilde{B}_{1,\varepsilon} = -1 + \varepsilon^2 B_{b,r},$$

with $E_{b,r} = E_{b,r}(\varepsilon, h_{c,\varepsilon})$ weakly nonlinear in the sense of Definition 7.16 and $B_{b,r} = B_{b,r}(\varepsilon, h_{c,\varepsilon})$ bounded as long as $h_{c,\varepsilon}$ is away from $\partial\Omega$. On the other side, one finds that $\tilde{M}_b = \varepsilon^2 \pi \text{Id}_2 + O(\varepsilon^4)$ and $\Gamma_{b,\varepsilon} = O(\varepsilon^4)$. The conclusion follows easily and this ends the proof of Proposition 7.17.

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