



HAL
open science

Derivation of nonlinear shell models combining shear and flexure: application to biological membranes

Olivier Pantz, Karim Trabelsi

► **To cite this version:**

Olivier Pantz, Karim Trabelsi. Derivation of nonlinear shell models combining shear and flexure: application to biological membranes. *Mathematics and Mechanics of Complex Systems*, 2015, 3 (2), pp.101–138. hal-01110273

HAL Id: hal-01110273

<https://inria.hal.science/hal-01110273>

Submitted on 27 Jan 2015

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

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

**DERIVATION OF NONLINEAR SHELL MODELS COMBINING
SHEAR AND FLEXURE. APPLICATION TO BIOLOGICAL
MEMBRANES.**

OLIVIER PANTZ AND KARIM TRABELSI

ABSTRACT. Biological membranes are often idealized as incompressible elastic surfaces whose strain energy only depends on their mean curvature and possibly on their shear. We show that this type of model can be derived using a formal asymptotic method by considering biological membranes to be thin, strongly anisotropic, elastic, locally homogeneous bodies.

The authors would like to thank the referees for their valuable comments which helped to improve the manuscript.

This work was partly supported by the ANR Geometrya.

O. Pantz is member of the DEFI team at INRIA Saclay Ile-de-France.

1. INTRODUCTION

Shells, plates and membranes are solid deformable bodies having one characteristic dimension small by comparison with the other two dimensions. Their behavior is fully described by standard three-dimensional laws of continuum mechanics. Nevertheless, it is tempting, at least from the modeling viewpoint, to consider them as two-dimensional structures and to replace the genuine mechanical laws by two-dimensional reduced versions. This immediately raises two questions: (1) What is the correct model? and (2) How can it be mathematically justified? To this end, we consider the thickness ε of the plate/shell/membrane as a parameter and identify the limit behavior of the structure as ε goes to zero. According to the dependence of the elasticity moduli on the thickness of the shell, a full zoology of models may be derived. Membrane, isometric bending and von Kármán theories have been justified (amongst others), first formally (see Fox, Raoult and Simo [Fox et al. \[1993\]](#)), then by means of Γ -convergence (see Le Dret and Raoult [Le Dret and Raoult \[1995, 1996\]](#), Pantz [Pantz \[2003\]](#), Müller, Friesecke and James [Friesecke et al. \[2002\]](#), Friesecke, James and Mora [Friesecke et al. \[2003\]](#), see also [Conti et al. \[2006\]](#)). In those works, elasticity coefficients are assumed to scale like a power of the thickness ε of the plate or shell, that is, like $\varepsilon^{-\alpha}$. Membrane theory corresponds to the case $\alpha = 1$, isometric bending to the case $\alpha = 3$ and von Kármán to $\alpha = 4$. Intermediate values of α have also been considered, and an almost exhaustive hierarchy of models has thus been produced (see Müller, Friesecke and James [Friesecke et al. \[2006\]](#)). Some cases remain to be treated, Conti and Maggi [Conti and Maggi \[2008\]](#), for instance, investigate the scaling of the energy corresponding to folds. The initial motivation for this article was the study of the mechanical behavior of Red Blood Cells (RBCs), and our aim was to determine whether the classical RBC model could be derived by the above procedure.

The mature anucleate RBCs¹ are made of two mechanical structures: The cytoskeleton – a two-dimensional network of protein filaments that extends throughout the interior of the cell – and a lipid bilayer. Both are bound together by proteins linking the nodes of the mesh of the cytoskeleton to the lipid bilayer via transmembrane proteins. Lipid bilayers are self-assembled structures of phospholipids which are small molecules containing a negatively charged phosphate group (called the head), and two highly hydrophobic fatty acid chains (called the tails). In an aqueous environment, phospholipids spontaneously form a double layer whose configuration enables to isolate the hydrophobic tails from the surrounding water molecules. Modifying the area of such a lipid bilayer is energy-costly because it exposes some of the tails to the environment.

A bilayer that supports no other mechanical structure, which is connected and has no boundary is called a vesicle. Vesicles are massively studied because they are easy to obtain experimentally. Moreover, they partially mimic the behavior of RBCs. Roughly speaking, they are RBCs without cytoskeleton (even if the RBC bilayer does embed a lot of different proteins responsible for different functions of the cell). They similarly resist to bending. However, a vesicle shows no resistance to shear stress, contrarily to RBCs owing to their cytoskeleton.

¹Every mention to RBCs in this article will implicitly refer to anucleate mature RBCs without further notice.

A widely used model consists in considering that a lipid bilayer may be endowed with an elastic energy depending solely on the mean curvature of the vesicle usually known as the Helfrich functional (named after Willmore in other contexts). It has been introduced, as far as we know independently, by Canham [Canham \[1970\]](#) and Helfrich [Helfrich \[1973\]](#) some forty years ago. Evans [Evans \[1974\]](#) has shown that the Helfrich functional can be derived by assuming a vesicle to be made of two interconnected elastic fluid membranes, each of them resisting to change of local area but not to bending itself. Jenkins [Jenkins \[1977a\]](#) has extended the analysis of Helfrich to general two-dimensional liquid crystals [Singer and Nicolson \[1972\]](#). In particular, he derives the Euler-Lagrange equations satisfied by the equilibrium states, and examines the consequences of fluidity on the form of the strain energy (see also [Steigmann \[1999\]](#)). As a means to take into account the various vesicle shapes observed, it is common to presume that the vesicle is endowed with a nonzero spontaneous curvature. The origin of this spontaneous curvature is usually attributed to different compositions of the outer and inner layers. Several refinements to this basic model have since been proposed as the so-called bilayer-couple model [Svetina and ek \[1989\]](#), that consists in allowing the two lipidic layers to slip on one another, and imposing that the total area of each layer remains constant (see also [Seifert et al. \[1991\]](#) for a comparison between the two models). Miao, Seifert, Wortis and Döbereiner [Miao et al. \[1994\]](#) proposed an intermediate model called area-difference elasticity model, where slight total area changes of each layer are allowed but still penalized.

As previously mentioned, the mechanical structure of the RBC is not only imputable to its bilayers. Their cytoskeleton endows them with resistance to shear stress. In most models, only the deformation of the RBC membrane is considered (that is of the bilayer). To take into account the presence of the cytoskeleton an additional term is added to the the total energy depending on the change of the metric of the membrane. Krishnaswamy [Krishnaswamy \[1996\]](#) proposed another model for which the deformations of the cytoskeleton and the fluid bilayer may differ.

Finally, the aforementioned models for vesicles and RBCs are backed up by numerous numerical studies that allow to reproduce various shapes observed experimentally. Amongst others, Deuling and Helfrich [Deuling, H.J. and Helfrich, W. \[1976\]](#) (see also Jenkins [Jenkins \[1977b\]](#) and Luke [Luke \[1982\]](#), [Luke and Kaplan \[1979\]](#)) have computed axisymmetric vesicle shapes of minimum energy with respect to the values of the reduced volume and spontaneous curvature. Seifert, Berndt and Lipowsky [Seifert et al. \[1991\]](#) have compared the axisymmetric solutions obtained using the spontaneous curvature model and the bilayer-couple model, whereas Agrawal and Steigmann [Agrawal and Steigmann \[2009\]](#) have included contact conditions between the vesicle and a substrate. Full three-dimensional simulations have been performed by Feng and Klug [Feng and Klug \[2006\]](#), Bonito, Nochetto and Pauletti [Bonito et al. \[2010, 2011\]](#), Dziuk [Dziuk \[2008\]](#) using a finite element method. Peng et al [Peng et al. \[2013\]](#) use a dissipative particle dynamic approach and focus on the interaction between the lipid bilayer and the cytoskeleton. Du, Chun and Xiaoqiang perform numerical computations based on a Phase Field Method [Du et al. \[2006, 2004\]](#), [Du and Zhang \[2008\]](#). Boundary Integral Methods have been used by Veerapaneni, Gueyffier and Zorin [Veerapaneni et al.](#)

[2009], Sohn, Tseng, Li, Voigt and Lowengrub [Sohn et al. \[2010\]](#). Another approach based on the Immersed Boundary Method has been investigated by Kim and Lai [Kim and Lai \[2010\]](#), Liu et al [Liu et al. \[2004\]](#), [Liu and Liu \[2006\]](#) and, together with a Lattice Boltzmann approach by Crowl and Fogleson [Crowl and Fogleson \[2010\]](#). Finally, Level Set Methods have also been implemented in this context by Salac and Miksis [Salac and Miksis \[2011\]](#), and Maitre, Milcent, Cottet, Raoult and Usson [Maitre et al. \[2009\]](#) (see also Doyeux et al. [Doyeux et al. \[2013\]](#)).

We prove in this article that the classical mechanical model of the RBCs can be recovered by means of a formal asymptotic analysis assuming that the RBCs' membrane is made of a locally homogeneous, albeit strongly anisotropic, non-linearly elastic material. The main difference with previous works on the justification of thin structures is that we assume different scalings for the elastic moduli in the tangential and normal directions to the midsection. Let us underline that our work cannot be considered as a justification of the classical RBC mechanical model. Indeed, the RBC is not a locally homogeneous elastic membrane. Firstly because it is made of two different structures: The lipid bilayer (responsible for the resistance to bending) and a cytoskeleton (responsible for resistance to shear). Even the lipid bilayer could hardly be considered as made of a homogeneous material, the scale of the phospholipids it contains being of the same order as the thickness of the membrane. The cytoskeleton, being a two-dimensional spectrin network, is no more a homogeneous elastic body. Even if it is not overt at first glance, our work is strongly related to the justification, already mentioned, proposed by Evans [Evans \[1974\]](#).

We have chosen to consider a rather general setting (presented in section 2) for which the modeling of the RBCs is obtained as a particular case (see section 6). The asymptotic analysis is performed in section 3. Assuming that the minimizers of the energy admit an asymptotic expansion with respect to the thickness (section 3.2), they converge toward the solutions of a two-dimensional problem (see section 3.3). The limit energy, computed in section 3, contains membrane and flexural terms. In section 4, we prove that under invariance assumptions on the stored energy of the material, the flexural term depends only upon the second fundamental form, or even only upon the mean curvature of the shell. The isometric bending shell, RBC and vesicle models are obtained as particular applications in section 6. The last section is devoted to some general remarks in particular on the relaxation of the formal energy limit.

Finally, let us specify some notations. If M is a differentiable manifold, we denote by TM and T^*M its tangent and cotangent bundles. Moreover, $T^*(M; \mathbb{R}^3)$ will stand for the triple Whitney sum $T^*M \oplus T^*M \oplus T^*M$. The tangent spaces of a product of manifolds will be implicitly identified with the product of the tangent spaces: so that if M_1 and M_2 are differentiable manifolds and $M = M_1 \times M_2$, the bundle TM will be implicitly identified with $TM_1 \times TM_2$. If M is an open subset of \mathbb{R}^N , TM will be identified with $M \times \mathbb{R}^N$. The corresponding identifications will also be made for T^*M and $T^*(M; \mathbb{R}^3)$. The set of reals \mathbb{R} and its dual \mathbb{R}' will also be often implicitly identified. Sets will always be displayed in capital letter (for instance the set of deformations ψ^ε will be denoted Ψ^ε). Sequences of terms of an asymptotic expansion are denoted using bold letters (for instance $\boldsymbol{\psi} = (\psi_k)_{k \in \mathbb{N}}$ stands for the asymptotic expansion of ψ^ε). Accordingly, the sets of asymptotic

expansions used both bold and capitalized letters (for instance $\boldsymbol{\psi} \in \boldsymbol{\Psi}$). Moreover, calligraphic letters will be exclusively used for fiber spaces. Two different reference configurations are used throughout our article, one is qualified to be abstract and the other geometrical. The same notations are used for both configurations, the only distinction being that a tilde is added over variables, sets and functionals defined on the geometric configuration (for instance $\tilde{\psi}^\varepsilon$ is the deformation defined on the geometric configuration, whereas ψ^ε stands for the deformation over the abstract one). All the notations introduced are recalled at the end of the article for convenience.

2. ELASTIC SHELLS – THREE DIMENSIONAL MODELING

We consider a thin nonlinearly elastic shell of midsurface S' and constant half thickness $\varepsilon > 0$, and choose $S^\varepsilon = S' \times (-\varepsilon, \varepsilon)$ to be the reference configuration of this elastic body. We assume S' to be a regular two-dimensional orientable submanifold of \mathbb{R}^3 with or without boundary. In the following, S' is implicitly endowed with the metric induced by the Euclidean metric in \mathbb{R}^3 . Let ψ^ε be the deformation of the shell, that is, a map from S^ε into \mathbb{R}^3 . The differential $D\psi^\varepsilon(x^\varepsilon)$ of ψ^ε at x^ε is a linear map from $T_{x^\varepsilon}S^\varepsilon$ into $T_{\psi^\varepsilon(x^\varepsilon)}\mathbb{R}^3$. Since $T_{\psi^\varepsilon(x^\varepsilon)}\mathbb{R}^3$ is canonically isomorphic to \mathbb{R}^3 , $D\psi^\varepsilon(x^\varepsilon)$ is identified with an element of the Whitney sum $T^*S^\varepsilon \oplus T^*S^\varepsilon \oplus T^*S^\varepsilon$ denoted by $T^*(S^\varepsilon; \mathbb{R}^3)$. We denote by $J_\varepsilon(\psi^\varepsilon)$ the elastic energy of the shell under the deformation ψ^ε . We assume that the elastic energy is local and depends only on the first derivatives of the deformation. In other words, there exists a map W^ε from $T^*(S^\varepsilon; \mathbb{R}^3)$ into $\overline{\mathbb{R}}^+$ such that

$$J_\varepsilon(\psi^\varepsilon) := \int_{S^\varepsilon} W^\varepsilon(D\psi^\varepsilon) dx^\varepsilon,$$

where $dx^\varepsilon = dx' \wedge dx_3^\varepsilon$ and dx' is the two-dimensional Hausdorff measure restricted to S' , whereas $D\psi^\varepsilon(x^\varepsilon)$ stands for the differential of ψ^ε at $x^\varepsilon \in S^\varepsilon$. Note that this representation enables us to consider inhomogeneous shells. The shell is assumed to be submitted to volumic dead body loads $f_\varepsilon \in L^2(S^\varepsilon)^3$, and we set

$$L_\varepsilon(\psi^\varepsilon) := \int_{S^\varepsilon} f_\varepsilon \cdot \psi^\varepsilon dx^\varepsilon.$$

The total energy of the system is accordingly given by

$$I_\varepsilon(\psi^\varepsilon) := J_\varepsilon(\psi^\varepsilon) - L_\varepsilon(\psi^\varepsilon).$$

Finally, boundary conditions may also be added. We set $\Gamma^\varepsilon = \gamma \times (-\varepsilon, \varepsilon)$, where $\gamma \subset \partial S'$ is the – possibly empty – part of the boundary where the shell is clamped, and by ϕ^ε the imposed deformation on this set. Our aim is to determine the behavior of the minimizers φ^ε of I_ε over

$$\Psi^\varepsilon := \{ \psi^\varepsilon \in W^{1,\infty}(S^\varepsilon)^3 \text{ such that } \psi^\varepsilon(x^\varepsilon) = \phi^\varepsilon(x^\varepsilon) \text{ for every } x^\varepsilon \in \Gamma^\varepsilon \}$$

as ε goes to zero. Note that the minimization problem of I_ε over Ψ^ε without any growth and polyconvex or quasiconvex assumptions on the stored energy function is generally not well posed. Here, we implicitly assume this problem to have a regular solution. Various assumptions have to be made regarding the dependence of the energy on the thickness for the needs of our analysis. These mainly concern the stored energy W^ε (see 2.1), but also the applied loads (see 2.2).

2.1. Dependence of the stored energy functions with respect to the thickness. We set $S = S^1$, and assume that the stored energy W^ε to be of the form

$$(1) \quad W^\varepsilon(F) = \varepsilon^{-1} (\varepsilon^{-2} W_2(F) + W_0(F)),$$

for every $F \in T^*(S^\varepsilon; \mathbb{R}^3)$ and $\varepsilon \leq 1$ where W_0 and W_2 are continuous nonnegative maps from $T^*(S; \mathbb{R}^3)$ into $\overline{\mathbb{R}}^+$. Note that we implicitly use the injection of $T^*(S^\varepsilon; \mathbb{R}^3) = T^*(S' \times (-\varepsilon, \varepsilon); \mathbb{R}^3) = T^*(S'; \mathbb{R}^3) \times T^*((-\varepsilon, \varepsilon); \mathbb{R}^3)$ into $T^*(S; \mathbb{R}^3) = T^*(S' \times (-1, 1); \mathbb{R}^3) = T^*(S'; \mathbb{R}^3) \times T^*((-1, 1); \mathbb{R}^3)$ in the definition (1) of W^ε . Standard analysis focuses on the case where only one element of this expansion is not zero. For instance, if $W_2 = 0$, we recover a nonlinear membrane model [Le Dret](#)

and Raoult [1996], and if $W_0 = 0$ we obtain the isometric bending one Frießecke et al. [2003].

Behavior of strongly extended fibers. We assume that the stored energy W_2 is bounded from below by a positive constant for strongly extended fibers, namely, there exists $\delta, c > 0$ such that

$$(2) \quad \forall F' \in T^*(S'; \mathbb{R}^3), \forall F_3 \in T^*((-1, 1); \mathbb{R}^3) \text{ such that } |F_3| \geq \delta, \\ W_2(F', F_3) \geq c.$$

Note that for every element F of a vector bundle endowed with a Riemann metric, the notation $|F|$ should be understood as the norm of the vectorial part of F . In particular, in the statement (2), $|F_3| = |v|$ if $F_3 = (x_3, v) \in (-1, 1) \times \mathbb{R}^3 = T^*((-1, 1); \mathbb{R}^3)$.

Regularity and zero set of W_2 . We assume that W_2 is a C^2 -nonnegative function and denote by \mathcal{M} the restriction of its zero set to the midsection, that is,

$$\mathcal{M} := \{F \in T^*(S'; \mathbb{R}^3) \times T_0^*((-1, 1); \mathbb{R}^3) \text{ such that } W_2(F) = 0\}.$$

Let \mathcal{M}' be the projection of \mathcal{M} onto $T^*(S'; \mathbb{R}^3)$, that is,

$$\mathcal{M}' := \{F' \in T^*(S'; \mathbb{R}^3) : \text{there exists } n_0 \in T_0^*((-1, 1); \mathbb{R}^3) \text{ with } (F', n_0) \in \mathcal{M}\}.$$

We assume that the projection of \mathcal{M} onto \mathcal{M}' is one to one. We denote by $n_0 : \mathcal{M}' \rightarrow T_0^*((-1, 1); \mathbb{R}^3)$ the function that maps the elements F' of \mathcal{M}' to the corresponding element F_3 of $T_0^*((-1, 1); \mathbb{R}^3)$, so that

$$(3) \quad \mathcal{M} = \{(F', n_0(F')) \in T^*(S; \mathbb{R}^3) : F' \in \mathcal{M}'\}.$$

We recall that $S = S' \times (-1, 1)$ and that $T^*(S; \mathbb{R}^3)$ is identified with $T^*(S'; \mathbb{R}^3) \times T^*((-1, 1); \mathbb{R}^3)$. The vectorial part of $n_0(F') \in T_0^*((-1, 1); \mathbb{R}^3) = \{0\} \times \mathbb{R}^3$ will be noted $n(F')$, so that $n_0(F') = (0, n(F'))$.

Local interpenetration. To avoid local interpenetration of matter, it is geometric to expect $D\psi^\varepsilon$ to be invertible. To this end, we require that $W_0(F) = \infty$ for every $F \in \mathcal{M}$ such that $\det F < 0$ and that

$$(4) \quad W_0(F) \rightarrow \infty \text{ if } F \in \mathcal{M} \text{ and } \det(F) \rightarrow 0.$$

2.2. Dependence of the applied loads with respect to the thickness. The volumic loads are assumed to scale as the inverse of the thickness of the shell and more precisely that there exists $f : S \rightarrow \mathbb{R}^3$ such that for every $\varepsilon \leq 1$,

$$(5) \quad f_\varepsilon(x) = \varepsilon^{-1} f(x) \text{ for every } x \in S^\varepsilon.$$

3. FROM 3D TO 2D – A FORMAL ASYMPTOTIC ANALYSIS

3.1. Rescaling. We set $\psi(\varepsilon)(x', x_3) = \psi^\varepsilon(x', \varepsilon x_3)$. Moreover, we define the rescaled energies

$$J(\varepsilon)(\psi(\varepsilon)) := J_\varepsilon(\psi^\varepsilon) \quad \text{and} \quad I(\varepsilon)(\psi(\varepsilon)) := I_\varepsilon(\psi^\varepsilon).$$

The minimization problem of I_ε over Ψ^ε is then equivalent to the minimization problem of $I(\varepsilon)$ over

$$\Psi(\varepsilon) := \{ \psi(\varepsilon) \in W^{1,\infty}(S)^3 \text{ such that } \psi(\varepsilon)(x) = \phi(\varepsilon)(x) \text{ for every } x \in \Gamma \},$$

where $\phi(\varepsilon)(x) = \phi^\varepsilon(x', \varepsilon x_3)$.

For every map $\psi^\varepsilon : S^\varepsilon \rightarrow \mathbb{R}^3$, we denote by $(D'\psi^\varepsilon, D_3\psi^\varepsilon)$ the decomposition of the differential ψ^ε along the sections of the cylinder S^ε and along its fibers respectively. In other words, for every $x^\varepsilon = (x', x_3) \in S^\varepsilon$, $D'\psi^\varepsilon(x^\varepsilon)$ and $D_3\psi^\varepsilon(x^\varepsilon)$ stand for the elements of $T_{x'}^*(S'; \mathbb{R}^3)$ and $T_{x_3}^*((-1, 1); \mathbb{R}^3)$ such that $D\psi^\varepsilon(x^\varepsilon) = (D'\psi^\varepsilon(x^\varepsilon), D_3\psi^\varepsilon(x^\varepsilon))$.

For every deformation $\psi(\varepsilon)$ of S , we define its partial derivation $\partial_3\psi(\varepsilon)$ with respect to the normal direction as

$$\partial_3\psi(\varepsilon)(x', x_3) = \lim_{t \rightarrow 0} \frac{\psi(\varepsilon)(x', x_3 + t) - \psi(\varepsilon)(x', x_3)}{t}.$$

Performing a simple change of variable, we get

$$\begin{aligned} J(\varepsilon)(\psi(\varepsilon)) &= \varepsilon^{-1} \int_{S^\varepsilon} (\varepsilon^{-2}W_2 + W_0) (D'\psi^\varepsilon(x^\varepsilon), D_3\psi^\varepsilon(x^\varepsilon)) dx^\varepsilon \\ &= \varepsilon^{-1} \int_{S^\varepsilon} (\varepsilon^{-2}W_2 + W_0) (D'\psi^\varepsilon(x^\varepsilon), (x_3^\varepsilon, \partial_3\psi^\varepsilon(x^\varepsilon))) dx^\varepsilon \\ &= \int_S (\varepsilon^{-2}W_2 + W_0) (D'\psi(\varepsilon)(x), (\varepsilon x_3, \varepsilon^{-1}\partial_3\psi(\varepsilon)(x))) dx. \end{aligned}$$

3.2. Ansatz. In order to perform our formal analysis, we assume that the minimizers $\varphi(\varepsilon)(x', x_3) = \varphi^\varepsilon(x', \varepsilon x_3)$ of the energy admit an asymptotic expansion

$$(6) \quad \varphi(\varepsilon)(x) = \sum_{k \geq 0} \varepsilon^k \varphi_k(x) \text{ for every } x \in S,$$

with $(\varphi_k) \in \ell^1(W^{1,\infty}(S)^3)$. Obviously, the same assumption has to be made on the applied Dirichlet boundary condition and we let $\phi = (\phi_k) \in \ell^1(W^{1,\infty}(S)^3)$ be the terms of the asymptotic expansion of the deformation $\phi(\varepsilon)(x) = \phi^\varepsilon(x', \varepsilon x_3)$ imposed on $\Gamma := \gamma \times (-1, 1)$, that is

$$(7) \quad \phi(\varepsilon)(x) = \sum_{k \geq 0} \varepsilon^k \phi_k(x) \text{ for every } x \in S.$$

The condition $\varphi^\varepsilon \in \Psi^\varepsilon$ reads as $\varphi_k(x) = \phi_k(x)$ for every $x \in \Gamma$. Consequently, we introduce the admissible set

$$\Psi := \{ \psi = (\psi_k) \in \ell^1(W^{1,\infty}(S)^3) \text{ such that } \psi_k = \phi_k \text{ for every } x \in \Gamma \},$$

and the rescaled energies $\mathbf{J}(\varepsilon)$ and $\mathbf{I}(\varepsilon)$ from Ψ into $\overline{\mathbb{R}}$ defined by

$$(8) \quad \mathbf{J}(\varepsilon)(\psi) := J(\varepsilon) \left(\sum_{k \geq 0} \varepsilon^k \psi_k \right) \quad \text{and} \quad \mathbf{I}(\varepsilon)(\psi) := I(\varepsilon) \left(\sum_{k \geq 0} \varepsilon^k \psi_k \right).$$

3.3. Limit of the total energy. The first step of our analysis consists in computing the limit of $\mathbf{J}(\varepsilon)(\psi)$ as ε goes to zero for $\psi \in \Psi$. As we shall see in Proposition 1, the limit of $\mathbf{J}(\varepsilon)$ contains two terms. Roughly speaking, one term measures the elastic energy due to the change of the metric of the midsection of the shell. It depends only on W_0 . The second term measures the elastic energy due to the variations of the orientation of its fibers. It depends on the second derivative of the stored energy function W_2 through a quadratic form $Q_{D'\psi_0}$.

In order to enhance the readability of the sequel, we introduce a practical notation. We recall that a section F of a vector bundle \mathcal{F} is a map from its base into \mathcal{F} such that $\pi_B(F)$ is the identity, where π_B stands for the projection of \mathcal{F} onto its base B . Given such a section, we define the bundle map

$$(9) \quad \mathcal{F} \rightarrow T\mathcal{F}, \quad G \mapsto G_F = \frac{d}{dt}(F(\pi_B(G)) + tG)|_{t=0}.$$

Roughly speaking, G_F is the element G of $T_{F(\pi_B(G))}\mathcal{F}$. Similarly, for every $(x, v) \in \mathbb{R}^N \times \mathbb{R}^N$, we will sometimes denote $(x, v) \in T_x\mathbb{R}^N$ by v_x . Let F' be a section of \mathcal{M}' , we set for every $(G', s, v) \in T^*(S'; \mathbb{R}^3) \times \mathbb{R} \times (\mathbb{R}')^3$,

$$(10) \quad Q_{F'}(G', s, v) := D^2W_2[G'_{F'}, s_0, v_{n(F')}]^2,$$

where $(G'_{F'}, s_0, v_{n(F')})$ is the element of $T_{(F', n_0(F'))}(T^*(S; \mathbb{R}^3))$ defined in (9) based on the decomposition of $T^*(S; \mathbb{R}^3) = T^*(S'; \mathbb{R}^3) \times (-1, 1) \times (\mathbb{R}')^3$, while D^2W_2 stands for the Hessian of W_2 . Namely, we have

$$(11) \quad (G'_{F'}, s_0, v_{n(F')}) = \frac{d\gamma}{dt}(0) \quad \text{where} \quad \gamma(t) = (F'_{\pi_{S'}(G')} + tG', ts, n(F') + tv).$$

At first glance, the meaning of $D^2W_2[\dot{\gamma}(0)]^2$ is unclear, considering that the Hessian of a map defined on a manifold is not, in general, intrinsically defined. Nevertheless, it is well known that this is consistent on the set of critical points, which is precisely what is considered here. Indeed, $\gamma(0)$ is equal to the value of the section $(F', n_0(F'))$ at $\pi_{S'}(G)$. Yet, F' is a section of \mathcal{M}' , hence $W_2(F', n_0(F)) = 0$, $W_2(\gamma(0)) = 0$ and $DW_2(\gamma(0)) = 0$. As a result, $D^2W_2[\dot{\gamma}(0)]^2$ is well defined and, accordingly,

$$(12) \quad D^2W_2[\dot{\gamma}(0)]^2 = 2 \lim_{t \rightarrow 0} t^{-2} W_2(\gamma(t)).$$

Note that the right-hand side of (12) only depends on $\dot{\gamma}(0)$, so that the particular choice of the representative $\gamma(t)$ of $\dot{\gamma}(0)$ is irrelevant as already mentioned.

We are now in a position to state the main result of this section.

Proposition 1. *Let Φ be the subset of the admissible set Ψ defined by*

$$(13) \quad \Phi := \{\psi \in \Psi : \partial_3\psi_0 = 0, D'\psi_0(x) \in \mathcal{M}' \text{ and } \partial_3\psi_1(x) = n(D'\psi_0(x)) \text{ for every } x \in S\}.$$

Let $\psi \in \Psi$, then

$$\lim_{\varepsilon \rightarrow 0} \mathbf{I}(\varepsilon)(\psi) = \begin{cases} I_0\left(\psi_0, \frac{1}{2} \int_{-1}^1 \psi_1 dx_3, \partial_3\psi_2\right) & \text{if } \psi \in \Phi, \\ +\infty & \text{if } \psi \notin \Phi, \end{cases}$$

where

$$I_0(\psi_0, u, v) := J_0(\psi_0, u, v) - 2 \int_{S'} f_0 \cdot \psi_0 dx,$$

$$J_0(\psi_0, u, v) := \frac{1}{2} \int_S Q_{D'\psi_0}(D'u + x_3 D'n, x_3, v) dx, + 2 \int_{S'} W_0(D'\psi_0, n_0) dx',$$

where n and n_0 stand for $n(D'\psi_0)$ and $n_0(D'\psi_0)$ for short, $Q_{D'\psi_0}$ is defined by (10) and $f_0(x') = f(x', 0)$.

Proof. We proceed in two steps. First, we prove that every sequence of deformations $\psi \in \Psi$ of finite elastic energy, namely which satisfies

$$\liminf_{\varepsilon \rightarrow 0} \mathbf{J}(\varepsilon)(\psi) < +\infty,$$

belongs to Φ . In particular, this implies that $\mathbf{J}(\varepsilon)(\psi)$ converges to infinity as ε goes to zero, for every ψ that is not in Φ . In a second step, we compute the limit of $\mathbf{J}(\varepsilon)(\psi)$ for every ψ in Φ .

Let $\psi \in \Psi$ be the asymptotic expansion of a deformation of finite elastic energy. From Fatou's lemma, we deduce

$$\begin{aligned} & \int_S \liminf_{\varepsilon \rightarrow 0} (\varepsilon^{-2} W_2) \left(\sum_{k \geq 0} \varepsilon^k (D'\psi_k, (\varepsilon x_3, \varepsilon^{-1} \partial_3 \psi_k)) \right) dx \\ & \leq \liminf_{\varepsilon \rightarrow 0} \int_S (\varepsilon^{-2} W_2 + W_0) \left(\sum_{k \geq 0} \varepsilon^k (D'\psi_k, (\varepsilon x_3, \varepsilon^{-1} \partial_3 \psi_k)) \right) dx \\ & = \liminf_{\varepsilon \rightarrow 0} \mathbf{J}(\varepsilon)(\psi) < \infty. \end{aligned}$$

Hence, we have

$$\liminf_{\varepsilon \rightarrow 0} W_2 \left(\sum_{k \geq 0} \varepsilon^k (D'\psi_k, (\varepsilon x_3, \varepsilon^{-1} \partial_3 \psi_k)) \right) = 0 \text{ a.e.}$$

From the assumption (2) made on the behavior of strongly extended fibers, it follows that for almost every $x \in S$, $\sum_{k \geq 0} \varepsilon^{k-1} \partial_3 \psi_k$ remains bounded (up to a subsequence), that is $\partial_3 \psi_0 = 0$. Now, since

$$(14) \quad \sum_{k \geq 0} \varepsilon^k (D'\psi_k, (\varepsilon x_3, \varepsilon^{-1} \partial_3 \psi_k)) \xrightarrow[\varepsilon \rightarrow 0]{L^\infty} (D'\psi_0, (0, \partial_3 \psi_1)),$$

and W_2 is assumed to be continuous, we have $W_2(D'\psi_0, (0, \partial_3 \psi_1)) = 0$ almost everywhere. From the hypothesis (3), we get $D'\psi_0 \in \mathcal{M}'$ and $\partial_3 \psi_1 = n(D'\psi_0)$. As a conclusion, every sequence of deformations of finite elastic energy belongs to Φ as announced.

We move on to the next step. Let us consider an element $\psi \in \Phi$ and its associated energy

$$\begin{aligned} \lim_{\varepsilon \rightarrow 0} \mathbf{J}(\varepsilon)(\psi) &= \lim_{\varepsilon \rightarrow 0} \varepsilon^{-2} \int_S W_2 \left(\sum_{k \geq 0} \varepsilon^k (D'\psi_k, (\varepsilon x_3, \varepsilon^{-1} \partial_3 \psi_k)) \right) dx \\ & \quad + \int_S W_0(D'\psi_0, (0, \partial_3 \psi_1)) dx. \end{aligned}$$

Considering that W_2 is a C^2 -function, and that $W_2(D'\psi_0, (0, \partial_3\psi_1)) = 0$, combined with (12) with $\gamma(t) = (D'\psi_0 + tD'\psi_1, tx_3, \partial_3\psi_1 + t\partial_3\psi_2)$, Lebesgue's theorem implies

$$\begin{aligned} \lim_{\varepsilon \rightarrow 0} \varepsilon^{-2} \int_S W_2 \left(\sum_{k \geq 0} \varepsilon^k (D'\psi_k, (\varepsilon x_3, \varepsilon^{-1} \partial_3 \psi_k)) \right) dx \\ = \frac{1}{2} \int_S D^2 W_2 \left[(D'\psi_1, D'\psi_0(x), x_3, \partial_3 \psi_2) \right]^2 dx \\ = \frac{1}{2} \int_S D^2 W_2 \left[(D'\psi_1, D'\psi_0(x), x_3, \partial_3 \psi_2) \right]^2 dx \\ = \frac{1}{2} \int_S Q_{D'\psi_0}(D'\psi_1, x_3, \partial_3 \psi_2) dx. \end{aligned}$$

The limit of the elastic energy $\mathbf{J}(\varepsilon)$ falls out from the fact that ψ_1 may be written as

$$\psi_1 = \frac{1}{2} \int_{-1}^1 \psi_1 dx_3 + x_3 n(D'\psi_0)$$

Finally, according to the definition (5) of f , we have

$$\mathbf{I}(\varepsilon)(\psi) = \mathbf{J}(\varepsilon)(\psi) - \int_S f(x', \varepsilon x_3) \cdot \left(\sum_{k \geq 0} \varepsilon^k \psi_k \right) dx.$$

The second term on the right-hand side converges towards $2 \int_S f \cdot \psi_0 dx$ as ε goes to zero. \square

Since the limit energy is finite only for elements ψ in Φ , ϕ has to be equal to an element of Φ on the subset Γ of the boundary where clamping conditions are imposed.

Corollary 1. *If the minimizers $\varphi(\varepsilon)$ of the total rescaled energy $I(\varepsilon)$ over $\Psi(\varepsilon)$ admit an asymptotic expansion as in (6), and if their total energy $I(\varepsilon)(\varphi(\varepsilon))$ remains bounded, then $\phi_0(x', x_3)$ depends only on $x' \in \Gamma$. In addition, there exists $u_\gamma, n_\gamma \in W^{1,\infty}(\gamma)^3$ such that*

$$\phi_1(x', x_3) = u_\gamma(x') + x_3 n_\gamma(x') \quad \text{for } x \in \Gamma.$$

Note: since ϕ_0 depends only on x' , we shall write $\phi_0(x')$ instead of $\phi_0(x', x_3)$ henceforth.

3.4. Convergence of the minimizers.

Lemma 1. *If the minimizers $\varphi(\varepsilon)$ of the total rescaled energy $I(\varepsilon)$ over $\Psi(\varepsilon)$ admit an asymptotic expansion as in (6), and if their total energy $I(\varepsilon)(\varphi(\varepsilon))$ remains bounded, then*

$$I_0 \left(\varphi_0, \frac{1}{2} \int_{-1}^1 \varphi_1 dx_3, \partial_3 \varphi_2 \right) \leq \inf_{(\psi_0, u, v) \in \Phi_0} I_0(\psi_0, u, v),$$

where

$$\begin{aligned} \Phi_0 := & \left\{ (\psi_0, u, v) \in W^{1,\infty}(S')^3 \times W^{1,\infty}(S')^3 \times W^{1,\infty}(S'; L^\infty(-1, 1)^3) : \right. \\ & D'\psi_0 \in \mathcal{M}' \text{ a.e., } n(D'\psi_0) \in W^{1,\infty}(S')^3, \psi_0(x') = \phi_0(x'), u(x') = u_\gamma(x'), \\ & \left. n(D'\psi_0(x')) = n_\gamma(x') \text{ for every } x' \in \gamma, \text{ and } v(x) = \partial_3 \phi_2(x) \text{ for } x \in \Gamma \right\}. \end{aligned}$$

Proof. Let (φ_k) be the asymptotic expansion of a minimizer $\varphi(\varepsilon)$ of the total energy $I(\varepsilon)$. For every $(\psi_0, u, v) \in \Phi_0$, we set

$$\psi_1 = u + x_3 n(D'\psi_0), \quad \psi_2(x) = \phi_2(x', 0) + \int_0^{x_3} v(x', s) ds, \quad \text{and } \psi_k = \phi_k \quad \forall k \geq 3.$$

Using the Leibniz integral rule, we get that $D'\psi_2 = D'\phi_2 + \int_0^{x_3} D'v(x', s) ds$ and $\partial_3 \psi_2 = v$. It follows that ψ_2 belongs to $W^{1,\infty}(S')^3$. As u and $n(D'\psi_0)$ are assumed to be Lipschitzian, so is ψ_1 and $\boldsymbol{\psi} \in \ell^1(W^{1,\infty}(S')^3)$. From Corollary 1, ϕ_0 depends only on x' and $\phi_1(x', x_3) = u_\gamma(x') + x_3 n_\gamma(x')$ on Γ . As $(\psi_0, u, v) \in \Phi_0$, we infer that $\psi_0 = \phi_0$ and $\psi_1 = \phi_1$ on Γ . Similarly, $\psi_2 = \phi_2$ on Γ . Thus, $\boldsymbol{\psi}$ belongs to $\boldsymbol{\Phi}$ and $\mathbf{I}(\varepsilon)(\boldsymbol{\varphi}) \leq \mathbf{I}(\varepsilon)(\boldsymbol{\psi})$. Letting ε goes to zero, we get from Proposition 1 that

$$(15) \quad I_0 \left(\varphi_0, \frac{1}{2} \int_{-1}^1 \varphi_1 dx_3, \partial_3 \varphi_2 \right) \leq I_0(\psi_0, u, v).$$

□

Lemma 2. *If the minimizers $\varphi(\varepsilon)$ of the total rescaled energy $I(\varepsilon)$ over $\Psi(\varepsilon)$ admit an asymptotic expansion as in (6), and if their total energy $I(\varepsilon)(\varphi(\varepsilon))$ remains bounded, then*

$$\left(\varphi_0, \frac{1}{2} \int_{-1}^1 \varphi_1 dx_3 \right) = \arg \min_{(\psi_0, u) \in \Phi_1} I_1(\psi_0, u),$$

where

$$\begin{aligned} \Phi_1 := & \left\{ (\psi_0, u) \in W^{1,\infty}(S')^3 \times W^{1,\infty}(S')^3 : D'\psi_0 \in \mathcal{M}' \text{ a.e.,} \right. \\ & n(D'\psi_0) \in W^{1,\infty}(S')^3, \psi_0(x') = \phi_0(x'), u(x') = u_\gamma(x'), \\ & \left. \text{and } n(D'\psi_0(x')) = n_\gamma(x') \text{ for } x' \in \gamma \right\}, \end{aligned}$$

$$\begin{aligned} I_1(\psi_0, u) := & \int_{S'} Q_{D'\psi_0}^0(D'u, 0) dx' + \frac{1}{3} \int_{S'} Q_{D'\psi_0}^0(D'n, 1) dx' \\ & + 2 \int_{S'} W_0(D'\psi_0, n_0(D'\psi_0)) dx' - 2 \int_{S'} f_0 \cdot \psi_0 dx', \end{aligned}$$

$f_0(x') = f(x', 0)$ for every $x' \in S'$, and

$$(16) \quad Q_{F'}^0(G', x_3) = \inf_{v \in \mathbb{R}^3} Q_{F'}(G', x_3, v).$$

Proof. From Lemma 1, we have

$$I_0 \left(\varphi_0, \frac{1}{2} \int_{-1}^1 \varphi_1 dx_3, \partial_3 \varphi_2 \right) \leq \inf_{(\psi_0, u, v) \in \Phi_0} I_0(\psi_0, u, v).$$

Moreover, from Proposition 1, we have $\varphi \in \Phi$ which implies that

$$\left(\varphi_0, \frac{1}{2} \int_{-1}^1 \varphi_1 dx_3 \right) \in \Phi_1.$$

Since $\Phi_0 = \Phi_1 \times V$ with

$$V := \{v \in W^{1,\infty}(S'; L^\infty(-1,1)^3) : v(x) = \partial_3 \phi_2(x) \text{ for } x \in \Gamma\},$$

it follows that

$$I_1 \left(\varphi_0, \frac{1}{2} \int_{-1}^1 \varphi_1 dx_3 \right) = \inf_{(\psi_0, u) \in \Phi_1} \inf_{v \in V} I_0(\psi_0, u, v).$$

To complete the proof, we need to show that for every $(\psi_0, u) \in \Phi_1$, we have

$$(17) \quad \inf_{v \in V} I_0(\psi_0, u, v) = I_1(\psi_0, u).$$

We recall that for every $(\psi_0, u) \in \Phi_1$ and every $v \in V$, we have

$$\begin{aligned} I_0(\psi_0, u, v) &= \frac{1}{2} \int_S Q_{D'\psi_0}(D'u + x_3 D'n, x_3, v) dx \\ &\quad + 2 \int_{S'} W_0(D'\psi_0, n_0(D'\psi_0)) dx' - 2 \int_{S'} f_0 \cdot \psi_0 dx'. \end{aligned}$$

Next, the definition of $Q_{D'\psi_0}^0$ entails that

$$\begin{aligned} I_0(\psi_0, u, v) &\geq \frac{1}{2} \int_S Q_{D'\psi_0}^0((D'u, 0) + x_3(D'n, 1)) dx \\ &\quad + 2 \int_{S'} W_0(D'\psi_0, n_0(D'\psi_0)) dx' - 2 \int_{S'} f_0 \cdot \psi_0 dx'. \end{aligned}$$

Furthermore, for every $x' \in S'$ and every $F \in T_{(x', 0)}^*(S; \mathbb{R}^3)$, the quadratic form $Q_{F'}^0$ derives from a bilinear form. Hence,

$$\begin{aligned} \int_{-1}^1 Q_{D'\psi_0}^0((D'u, 0) + x_3(D'n, 1)) dx_3 &= \int_{-1}^1 [Q_{D'\psi_0}^0(D'u, 0) + x_3^2 Q_{D'\psi_0}^0(D'n, 1)] dx_3 \\ &= 2Q_{D'\psi_0}^0(D'u, 0) + \frac{2}{3} Q_{D'\psi_0}^0(D'n, 1). \end{aligned}$$

Accordingly, we obtain that $I_0(\psi_0, u, v) \geq I_1(\psi_0, u)$, so that $\inf_{v \in V} I_0(\psi_0, u, v) \geq I_1(\psi_0, u)$. It remains to prove the converse inequality to establish (17). For every $\delta \geq 0$, we have

$$I_0(\psi_0, u, v) \leq I_0(\psi_0, u, v) + \int_S \delta |v|^2 dx.$$

As a consequence

$$\inf_{v \in V} I_0(\psi_0, u, v) \leq \inf_{v \in V} \left(I_0(\psi_0, u, v) + \int_S \delta |v|^2 dx \right).$$

From the assumptions made on W_2 , it follows that the quadratic form $Q_{D'\psi_0}$ is positive semidefinite almost everywhere. As a consequence, the map $v \mapsto Q_{D'\psi_0}(D'u + x_3 D'n, x_3, v) + \delta |v|^2$ admits a unique minimizer for almost every $x \in S$. Let $v_\delta : S \rightarrow \mathbb{R}^3$ be the map such that $v_\delta = \arg \min_v Q_{D'\psi_0}(D'u + x_3 D'n, x_3, v) + \delta |v|^2$.

Since W_2 is assumed to be of class C^2 and $D'\psi_0$ is bounded, the norm of the quadratic form $Q_{D'\psi_0}$ is uniformly bounded. As a result, v_δ is measurable and belongs to $L^\infty(S)^3$. Also, there exists a sequence v_δ^k in V converging towards v_δ

in $L^2(S)^3$ as k goes to infinity, due to the density of V in $L^2(S)^3$. For every k , we have

$$\begin{aligned} \inf_{v \in V} I_0(\psi_0, u, v) &\leq \frac{1}{2} \int_S [Q_{D'\psi_0}(D'u + x_3 D'n, x_3, v_\delta^k) + \delta |v_\delta^k|^2] dx \\ &\quad + 2 \int_{S'} W_0(D'\psi_0, n_0(D'\psi_0)) dx' - 2 \int_{S'} f_0 \cdot \psi_0 dx' \end{aligned}$$

Taking the limit with respect to k , we infer that

$$\begin{aligned} \inf_{v \in V} I_0(\psi_0, u, v) &\leq \frac{1}{2} \int_S [Q_{D'\psi_0}(D'u + x_3 D'n, x_3, v_\delta) + \delta |v_\delta|^2] dx \\ &\quad + 2 \int_{S'} W_0(D'\psi_0, n_0(D'\psi_0)) dx' - 2 \int_{S'} f_0 \cdot \psi_0 dx' \end{aligned}$$

Note that $Q_{D'\psi_0}(D'u + x_3 D'n, x_3, v_\delta) + \delta |v_\delta|^2$ is a decreasing sequence (as δ goes to zero) of nonnegative functions. Therefore, its integral over S converges towards its pointwise limit $Q_{D'\psi_0}^0(D'u + x_3 D'n, x_3)$, and the intended inequality follows

$$\begin{aligned} \inf_{v \in V} I_0(\psi_0, u, v) &\leq \frac{1}{2} \int_S Q_{D'\psi_0}^0(D'u + x_3 D'n, x_3) dx \\ &\quad + 2 \int_{S'} W_0(D'\psi_0, n) dx' - 2 \int_{S'} f_0 \cdot \psi_0 dx' = I_1(\psi_0, u). \end{aligned}$$

□

3.5. Boundary conditions. An interesting feature of the limit energy is that it depends on both ψ_0 and $u = \frac{1}{2} \int \psi_1 dx_3$. For general boundary conditions, it implies a coupling between both quantities through the term $\int_{S'} Q_{D'\psi_0}^0(D'u, 0) dx'$ of $I_1(\psi_0, u)$. Hence, small perturbations scaling as the thickness of the shell may have an influence on the deformation ψ_0 of the midsection. In the literature, the boundary conditions are usually chosen to satisfy $u_\gamma = 0$, that is,

$$(18) \quad \phi_0(x) = \phi_0(x') \quad \text{and} \quad \phi_1(x) = x_3 n_\gamma(x') \quad \text{for every } (x', x_3) \in \Gamma^\varepsilon,$$

where n_γ is a unit vector. In this case, the minimization of $I_1(\psi_0, u)$ with respect to u is trivial and the limit energy can be expressed solely in terms of ψ_0 .

Proposition 2. *If the minimizers $\varphi(\varepsilon)$ of the total rescaled energy $I(\varepsilon)$ over $\Psi(\varepsilon)$ admit an asymptotic expansion as in (6), and if their total energy $I(\varepsilon)(\varphi(\varepsilon))$ remains bounded with $u_\gamma = 0$ on γ , then*

$$\varphi_0 = \arg \min_{\psi_0 \in \Psi_0} I_0(\psi_0),$$

where

$$(19) \quad \begin{aligned} I_0(\psi_0) &:= \frac{1}{3} \int_{S'} Q_{D'\psi_0}^0(D'n, 1) dx' \\ &\quad + 2 \int_{S'} W_0(D'\psi_0, n_0(D'\psi_0)) dx' - 2 \int_{S'} f_0 \cdot \psi_0 dx', \end{aligned}$$

$f_0(x') = f(x', 0)$ for every $x' \in S'$, and

$$\Psi_0 := \left\{ \psi_0 \in W^{1,\infty}(S')^3 : n = n(D'\psi_0) \in W^{1,\infty}(S')^3, \right. \\ \left. D'\psi_0 \in \mathcal{M}', \psi_0(x') = \phi_0(x'), \text{ and } n = n_\gamma(x') \text{ for every } x' \in \gamma \right\}.$$

4. INVARIANCE AND FLEXURAL ENERGY

Under several assumptions on the stored energy function W_2 , the expression of the flexural part

$$(20) \quad I_{flex}(\psi_0) := \frac{1}{3} \int_{S'} Q_{D'\psi_0}^0(D'n, 1) dx$$

of the total limit energy $I_0(\psi_0)$ may be reduced. More precisely, we shall consider the implications of homogeneity along the fibers, frame-indifference (left invariance under $\text{SO}(3)$), planar isotropy (right invariance under in-plane rotations), and finally right invariance under the special linear group of TS' of the stored energy.

4.1. Homogeneity along the fibers. We say that the shell is homogeneous along the fibers if for every $(F', x_3, v) \in T^*(S; \mathbb{R}^3) = T^*(S'; \mathbb{R}^3) \times (-1, 1) \times (\mathbb{R}^3)^3$, we have $W_2(F', x_3, v) = W_2(F', 0, v)$. In this case, we have for every $(G', s, v) \in T^*(S'; \mathbb{R}^3) \times \mathbb{R} \times (\mathbb{R}^3)^3$ and every section F' of \mathcal{M}' , $D^2W_2[G'_{F'}, s_0, v_{n(F')}]^2 = D^2W_2[G'_{F'}, 0, v_{n(F')}]^2$. It follows that $Q_{F'}^0(G', s)$ is independent of s , and is simply denoted by $Q_{F'}^0(G')$ so that

$$I_{flex}(\psi_0) = \frac{1}{3} \int_{S'} Q_{D'\psi_0}^0(D'n) dx.$$

4.2. Frame-indifference. The principle of frame-indifference states that the space is invariant under rotation which translates in our case in the following condition on the stored energy function W^ε , $W^\varepsilon(F) = W^\varepsilon(RF)$ for every rotation $R \in \text{SO}(3)$. This is assumed in the sequel. Accordingly, the same property is satisfied by W_2 , i.e., $W_2(F) = W_2(RF)$, for every $R \in \text{SO}(3)$.

In the following, we denote by $\mathcal{E}_{S'}$ the set of symmetric bilinear forms on TS' , that is, the fiber bundle of base space S' and whose fiber $(\mathcal{E}_{S'})_{x'}$ at $x' \in S'$ is the set of symmetric bilinear forms on $T_{x'}S'$. The fiber bundle \mathcal{E}_S is defined in a similar way and \mathcal{E}'_S stands for its restriction to S' . Moreover, if $F \in T_x^*(S; \mathbb{R}^3)$, $F^T F$ stands for the element of $(\mathcal{E}_{S'})_x$ that maps every element (u, v) of $(T_x S)^2$ to the scalar product between Fu and Fv . A similar notation is used to defined $(F')^T F' \in \mathcal{E}_{S'}$ for every $F' \in T^*(S'; \mathbb{R}^3)$.

Lemma 3. *If the stored energy W_2 is frame-indifferent, then for every $F' \in \mathcal{M}'$ of maximum rank and every $R \in \text{SO}(3)$, we have*

$$RF' \in \mathcal{M}' \quad \text{and} \quad n(RF') = Rn(F').$$

Moreover, there exists a bundle map $\tau' : \mathcal{M}' \rightarrow TS'$ and a map $\tau_3 : \mathcal{M}' \rightarrow \mathbb{R}$ such that for every $F' \in \mathcal{M}'$ of maximal rank,

$$n(F') = F' \tau'(F') + n_{F'} \tau_3(F'),$$

with both $\tau'(F')$ and $\tau_3(F')$ depending only on $C' = (F')^T F'$ and $n_{F'} \in \mathbb{R}^3$ is defined by

$$(21) \quad n_{F'} \cdot w = \det(F', w) \quad \text{for every } w \in \mathbb{R}^3.$$

Lastly, $C = (F', n_0(F'))^T (F', n_0(F'))$ depends only on C' .

Proof. The first part of the proposition is obvious. Next, since F' is of maximum rank, $(F', n_{F'})$ is invertible so we can set $(\tau'(F'), \tau_3(F')) = (F', n_{F'})^{-1}n(F')$. Moreover, we can check that

$$\begin{aligned} (\tau'(RF'), \tau_3(RF')) &= (RF', n_{RF'})^{-1}n(RF) = (RF', R(n_{F'}))^{-1}Rn(F) \\ &= (R(F', n_{F'}))^{-1}Rn(F) = (F', n_{F'})^{-1}n(F') = (\tau'(F'), \tau_3(F')), \end{aligned}$$

whence both τ' and τ_3 only depend on $(F')^T F'$. Finally, it is readily verified that both $n(F')^T n(F')$ and $n(F')^T F'$ are invariant under rotations of F' . As a result, $C = (F', n_0(F'))^T (F', n_0(F'))$ depends only on $F'^T F'$ as well. \square

Since $T_{(x', x_3)}S = T_{x'}S' \times T_{x_3}(-1, 1) = T_{x'}S' \times \mathbb{R}$, every element $C \in (\mathcal{E}_S)_x$ can be decomposed univocally into $(C', C_3, C_{33}) \in (\mathcal{E}_{S'})_{x'} \times T_{x'}^*S' \times \mathbb{R}$ such that for every (u', u_3) and every (v', v_3) in $T_{x'}S' \times \mathbb{R} = T_x S$,

$$C((u', u_3), (v', v_3)) = C'(u', v') + u_3 C_3(v') + v_3 C_3(u') + C_{33}u_3v_3.$$

In addition, we write this decomposition as follows

$$C = \begin{pmatrix} C' & C_3^T \\ C_3 & C_{33} \end{pmatrix}.$$

Let us introduce the fiber bundle \mathcal{P} of base space S' , and whose fiber at $x' \in S'$ is the set of polynomials of degree lower than or equal to two on $(\mathcal{E}_{S'})_{x'}$.

Proposition 3. *If the stored energy function W_2 is frame-indifferent, then there exists a bundle map $P : C' \mapsto P_{C'}$ over S' from $\mathcal{E}_{S'}$ into \mathcal{P} such that for every deformation ψ_0 of finite limit energy $I_0(\psi_0)$, we have for every $G' \in T^*(S'; \mathbb{R}^3)$,*

$$Q_{D'\psi_0}^0(G', 1) = P_{C'}(D'\psi_0^T G' + G'^T D'\psi_0),$$

where $C' = D'\psi_0^T D'\psi_0$ for short. Moreover, if W_2 is homogeneous along the fibers, then $P_{C'}$ is homogeneous of degree two.

Proof. Let $\mathcal{M}^+ = \{F \in \mathcal{M} : \det F > 0\}$. Since W_2 is assumed to be frame-indifferent, there exists a map $W_S : \mathcal{E}_S \rightarrow \mathbb{R}$ such that for every F in a neighborhood of \mathcal{M}^+ ,

$$(22) \quad W_2(F) = W_S \circ m(F),$$

where $m : T^*(S; \mathbb{R}^3) \rightarrow \mathcal{E}_S$ is the bundle map defined by $m(G) = G^T G$. Let F' be a section of \mathcal{M}' , $s \in \mathbb{R}$ and $G' \in T^*(S'; \mathbb{R}^3)$ such that $(F', n_0(F')) \in \mathcal{M}^+$ a.e. Then, definition (10) combined with (22), gives

$$Q_{F'}^0(G', s) = \inf_{v \in \mathbb{R}^3} D^2 W_2[G'_{F'}, s_0, v_{n(F')}]^2 = \inf_{v \in \mathbb{R}^3} D^2 W_S[Dm(G'_{F'}, s_0, v_{n(F')})]^2.$$

Since $\mathcal{E}_S = (-1, 1) \times \mathcal{E}'_S$, we can identify $T\mathcal{E}_S$ with $T(-1, 1) \times T\mathcal{E}'_S$. Doing so, we obtain that

$$\begin{aligned} Dm(G'_{F'}, s_0, v_{n(F')}) &= \frac{d}{dt}(ts, C + tE)(t=0) \\ &= (s_0, E_C), \end{aligned}$$

where

$$E = \begin{pmatrix} (F')^T G' + (G')^T F' & (F')^T v + (G')^T n \\ v^T F' + n^T G' & n^T v + v^T n \end{pmatrix},$$

$n(F')$ is denoted n for short, and $C = m(F)$. Since $(F', n(F'))(x')$ is assumed to be invertible for every $x' \in S'$, setting $w' = F'^T v + G'^T n$ and $w_3 = n^T v + v^T n$, we get

$$(23) \quad Q_{F'}^0(G', 1) = P_{C'}((F')^T G' + (G')^T F'),$$

where $P_{C'}$ is the section of \mathcal{P} is defined for every $M \in \mathcal{E}_{S'}$ by

$$(24) \quad P_{C'}(M) = \inf_{w \in \mathbb{R}^3} D^2 W_S \left[1_0, \begin{pmatrix} M & w' \\ (w')^T & w_3 \end{pmatrix}_C \right]^2,$$

and $C = (F', n_0(F'))^T (F', n_0(F'))$, which according to Lemma 3 depends only on C' . Finally, if ψ_0 is a deformation satisfying $I_0(\psi_0) < \infty$, owing to the non-interpenetration assumptions made, we know that $D\psi_0 \in \mathcal{M}^+$ a.e. As a result, $Q_{D'\psi_0}^0(D'n, 1) = P_{C'}(M)$ a.e. on S' with $M = F^T D'n + D'n^T F'$ as claimed. Moreover, if the shell is homogeneous along its fibers, (24) reduces to

$$(25) \quad P_{C'}(M) = \inf_{w \in \mathbb{R}^3} D^2 W_S \left[0, \begin{pmatrix} M & w' \\ (w')^T & w_3 \end{pmatrix}_C \right]^2,$$

which is homogeneous of degree two with respect to M . \square

Remark 1. Note that $D'\psi_0^T D'n + D'n^T D'\psi_0$ is not, in general, the second fundamental form of the deformation, except in the case where $n(D'\psi_0)$ is the normal vector to $D'\psi_0$.

4.3. Planar isotropy. We say that the material is isotropic along the midsection of the shell, if for every planar rotation R in the set $\text{SO}(T_{x'} S')$ of rotations in $T_{x'} S'$, and every $F = (F', F_3) \in T_{x'}^*(S^\varepsilon; \mathbb{R}^3)$, we have

$$(26) \quad W^\varepsilon(F', F_3) = W^\varepsilon(F' R, F_3).$$

As a consequence, for deformations of finite energy, the fibers of the shell remain normal to its section.

Lemma 4. Assume that the shell is isotropic along its midsection, then there exists a map $\tau_3 : \mathcal{M}' \rightarrow \mathbb{R}$ such that for every $F' \in \mathcal{M}'$ of maximal rank,

$$n(F') = n_{F'} \tau_3(F'),$$

where $n_{F'}$ is defined by (21). Moreover, $\tau_3(F')$ depends only on the metric $C' = F'^T F'$.

Proof. Let F' be an element of \mathcal{M}' of maximal rank. By definition, we have

$$n_0(F') = \arg \min_{F_3 \in T_0^*((-1,1); \mathbb{R}^3)} W_2(F', F_3).$$

For every rotation $R \in \text{SO}(T S')$, the isotropy property yields

$$\begin{aligned} n_0(F' R) &= \arg \min_{F_3 \in T_0^*((-1,1); \mathbb{R}^3)} W_2(F' R, F_3) \\ &= \arg \min_{F_3 \in T_0^*((-1,1); \mathbb{R}^3)} W_2(F', F_3) = n_0(F'). \end{aligned}$$

In particular, this entails that $n(-F') = n(F')$. What is more, due to frame-indifference, we have from Lemma 3

$$n(F') = F' \tau'(F') + n_{F'} \tau_3(F'),$$

where τ' is a bundle map from \mathcal{M}' into TS' and τ_3 a map from \mathcal{M}' into \mathbb{R} , both of them depending only on the metric $C' = F'^T F'$. Thus,

$$\begin{aligned} F'\tau'(F') + n_{F'}\tau_3(F') &= n(F') = n(-F') = -F'\tau'(-F') + n_{-F'}\tau_3(-F') \\ &= -F'\tau'(F') + n_{F'}\tau_3(F'). \end{aligned}$$

Consequently, $F'\tau'(F') = 0$ and $n(F') = n_{F'}\tau_3(F')$ as claimed. \square

Proposition 4. *Assume that the shell is isotropic along its midsection, then the flexural energy $I_{flex}(\psi_0)$ depends only on the metric and the second fundamental form of the deformed surface. Namely, we have*

$$I_{flex}(\psi_0) = \frac{1}{3} \int_{S'} P_{C'}(|n(D'\psi_0)|b_{D'\psi_0}) dx',$$

where $b_{D'\psi_0}$ is the second fundamental form of ψ_0 , i.e.,

$$(27) \quad b_{F'} = D'N^T F' + F'^T D'N,$$

with $N = n_{F'}/|n_{F'}|$ and $P_{C'}$ is defined by Proposition 3.

Proof. Since $n(D'\psi_0)$ (denoted n for short) is normal colinear to $n_{D'\psi_0}$ and thus to the normal N to the deformed surface, we get

$$D'n^T D'\psi_0 + D'\psi_0^T D'n = (n \cdot N)(D'N^T D'\psi_0 + D'\psi_0^T D'N).$$

\square

4.4. Right invariance under the special linear group. We denote by $SL(TS')$ the special linear group over TS' , that is the fiber bundle over S' whose fiber at x' is the linear diffeomorphisms of $T_{x'}S'$ of determinant equal to one. In this section, we consider the case where the energy W_2 is right invariant under the special linear group, that is $W_2(F', F_3) = W_2(F'U, F_3)$, for every $x' \in S'$, $U \in SL(T_{x'}S')$ and $(F', F_3) \in T_{x'}^*(S'; \mathbb{R}^3) \times T^*((-1, 1); \mathbb{R}^3)$.

Proposition 5. *Assume that W_2 is right invariant under $SL(TS')$, then the flexural energy $I_{flex}(\psi_0)$ depends only on the metric, and on the mean curvature $H = \text{Tr}(C'^{-1/2}b_{D'\psi}C'^{-1/2})$ of the deformation. More precisely, we have*

$$(28) \quad I_{flex}(\psi_0) = \frac{1}{3} \int_{S'} K_{x', \det(C')}(H) dx',$$

where $K : S' \times \mathbb{R} \rightarrow \Pi_2; \Pi_2$ is the set of polynomials of degree lower or equal to 2. Moreover, if the shell is homogeneous along its fibers, then

$$(29) \quad I_{flex}(\psi_0) = \frac{1}{3} \int_{S'} \kappa_{x', \det(C')} |H|^2 dx',$$

where κ is a map from $S' \times \mathbb{R}^+$ into \mathbb{R}^+ .

Proof. Let \mathcal{O} be the fiber bundle over S' whose fibers are the maps from $T_{x'}S'$ into itself of zero trace. For every $O \in \mathcal{O}$ and $x' = \pi_{S'}(O)$, there exists a regular map $U : (0, 1) \rightarrow SL(T_{x'}S')$ such that $\dot{U}(0) = O$ and $U(0) = \text{Id}$. Let F be a section of \mathcal{M} , $(G', s, v) \in T_{x'}^*(S'; \mathbb{R}^3) \times \mathbb{R} \times (\mathbb{R}^3)^3$, and let $\gamma(t) = (\gamma'(t), \gamma_3(t))$ be a curve in $T^*(S; \mathbb{R}^3)$ such that $\dot{\gamma}(0) = (G'_{F'}, s_0, v_{n(F')})$ as in (11). From the right

invariance under the special group, for every $U_0 \in \mathrm{SL}(T_{x'}S')$, we have $W_2(\gamma(t)) = W_2(\gamma'(t)U_0U(t), \gamma_3(t))$. As a consequence,

$$(30) \quad D^2W_2 [\dot{\gamma}(0)]^2 = D^2W_2 \left[\frac{d}{dt}(\gamma'U_0U, \gamma_3)|_{t=0} \right]^2.$$

Then, a simple computation yields

$$\frac{d}{dt}(\gamma'U_0U, \gamma_3)|_{t=0} = ((G'U_0 + F'U_0O)_{F'U_0}, s_0, v_n(F')),$$

which, owing to (30) and (10), leads to $Q_{F'}^0(G', s) = Q_{F'U_0}^0(G'U_0 + F'U_0O, s)$. From (23), recalling that $C' = F'^T F'$, we get

$$\begin{aligned} P_{C'}(F'^T G' + G'^T F') &= Q_{F'}^0(G', 1) = Q_{F'U_0}^0(G'U_0 + F'U_0O, 1) \\ &= P_{U_0^T F'^T F' U_0}((F'U_0)^T (G'U_0 + F'U_0O) + (G'U_0 + F'U_0O)^T (F'U_0)) \\ &= P_{U_0^T C' U_0}(U_0^T (F'^T G' + G'^T F') U_0 + U_0^T C' U_0 O + O^T U_0^T C' U_0) \\ &= P_{C'_0} \left((C'_0)^{1/2} [C'_0^{-1/2} U_0^T (F'^T G' + G'^T F') U_0 C'_0^{-1/2} \right. \\ &\quad \left. + C'_0^{1/2} O C'_0^{-1/2} + C'_0^{-1/2} O^T C'_0^{1/2}] C'_0^{1/2} \right) \end{aligned}$$

where $C' = F'^T F'$ and $C'_0 = U_0^T C' U_0$. Since the map $O \mapsto C'_0^{1/2} O C'_0^{-1/2} + C'_0^{-1/2} O^T C'_0^{1/2}$ is a diffeomorphism over the set of symmetric trace-free matrices, the above expression leads to

$$P_{C'}(F'^T G' + G'^T F') = P_{C'_0} \left(\mathrm{Tr} \left(C'_0^{-1/2} U_0^T (F'^T G' + G'^T F') U_0 C'_0^{-1/2} \right) C'_0 / 2 \right).$$

In addition,

$$\mathrm{Tr} \left(C'_0^{-1/2} U_0^T (F'^T G' + G'^T F') U_0 C'_0^{-1/2} \right) = \mathrm{Tr} \left(C'^{-1/2} (F'^T G' + G'^T F') C'^{-1/2} \right),$$

so that we may write

$$P_{C'}(F'^T G' + G'^T F') = P_{C'_0} \left(\mathrm{Tr} \left(C'^{-1/2} (F'^T G' + G'^T F') C'^{-1/2} \right) C'_0 / 2 \right).$$

Since C' is symmetric and nonnegative, there exists a rotation $R \in \mathrm{SO}(T_{x'}S')$ and λ_1, λ_2 nonnegative reals such that

$$C' = R^T \begin{pmatrix} \lambda_1 & 0 \\ 0 & \lambda_2 \end{pmatrix} R.$$

Let us choose $U_0 \in \mathrm{SL}(T_{x'}S')$ in this fashion

$$U_0 = (\det C')^{1/4} R^T \begin{pmatrix} \lambda_1^{-1/2} & 0 \\ 0 & \lambda_2^{-1/2} \end{pmatrix}.$$

Hence, $C'_0 = U_0^T C' U_0 = (\det C')^{1/2} \mathrm{Id}$, so that

$$(31) \quad P_{C'}(F'^T G' + G'^T F') = P_{(\det C')^{1/2} \mathrm{Id}} \left(\mathrm{Tr} \left(C'^{-1/2} (F'^T G' + G'^T F') C'^{-1/2} \right) (\det C')^{1/2} \mathrm{Id} / 2 \right).$$

Using the definition of I_{flex} , and $D'\psi_0^T D'n + D'n^T D'\psi_0 = |n(D'\psi_0)|b_{D'\psi_0}$, we infer that

$$\begin{aligned} I_{flex}(\psi_0) &= \frac{1}{2} \int_{S'} P_{(\det C')^{1/2} \text{Id}} \left(|n(D'\psi_0)| \text{Tr} \left(C'^{-1/2} b_{D'\psi} C'^{-1/2} \right) (\det C')^{1/2} \text{Id} / 2 \right) dx' \\ &= \frac{1}{2} \int_{S'} P_{(\det C')^{1/2} \text{Id}} \left(|n(D'\psi_0)| H(\det C')^{1/2} \text{Id} / 2 \right) dx'. \end{aligned}$$

Finally, the right invariance of W_2 with respect to $\text{SL}(TS')$ implies that $|n(F')|$ depends only on $\det C'$. Setting

$$K_{x', \det C'}(H) = P_{(\det C')^{1/2} \text{Id}}(|n(F')| H(\det C')^{1/2} \text{Id} / 2),$$

we get (28). Moreover, if the shell is homogeneous along its fibers, then $P_{(\det C')^{1/2} \text{Id}}$ is homogeneous of degree two and, accordingly, $K_{x', \det C'}(H)$ is a monomial. \square

5. GEOMETRIC CONFIGURATION

Classically, the energy of an elastic body is not written in terms of the deformation ψ^ε of S^ε , but in terms of the deformation $\tilde{\psi}^\varepsilon$ of the *geometric* configuration $\tilde{S}^\varepsilon := g_\varepsilon(S^\varepsilon)$, where

$$\begin{aligned} g_\varepsilon : S' \times (-\varepsilon, \varepsilon) &\rightarrow \mathbb{R}^3 \\ (x', x_3) &\mapsto x' + x_3 n'(x'), \end{aligned}$$

$n' : S' \rightarrow \mathbb{R}^3$ being the normal to S' . We set $\tilde{S} := \tilde{S}^{\varepsilon_0}$, for a small enough ε_0 such that g_{ε_0} is one to one. In the following, we will always assume that $\varepsilon \leq \varepsilon_0$. We intend to recast our results in this geometric configuration. This is easily achieved by a mere change of variables. To begin with, we have to recast our initial three-dimensional problem in the geometric configuration.

5.1. Recast of the problem. We denote by $\tilde{J}_\varepsilon(\tilde{\psi}^\varepsilon)$ the elastic energy of a deformation $\tilde{\psi}^\varepsilon$ of \tilde{S}^ε and assume that it has the following form

$$\tilde{J}_\varepsilon(\tilde{\psi}^\varepsilon) := \int_{\tilde{S}^\varepsilon} \tilde{W}^\varepsilon(\tilde{x}, \nabla \tilde{\psi}^\varepsilon(\tilde{x})) \, d\tilde{x},$$

where \tilde{W}^ε stands for the stored energy function of the solid. Furthermore, we assume the shell to be submitted to dead body loads \tilde{f}_ε , so that the total energy of the system is given by

$$(32) \quad \tilde{I}_\varepsilon(\tilde{\psi}^\varepsilon) := \tilde{J}_\varepsilon(\tilde{\psi}^\varepsilon) - \int_{\tilde{S}^\varepsilon} \tilde{f}_\varepsilon \cdot \tilde{\psi}^\varepsilon \, d\tilde{x}.$$

Finally, clamping boundary conditions are added on a part of the boundary $g_\varepsilon(\Gamma^\varepsilon) = g_\varepsilon(\gamma \times (-\varepsilon, \varepsilon))$, where $\gamma \subset \partial S'$.

Our aim is to determine the behavior of the minimizers $\tilde{\varphi}^\varepsilon$ of \tilde{I}_ε over

$$\tilde{\Psi}^\varepsilon := \left\{ \tilde{\psi}^\varepsilon \in W^{1,\infty}(\tilde{S}^\varepsilon)^3 : \tilde{\psi}^\varepsilon \circ g_\varepsilon(x^\varepsilon) = \phi^\varepsilon(x^\varepsilon), x^\varepsilon \in \Gamma^\varepsilon \right\},$$

as ε goes to zero under the assumptions, on the stored energy and the applied loads, made hereunder.

In order to apply our results, several assumptions, similar to the ones we made on W^ε and φ^ε , have to be imposed on \tilde{W}^ε and on the minimization sequences $\tilde{\varphi}^\varepsilon$. Dependence of the stored energy with respect to the thickness. We assume that there exists continuous nonnegative maps \tilde{W}_2 and \tilde{W}_0 such that for every $(\tilde{x}, F) \in \tilde{S} \times \mathbb{R}^{3 \times 3}$, we have

$$\tilde{W}^\varepsilon(\tilde{x}, F) = \varepsilon^{-1} \left(\varepsilon^{-2} \tilde{W}_2(\tilde{x}, F) + \tilde{W}_0(\tilde{x}, F) \right).$$

Behavior of strongly extended fibers. We assume that the stored energy \tilde{W}_2 is bounded from below by a positive constant for strongly extended fibers, namely, there exists $\delta, c > 0$ such that

$$(33) \quad \forall \tilde{x} \in \tilde{S}, \forall F' \in T_{x'}^*(S'; \mathbb{R}^3), \forall F_3 \in \mathbb{R}^3 \text{ such that } |F_3| \geq \delta,$$

$$\text{we have } \tilde{W}_2(\tilde{x}, F' \circ \pi'_{x'} + F_3 \otimes n'(x')) \geq c,$$

where x' is the projection of \tilde{x} onto S' and $\pi'_{x'}$ is the projection of \mathbb{R}^3 onto $T_{x'}(S'; \mathbb{R}^3)$.

Regularity and zero set of \widetilde{W}_2 . We assume that \widetilde{W}_2 is a C^2 nonnegative function and denote by $\widetilde{\mathcal{M}}$ the restriction of its zero set to $S' \times \mathbb{R}^{3 \times 3}$, that is

$$\widetilde{\mathcal{M}} := \left\{ (x', F) \in S' \times \mathbb{R}^{3 \times 3} \text{ such that } \widetilde{W}_2(x', F) = 0 \right\}.$$

Let \mathcal{M}' be the projection of $\widetilde{\mathcal{M}}$ onto $T^*(S'; \mathbb{R}^3)$, that is,

$$\mathcal{M}' := \bigcup_{x' \in S'} \left\{ F' \in T_{x'}^*(S'; \mathbb{R}^3) : \text{there exists } n \in \mathbb{R}^3 \text{ such that} \right. \\ \left. (x', F' \circ \pi'_{x'} + n \otimes n'(x')) \in \widetilde{\mathcal{M}} \right\}.$$

Once again, we assume that the projection of $\widetilde{\mathcal{M}}$ onto \mathcal{M}' to be one to one, that is there exists map $n : \mathcal{M}' \rightarrow \mathbb{R}^3$ such that

$$(34) \quad \widetilde{\mathcal{M}} = \left\{ (x', F' \circ \pi'_{x'} + n(F') \otimes n'(x')) \in S' \times \mathbb{R}^{3 \times 3} : F' \in \mathcal{M}' \right\}.$$

Interpenetration. To avoid interpenetration of matter, it is geometric to expect $D\widetilde{\psi}^\varepsilon$ to be invertible. To this end, we require that $\widetilde{W}_0(x', F) = \infty$ for every $(x', F) \in \widetilde{\mathcal{M}}$ such that $\det F < 0$, and that

$$\widetilde{W}_0(x', F) \rightarrow \infty \quad \text{if } (x', F) \in \widetilde{\mathcal{M}} \quad \text{and} \quad \det F \rightarrow 0.$$

Applied loads. The volumic loads are assumed to scale as the inverse of the thickness of the shell and more precisely that there exists $\tilde{f} : \tilde{S} \rightarrow \mathbb{R}^3$ such that

$$\tilde{f}_\varepsilon(\tilde{x}) = \varepsilon^{-1} \tilde{f}(\tilde{x}) \quad \text{for every } \tilde{x} \in \tilde{S}.$$

Ansatz. We assume that the minimizers of the energy admit an asymptotic expansion in this fashion

$$(35) \quad \tilde{\varphi}^\varepsilon(x' + \varepsilon x_3 n') = \sum_{k \geq 0} \varepsilon^k \varphi_k(x', x_3).$$

5.2. Change of variable. In order to apply our result, we first have to rewrite the energy in terms of the associated deformation $\psi^\varepsilon = \tilde{\psi}^\varepsilon \circ g_\varepsilon$ of S^ε . We have

$$(36) \quad \tilde{J}(\tilde{\psi}^\varepsilon) = J(\psi^\varepsilon) = \int_{S^\varepsilon} W^\varepsilon(D\psi^\varepsilon) dx,$$

with $W^\varepsilon = \varepsilon^{-1} (\varepsilon^{-2} W_2 + W_0)$ and for every $F \in T_x^*(S; \mathbb{R}^3)$,

$$(37) \quad W_k(F) = \widetilde{W}_k(F \circ (Dg_\varepsilon(x))^{-1}) \det(Dg_\varepsilon(x)), \quad k = 0, 2.$$

Note that W_2 and W_0 are independent of ε since $Dg_\varepsilon = (\text{Id}', n') + x_3(D'n', 0)$ (which is denoted by Dg hereafter). In addition, these energies satisfy the assumptions made in section 2.1. Finally, the minimizers $\varphi^\varepsilon = \tilde{\varphi}^\varepsilon \circ g_\varepsilon$ admit the same asymptotic expansion than $\tilde{\varphi}^\varepsilon$. Thus, all of the results of Sections 3 and 4 apply and may be expressed in terms of \widetilde{W}_0 and \widetilde{W}_2 up to a change of variable. Moreover, the definitions of \mathcal{M}' and of the map $n : \mathcal{M}' \rightarrow \mathbb{R}^3$ are independent of the chosen approach.

Lemma 5. *If function W_2 is defined by (37), and F' is a section of \mathcal{M}' , then for every $G' \in T_{x'}^*(S'; \mathbb{R}^3)$ and $s \in \mathbb{R}$, we have*

$$Q_{F'}^0(G', s) = \widetilde{Q}_{F'}^0(G' - sF'D'n', s),$$

where

$$\tilde{Q}_{F'}^0(G', s) := \inf_{v \in \mathbb{R}^3} D^2 \tilde{W}_2(x', F) [sn', G' \pi'_{x'} + v \otimes n']^2,$$

where $\pi'_{x'}$ is the projection of \mathbb{R}^3 onto $T_{x'} S'$ and $F = F'(x') \circ \pi'_{x'} + n(F'(x')) \otimes n'(x')$.

Proof. Let F' be a section of \mathcal{M}' , x' be an element of S' and $G' \in T_{x'}^*(S'; \mathbb{R}^3)$. From the definition of $Q_{F'}^0$, we have

$$Q_{F'}^0(G', s) = \inf_{v \in \mathbb{R}^3} D^2 W_2[G'_{F'}, s_0, v_{n(F')}]^2,$$

where $(G'_{F'}, s_0, v_{n(F')}) = \dot{\gamma}(0)$ and $\gamma(t) = (F'(x') + tG', ts, n(F'(x')) + vt) \in T^*(S'; \mathbb{R}^3) \times (-1, 1) \times (\mathbb{R}^3)^3 = T^*(S'; \mathbb{R}^3) \times T^*((-1, 1); \mathbb{R}^3) = T^*(S; \mathbb{R}^3)$. From (37), we deduce

$$(38) \quad Q_{F'}^0(G', s) = \det(Dg(x', 0)) \inf_{v \in \mathbb{R}^3} D^2 \tilde{W}_2[\dot{\gamma}(0)]^2 = \inf_{v \in \mathbb{R}^3} D^2 \tilde{W}_2[\dot{\gamma}(0)]^2,$$

where $\tilde{\gamma}(t) = \gamma(t) \circ Dg(x', ts)^{-1}$. On the other hand, $Dg = (\text{Id}', n') + x_3(D'n', 0)$, and

$$\begin{aligned} (Dg)^{-1} &= ((\text{Id}', n') + x_3(D'n', 0))^{-1} \\ &= ((\text{Id}', n')(\text{Id} + x_3(\text{Id}', n')^{-1}(D'n', 0)))^{-1} \\ &= (\text{Id} - x_3(\text{Id}', n')^{-1}(D'n', 0))(\text{Id}', n')^{-1} + o(x_3). \end{aligned}$$

Since $(\text{Id}' + n' \otimes e_3)^{-1} = \begin{pmatrix} \pi' \\ n'^T \end{pmatrix}$, the above identity reads

$$\begin{aligned} (Dg)^{-1} &= \begin{pmatrix} \pi' \\ n'^T \end{pmatrix} - x_3 \begin{pmatrix} \pi' \\ n'^T \end{pmatrix} (D'n', 0) \begin{pmatrix} \pi' \\ n'^T \end{pmatrix} + o(x_3) \\ &= \begin{pmatrix} \pi' \\ n'^T \end{pmatrix} - x_3 \begin{pmatrix} \pi' D'n' \pi' \\ 0 \end{pmatrix} + o(x_3). \end{aligned}$$

It follows that $\tilde{\gamma}(t) = (x(t), F(t)) + o(t)$ with $x(t) = x' + tsn'$ and, using the notation F' for $F'(x')$ for short,

$$\begin{aligned} F(t) &= ((F', n(F')) + t(G', v)) \left(\begin{pmatrix} \pi' \\ n'^T \end{pmatrix} - ts \begin{pmatrix} \pi' D'n' \pi' \\ 0 \end{pmatrix} \right) \\ &= (F', n(F')) \begin{pmatrix} \pi' \\ n'^T \end{pmatrix} + t(G' \pi' + v \otimes n') - ts F' D'n' \pi'. \end{aligned}$$

Consequently,

$$\begin{aligned} \dot{\gamma}(0) &= \left[\begin{pmatrix} x', (F', n(F')) \begin{pmatrix} \pi' \\ n'^T \end{pmatrix} \end{pmatrix}, (sn', G' \pi' + v \otimes n' - s(F' D'n' \pi')) \right] \\ &= [(x', F' \pi' + n(F') \otimes n'), (sn', (G' - sF' D'n') \pi' + v \otimes n')] \\ &= [(x', F), (sn', (G' - sF' D'n') \pi' + v \otimes n')] \end{aligned}$$

The conclusion follows from (38). \square

From now on, we limit our analysis to the case where standard boundary conditions (18) are applied. From Proposition 2, we immediately infer the next result.

Proposition 6. *Assume that the standard boundary conditions (18) are applied to the shell. Let $\tilde{\varphi}^\varepsilon$ be the minimizer of the total energy $\tilde{I}_\varepsilon(\tilde{\varphi}^\varepsilon)$ over the space of*

admissible deformations. If $\tilde{\varphi}^\varepsilon$ admits an asymptotic expansion as in (35), and if the total energy $\tilde{I}_\varepsilon(\tilde{\varphi}^\varepsilon)$ remains bounded, then

$$\varphi_0 = \arg \min_{\psi_0 \in \Psi_0} \tilde{I}_0(\psi_0),$$

where

$$\begin{aligned} \tilde{I}_0(\psi_0) = & \frac{1}{3} \int_{S'} \tilde{Q}_{D'\psi_0}^0(D'n - D'\psi_0 D'n', 1) dx' + 2 \int_{S'} \tilde{W}_0(x', (D'\psi_0, n)) dx' \\ & - 2 \int_{S'} \tilde{f}_0 \cdot \psi_0 dx', \end{aligned}$$

$\tilde{f}_0(x') = \tilde{f}(x', 0)$ for every $x' \in S'$, $n = n(D'\psi_0)$, and

$$\Psi_0 := \left\{ \psi_0 \in W^{1,\infty}(S')^3 : D'\psi_0 \in \mathcal{M}', n = n(D'\psi_0) \in W^{1,\infty}(S')^3 \right.$$

such that $\psi_0(x') = \phi_0(x')$ and $n(x') = n_\gamma(x')$ for every $x' \in \gamma$.

Note that a result equivalent to Lemma 2 may similarly be stated in the geometric configuration which applies to slightly more general Dirichlet conditions.

5.3. Homogeneity along the fibers. We say that the shell is homogeneous along its fibers in the geometric configuration if for every $x' \in S'$, $s \in (-1, 1)$ and $F \in \mathbb{R}^{3 \times 3}$, we have $\tilde{W}_2(x' + sn', F) = \tilde{W}_2(x', F)$.

Proposition 7. *If the shell is homogeneous along its fibers in the geometric configuration, then $\tilde{Q}_{F'}^0(G', s)$ is independent of s , and is denoted by $\tilde{Q}_{F'}^0(G')$.*

5.4. Frame-indifference. In the following, we assume the stored energy to be frame-indifferent, that is, $\tilde{W}^\varepsilon(\tilde{x}, RF) = \tilde{W}^\varepsilon(\tilde{x}, F)$, for every $(\tilde{x}, F) \in \tilde{S}^\varepsilon \times \mathbb{R}^{3 \times 3}$ (with $\varepsilon > 0$ small enough), and every rotation $R \in \text{SO}(3)$. Note that it is equivalent to the frame-indifference of W^ε .

Proposition 8. *If the stored energy function \tilde{W}_2 is frame-indifferent, then there exists a bundle map $\tilde{P} : C' \mapsto \tilde{P}_{C'}$ over S' from $\mathcal{E}_{S'}$ into \mathcal{P} such that for every deformation ψ_0 of finite energy $\tilde{I}_0(\psi_0)$, we have for every $G' \in T^*(S'; \mathbb{R}^3)$*

$$\tilde{Q}_{D'\psi_0}^0(G', 1) = \tilde{P}_{C'}(D'\psi_0^T G' + G'^T D'\psi_0),$$

with $C' = D'\psi_0^T D'\psi_0$ and $n = n(D'\psi_0)$. Moreover, if the shell is homogeneous along its fibers in the geometric configuration, then $\tilde{P}_{C'}$ is homogeneous of degree two.

Proof. The proof is similar to the one devised for the abstract configuration. Once again, there exists a map \tilde{W}_S such that, at least in a neighborhood of $\tilde{\mathcal{M}}^+ = \{F \in \tilde{\mathcal{M}} : \det F > 0\}$, we may write $\tilde{W}_2(x, F) = \tilde{W}_S(x, F^T F)$. After some computations, we derive the claimed result with

$$(39) \quad \tilde{P}_{C'}(M) = \inf_{\substack{w' \in T_{x'} S' \\ w_3 \in \mathbb{R}}} D^2 \tilde{W}_S(x', \tilde{C}) \left[n', \pi'^T M \pi' + n' w'^T \pi' + \pi'^T w' n'^T + n' w_3 n'^T \right]^2,$$

where $\tilde{C} = (D'\psi_0 \pi' + n(F) \otimes n')^T (D'\psi_0 \pi' + n(F) \otimes n')$. Moreover, frame indifference implies also that \tilde{C} depends only on C' . Finally, if the shell is homogeneous along

its fibers in the geometric configuration, we have
(40)

$$\tilde{P}_{C'}(M) = \inf_{\substack{w' \in T_{x', S'} \\ w_3 \in \mathbb{R}}} \frac{\partial^2 \tilde{W}_S}{\partial \tilde{C}^2} \left(x', \tilde{C} \right) \left[\pi'^T M \pi' + n' w'^T \pi' + \pi'^T w' n'^T + n' w_3 n'^T \right]^2.$$

□

5.5. Planar isotropy. We say the shell is isotropic along its midsection, if for every $x' \in S'$, and $(F', F_3) \in T_{x'}^*(S'; \mathbb{R}^3) \times \mathbb{R}^3$, we have

$$\tilde{W}^\varepsilon(x', F' R \pi'_{x'} + F_3 \otimes n') = \tilde{W}^\varepsilon(x', F' \pi'_{x'} + F_3 \otimes n'),$$

for every $R \in \text{SO}(T_{x'} S')$. This is equivalent to the definition used in the abstract configuration. We investigate the consequences of planar isotropy on the flexural part of the energy

$$\tilde{I}_{flex}(\psi_0) := \frac{1}{3} \int_{S'} \tilde{Q}_{D' \psi_0}^0(D'n - D'\psi_0 D'n', 1) dx'.$$

Proposition 9. *If the shell is isotropic along its midsection, then*

$$\tilde{I}_{flex}(\psi_0) = \frac{1}{3} \int_{S'} \tilde{P}_{C'}(|n(D'\psi_0)| b_{D' \psi_0} - (C' D'n' + (D'n')^T C')) dx',$$

where $b_{D' \psi_0}$ is the second fundamental form of the deformed surface, given by (27).

Proof. The proof is similar to Proposition 4. □

5.6. Right-invariance under the special linear group. We say that the stored energy \tilde{W}_2 is invariant under the special linear group, if for every $\tilde{x} = x' + x_3 n'$, and $(F', F_3) \in T_{x'}^*(S'; \mathbb{R}^3) \times \mathbb{R}^3$, we have

$$\tilde{W}_2(\tilde{x}, F' U \pi'_{x'} + F_3 \otimes n') = \tilde{W}_2(\tilde{x}, F' \pi'_{x'} + F_3 \otimes n'),$$

for every $U \in \text{SL}(T_{x'} S')$. Note that this definition is equivalent to the one given in the abstract configuration.

Proposition 10. *Assume that \tilde{W}_2 is right-invariant under the special linear group, then*

$$(41) \quad \tilde{I}_{flex}(\psi_0) = \frac{1}{3} \int_{S'} \tilde{K}_{x', \det(C')}(|n(D'\psi_0)| H - H_0) dx',$$

where H and H_0 are respectively the mean curvatures of the deformed shell $\psi_0(S')$ and undeformed shell S' . Moreover, if the shell is homogeneous along its fibers, then

$$\tilde{I}_{flex}(\psi) = \frac{1}{3} \int_{S'} \tilde{\kappa}_{x', \det(C')}(|n(D'\psi_0)| H - H_0)^2 dx'.$$

Proof. For all section F' of $T^*(S'; \mathbb{R}^3)$ and $G' \in T_{x'}^*(S'; \mathbb{R}^3)$, we have

$$\tilde{P}_{C'}(F'^T(G' - F'D'n') + (G' - F'D'n')^T F') = P_{C'}(F'^T G' + G'^T F').$$

From Proposition 5 and (31), we deduce that

$$\begin{aligned} \tilde{P}_{C'}(F'^T(G' - F'D'n') + (G' - F'D'n')^T F') = \\ \tilde{P}_{(\det C')^{1/2} \text{Id}} \left(\alpha(\det C')^{1/2} \text{Id} / 2 \right), \end{aligned}$$

with

$$\alpha = \text{Tr} \left(C'^{-1/2} (F'^T (G' - F' D' n') + (G' - F' D' n')^T F') C'^{-1/2} \right).$$

We thus obtain (41) with

$$(42) \quad \tilde{K}_{x', \det(C')} (H) = \tilde{P}_{\det(C') \text{Id}} (H (\det C')^{1/2} \text{Id} / 2).$$

If the shell is homogeneous along its fibers, then the $\tilde{P}_{\det(C') \text{Id}}$ is homogeneous of degree two, wherefrom the conclusion in this case. \square

6. EXAMPLES

We are now in position to apply our formal convergence result to derive different modelings for isometric bending shells, vesicles and RBCs. Note that, in our setting, we do not derive the nonlinear membrane shell model (see [Le Dret and Raoult \[1996\]](#)) since W_2 cannot be chosen to be equal to zero. In all this section, we assume that W^ε or $\widetilde{W}^\varepsilon$ satisfy assumptions (2,3) and (33,34), that the Dirichlet boundary conditions on Γ^ε are given by (7,18) and that the minimizers φ^ε or $\widetilde{\varphi}^\varepsilon$ of the total energy do admit an asymptotic expansion as in (6,35), while their total energies $I_\varepsilon(\varphi^\varepsilon)$ and $\widetilde{I}_\varepsilon(\widetilde{\varphi}^\varepsilon)$ remain bounded. Moreover, the stored energies are assumed to be frame indifferent.

6.1. Isometric bending shells. In this section, we recover the isometric bending shell model by choosing $\widetilde{W}_0 = 0$ and the set $\widetilde{\mathcal{M}}$ of the zeros of \widetilde{W}_2 restricted to the midsection to be equal to

$$(43) \quad \widetilde{\mathcal{M}}_{iso} := S' \times \text{SO}(3).$$

The sequence of minimizers of the energy converges toward the minimizer of an energy whose elastic part depends only on the difference between the second fundamental form of the deformed shell and that of its reference configuration.

Proposition 11. *If $\widetilde{W}^\varepsilon = \varepsilon^{-3}\widetilde{W}_2$; if \widetilde{W}_2 is such that $\widetilde{\mathcal{M}} = \widetilde{\mathcal{M}}_{iso}$ given by (43) then*

$$\varphi_0 = \arg \min_{\psi_0 \in \Psi_0} \widetilde{I}_0(\psi_0),$$

where

$$(44) \quad \widetilde{I}_0(\psi_0) = \frac{1}{3} \int_{S'} \widetilde{P}_{x', \text{Id}} ((D'\psi_0)^T D'n + D'n^T D'\psi_0 - (D'n' + D'n'^T)) dx' - 2 \int_{S'} f_0 \cdot \psi_0 dx',$$

$f_0(x') = f(x', 0)$ for every $x' \in S'$, $n = n(D'\psi_0)$, $\widetilde{P}_{x', \text{Id}}$ is a polynomial of degree at most equal to 2 given by (39) and

$$\Psi_0 = \left\{ \psi_0 \in W^{1,\infty}(S')^3 : (D'\psi_0)^T D'\psi_0 = \text{Id}, n = n(D'\psi_0) \in W^{1,\infty}(S')^3, \right. \\ \left. \psi_0(x') = \phi_0(x') \text{ and } n(x') = n_\gamma(x') \text{ for every } x' \in \gamma \right\}.$$

Moreover, if the shell is homogeneous along its fibers, then $\widetilde{P}_{x', \text{Id}}$ is homogeneous of degree 2.

Proof. Its is a straightforward application of Proposition 6 and Proposition 8. \square

Example. Let us give a practical example. For instance, one can choose the Saint Venant-Kirchhoff nonlinearly elastic stored energy function

$$\widetilde{W}_2(F) = \mu \text{Tr}((F^T F - \text{Id})^2) + \frac{\lambda}{2} \text{Tr}(F^T F - \text{Id})^2.$$

A simple computation leads to the energy

$$\widetilde{I}_0(\psi_0) = \frac{1}{3} \int_{S'} 2\mu \text{Tr}((b - b_{ref})^2) + \frac{\lambda\mu}{2\mu + \lambda} \text{Tr}(b - b_{ref})^2 - f_0 \cdot \psi_0 dx'$$

where $b = (D'\psi_0)^T D'n + (D'n)^T D'\psi_0$ is the second fundamental form of the deformed shell and $b_{ref} = D'n'^T + D'n'$ is the second fundamental form of the undeformed shell.

6.2. Vesicles. In this section, we derive Helfrich functionals, with or without spontaneous curvature, from three-dimensional elasticity. The main difference with the isometric case lies in the fact that we assume the energy to be right-invariant under the special linear group $SL(TS')$. Note that this readily implies that it may not be chosen to be isotropic. The Helfrich functional without spontaneous curvature is derived using the abstract configuration S , while the one with spontaneous curvature is obtained using the geometric configuration \tilde{S} of the shell.

6.2.1. Without spontaneous curvature. In this section, we consider the case, where the zero set of W_2 restricted to the midsection is given by

$$(45) \quad \mathcal{M}_H := \{(F', F_3) \in T^*(S'; \mathbb{R}^3) \times T_0^*((-1, 1); \mathbb{R}^3) : \\ \det(F) = 1, F_3 \cdot v = \det(F', v) \text{ for every } v \in T_0^*((-1, 1); \mathbb{R}^3)\}.$$

From Propositions 2 and 5, we obtain that the minimizers of the energy formally converge toward the Helfrich functional with no spontaneous curvature.

Proposition 12. *If $W^\varepsilon = \varepsilon^{-3}W_2$; if W_2 is right-invariant under $SL(TS')$, and such that $\mathcal{M} = \mathcal{M}_H$ given by (45); If the shell is homogeneous along its fibers in the abstract configuration then*

$$\varphi_0 = \arg \min_{\psi_0 \in \Psi_0} I_0(\psi_0),$$

where

$$(46) \quad I_0(\psi_0) = \frac{1}{3} \int_{S'} \kappa |H|^2 dx' - 2 \int_{S'} f_0 \cdot \psi_0 dx',$$

$f_0(x') = f(x', 0)$ for every $x' \in S'$, H is the mean curvature of the deformed shell $\psi_0(S')$, $\kappa(x') = P_{\text{Id}}(\text{Id}/2)$, where P_{Id} is given by (25), and

$$\Psi_0 = \{\psi_0 \in W^{1,\infty}(S')^3 : n \in W^{1,\infty}(S')^3 \text{ while } \det((D'\psi_0)^T D'\psi_0) = 1, \\ \psi_0(x') = \phi_0(x') \text{ and } n(x') = n_\gamma(x') \text{ for } x' \in \gamma\},$$

and n is the normal to the deformed surface $\psi_0(S')$.

Example. Proposition 12 can be applied with

$$(47) \quad W_2(F) = W_S(C) = \alpha(\det(C) - 1)^2 + \beta|Ce_3 - e_3|^2,$$

where $C = F^T F$ and α and β are positive real constants. A simple computation leads to

$$D^2W_S \left[0, \begin{pmatrix} \text{Id}/2 & w' \\ w' & w_3 \\ & & \text{Id} \end{pmatrix} \right]^2 = 2\alpha(1 + w_3)^2 + 2\beta|(w', w_3)|^2.$$

Then, from the expression (25) of P_{Id} , we get

$$\kappa = P_{\text{Id}}(\text{Id}/2) = \inf_w 2\alpha(1 + w_3)^2 + 2\beta|(w', w_3)|^2 = 2(\alpha^{-1} + \beta^{-1})^{-1}.$$

Hence, the limit energy in this case is

$$I_0(\psi_0) = \frac{2}{3} \int_{S'} (\alpha^{-1} + \beta^{-1})^{-1} |H|^2 dx' - 2 \int_{S'} f_0 \cdot \psi_0 dx'.$$

6.2.2. *With spontaneous curvature.* In this section, we derive from three-dimensional elasticity a model of shells whose limit energy is the Helfrich functional with nonzero spontaneous curvature. Basically, such a model is obtained by using the same assumptions as in the previous case but cast in the geometric configuration, with a set of zeros restricted to the midsection for \widetilde{W}_2 given by

$$(48) \quad \widetilde{\mathcal{M}}_H := \{(x', F) \in S' \times \mathbb{R}^{3 \times 3} : \det(F) = 1, \text{ and } (\text{Cof } F - F)n' = 0\}.$$

The following Proposition is a direct application of Proposition 6 and Proposition 10.

Proposition 13. *If $\widetilde{W}^\varepsilon = \varepsilon^{-3}\widetilde{W}_2$; if \widetilde{W}_2 is right-invariant under $\text{SL}(TS')$, and such that $\widetilde{\mathcal{M}} = \widetilde{\mathcal{M}}_H$ given by (48); if the shell is homogeneous along its fibers in the geometric configuration then*

$$\varphi_0 = \arg \min_{\psi_0 \in \Psi_0} \widetilde{I}_0(\psi_0),$$

where

$$(49) \quad \widetilde{I}_0(\psi_0) = \frac{1}{3} \int_{S'} \widetilde{\kappa} |H - H_0|^2 dx' - 2 \int_{S'} f_0 \cdot \psi_0 dx',$$

$f_0(x') = \widetilde{f}(x', 0)$ for every $x' \in S'$, where H is the mean curvature of the deformed shell $\psi_0(S')$ and H_0 is the mean curvature of S' , with

$$\widetilde{\kappa}(x') = \widetilde{P}_{\text{Id}}(\text{Id}/2),$$

where $\widetilde{P}_{\text{Id}}$ is given by (40) and

$$(50) \quad \Psi_0 = \{\psi_0 \in W^{1,\infty}(S')^3 : n \in W^{1,\infty}(S')^3 \text{ while } \det(D'\psi_0^T D'\psi_0) = 1, \\ \psi_0(x') = \phi_0(x') \text{ and } n(x') = n_\gamma(x') \text{ for every } x' \in \gamma\},$$

where n is the normal to the deformed surface $\psi_0(S')$.

Example. The stored energy function

$$\widetilde{W}_2(x', F) = \widetilde{W}_S(x', \widetilde{C}) = \alpha(\det(F^T F) - 1)^2 + \beta|F^T F n' - n'|^2.$$

satisfies the assumptions of Proposition 13, and we have

$$\widetilde{\kappa} = \inf_{w' \in T_{x', S'}, w_3 \in \mathbb{R}} \frac{\partial \widetilde{W}_S}{\partial \widetilde{C}^2}(x', \text{Id}) [\pi'^T \pi' / 2 + n' w'^T \pi' + \pi'^T w' n'^T + n' w_3 n'^T]^2.$$

Furthermore, we have

$$\frac{\partial \widetilde{W}_S}{\partial \widetilde{C}^2}(x', \text{Id}) [\delta C]^2 = 2(\alpha \text{Tr}(\delta C)^2 + \beta |\delta C n'|^2),$$

so that,

$$\widetilde{\kappa} = \inf_{w' \in T_{x', S'}, w_3 \in \mathbb{R}} 2(\alpha(1 + w_3)^2 + \beta(|w'|^2 + w_3^2)) = 2(\alpha^{-1} + \beta^{-1})^{-1}.$$

For such a choice of \widetilde{W}_2 , and under the assumptions made in Proposition 13, the sequence of minimizers $\widetilde{\varphi}^\varepsilon$ formally converges toward a minimizer of

$$\widetilde{I}_0(\psi_0) = \frac{2}{3} \int_{S'} (\alpha^{-1} + \beta^{-1})^{-1} |H - H_0|^2 dx' - 2 \int_{S'} \widetilde{f}_0 \cdot \psi_0 dx',$$

over Ψ_0 given by (50). Note that this is the set of deformations that preserve the local area of the shell supplemented with boundary conditions on a subset of the boundary.

6.3. Red Blood Cells. The mechanical behavior of a Red Blood Cell (RBC) is driven by the nature of its membrane which is mainly made of a lipid bilayer. Note that in addition to the lipid bilayer, RBCs are also composed of a protein skeleton. This skeleton ensures a small resistance of the RBCs to shear stress. Such a model may be obtained as the limit of the three-dimensional elasticity. In this section, we derive a model of the mechanical behavior of RBCs as the limit of genuine three-dimensional elasticity. To this end, we consider a stored energy W^ε whose asymptotic assumption reads as

$$W^\varepsilon = \varepsilon^{-1}(\varepsilon^{-2}W_2 + W_0),$$

where W_2 satisfies the same assumption as in the study of vesicles without spontaneous curvature (see section 6.2.1), namely, its zero set restricted to the midsection is given by (45). We get that the sequence φ^ε of minimizers formally converges toward

$$\varphi_0 = \arg \min_{\psi_0 \in \Psi_0} I_0(\psi_0),$$

where

$$I_0(\psi_0) = \frac{1}{3} \int_{S'} k|H|^2 dx' + 2 \int_{S'} W_0(D'\psi_0, n) dx' - 2 \int_{S'} f_0 \cdot \psi_0 dx',$$

$f_0(x') = f(x', 0)$, n is the normal to the deformed shell $\psi_0(S')$ and Ψ_0 is the set of deformations that preserve the local area of the shell and satisfy the boundary conditions

$$\psi_0(x') = \phi_0(x') \quad \text{and} \quad n(x') = n_\gamma(x') \quad \text{for every } x' \in \gamma.$$

Example. As an example, we can choose the nonlinearly elastic Saint Venant-Kirchhoff stored energy function

$$W_0(F) = \mu \operatorname{Tr}((C - \operatorname{Id})^2) + \frac{\lambda}{2} \operatorname{Tr}(C - \operatorname{Id})^2, \quad \text{with } C = F^T F,$$

and $W_2(F)$ as in (47). This leads to a limit energy

$$\begin{aligned} I_0(\psi_0) = & \frac{2}{3} \int_{S'} (\alpha^{-1} + \beta^{-1})^{-1} |H|^2 dx' + 2 \int_{S'} \mu \operatorname{Tr}((D'\psi_0^T D'\psi_0 - \operatorname{Id}')^2) \\ & + \frac{\lambda}{2} \operatorname{Tr}(D'\psi_0^T D'\psi_0 - \operatorname{Id}')^2 dx' - 2 \int_{S'} f_0 \cdot \psi_0 dx'. \end{aligned}$$

7. CONCLUSION

In this article, we prove, using a formal approach, that new nonlinearly elastic shell models may be derived assuming the shell to be highly anisotropic. Notably, it enables us to derive some models used in the study of vesicles and RBCs. Part of the results presented in this article have since been proved by a Γ -convergence approach in an Eulerian setting for the justification of the modeling of vesicles by Merlet [Merlet \[2013a,b\]](#). Finally, let us recall and emphasize the fact that the computation of the limit energy should include a relaxation step that is not taken into account in our formal framework. The only interesting case being the one where the flexural term $Q_{F'}^0(G, s)$ is not fully degenerate, that is not independent of G . In such a case, a relaxation of the membrane term of the limit energy is expected to take place. The correct limit energy in Proposition 2 should read

$$I_0'(\psi_0) = \frac{1}{3} \int_{S'} Q_{D'\psi_0}^0(D'n, 1) dx' + 2 \int_{S'} \mathcal{Q}'W_0(D'\psi_0, n) dx' - 2 \int_{S'} f_0 \cdot \psi_0 dx',$$

where $\mathcal{Q}'W_0$ is the in-plane quasiconvexification of W_0 , defined for every element F' of $T^*(S'; \mathbb{R}^3)$ by

$$\mathcal{Q}'W_0(F', n) = \inf_{\varphi \in C_0^\infty(\omega; T_{x'}S')} |\omega|^{-1} \int_{\omega} W_0(F'(\text{Id}' + D'\varphi), n) dy',$$

where $x' = \pi_{S'}(F')$ and ω is a bounded regular open set of $T_{x'}S'$.

NOTATIONS

- \mathbb{R} , set of reals
- \mathbb{R}' , dual set of reals
- \mathbb{N} , set of non-negative integers
- S' , midsurface of the shell in the reference configuration
- ε , thickness of the shell
- $S^\varepsilon := S' \times (-\varepsilon, \varepsilon)$, abstract reference configuration of the shell
- $S := S^1$, rescaled abstract reference configuration of the shell
- \tilde{S}^ε , geometric reference configuration
- \tilde{S} , geometric reference configuration of maximum thickness
- x^ε , element of S^ε
- \tilde{x} , element of \tilde{S}
- x' , element of S'
- TM , tangent space to M
- $T_x M$, tangent fiber to M at x
- T^*M , cotangent space to M
- T_x^*M , cotangent space to M at x
- $T(M; \mathbb{R}^3) := TM \oplus TM \oplus TM$, Whitney triple sum of TM
- $T_x(M; \mathbb{R}^3) := T_x M \oplus T_x M \oplus T_x M$, Whitney triple sum of $T_x M$
- $T^*(M; \mathbb{R}^3) := T^*M \oplus T^*M \oplus T^*M$, Whitney triple sum of T^*M
- $T_x^*(M; \mathbb{R}^3) := T_x^*M \oplus T_x^*M \oplus T_x^*M$, Whitney triple sum of T_x^*M
- $\pi_B : \mathcal{F} \rightarrow B$, projection of a fiber bundle \mathcal{F} onto its base B
- $\pi_{x'} : \mathbb{R}^3 \rightarrow T_{x'} S'$, projection onto $T_{x'} S'$
- ψ^ε , deformation of S^ε
- $\psi(\varepsilon) : S \rightarrow \mathbb{R}^3$, rescaled map of the deformation ψ^ε
- $\psi = (\psi_k)$, expansion of the deformation $\psi(\varepsilon) = \sum_k \varepsilon^k \psi_k$ with respect to the thickness
- $\varphi = (\varphi_k)$, expansion of the minimizers $\varphi(\varepsilon) = \sum_k \varepsilon^k \varphi_k$ with respect to the thickness
- $D\psi^\varepsilon$, differential of $\psi^\varepsilon : S^\varepsilon \rightarrow \mathbb{R}^3$
- $D\psi^\varepsilon(x^\varepsilon)$, differential of $\psi^\varepsilon : S^\varepsilon \rightarrow \mathbb{R}^3$ at x^ε
- $(D'\psi^\varepsilon, D_3\psi^\varepsilon)$, decomposition of $D\psi^\varepsilon \in T^*(S^\varepsilon; \mathbb{R}^3)$ on $T^*(S'; \mathbb{R}^3) \times T^*((-\varepsilon, \varepsilon); \mathbb{R}^3)$ where $\psi^\varepsilon : S^\varepsilon \rightarrow \mathbb{R}^3$
- $(D'\psi^\varepsilon(x), D_3\psi^\varepsilon(x))$, decomposition of $D\psi^\varepsilon(x) \in T_x^*(S^\varepsilon; \mathbb{R}^3)$ on $T_{x'}^*(S'; \mathbb{R}^3) \times T_{x_3}^*((-\varepsilon, \varepsilon); \mathbb{R}^3)$ where $\psi^\varepsilon : S^\varepsilon \rightarrow \mathbb{R}^3$
- ∂_3 , partial differentiation along the fibers
- $J_\varepsilon(\psi^\varepsilon)$, elastic energy of a deformation $\psi^\varepsilon : S^\varepsilon \rightarrow \mathbb{R}^3$
- $I_\varepsilon(\psi^\varepsilon)$, total energy of a deformation $\psi^\varepsilon : S^\varepsilon \rightarrow \mathbb{R}^3$
- $L_\varepsilon(\psi^\varepsilon)$, work of the external loads
- $J(\varepsilon)(\psi(\varepsilon))$, rescaled elastic energy of the deformation ψ^ε
- $I(\varepsilon)(\psi(\varepsilon))$, rescaled total energy of the deformation ψ^ε
- $\mathbf{J}(\varepsilon)(\boldsymbol{\psi})$, elastic energy of the deformation of asymptotic expansion $\boldsymbol{\psi}$
- $\mathbf{I}(\varepsilon)(\boldsymbol{\psi})$, total energy of the deformation of asymptotic expansion $\boldsymbol{\psi}$
- $I_0(\psi_0)$, limit of the total energy (for standard boundary conditions)
- f_ε , external loads
- $W^\varepsilon : T^*(S^\varepsilon; \mathbb{R}^3) \rightarrow \overline{\mathbb{R}}^+$, stored energy
- Ψ^ε , set of admissible deformations of S^ε
- Ψ , set of admissible asymptotic expansions for the deformations of S^ε

- $W_k : T^*(S; \mathbb{R}^3) \rightarrow \mathbb{R}^+$, k -th term of the asymptotic expansion of the stored energy
- $\mathcal{M} \subset T^*(S; \mathbb{R}^3)$, the restriction of the zero set of W_2 to the midsection S'
- $\mathcal{M}' \subset T^*(S'; \mathbb{R}^3)$, projection of \mathcal{M} on $T^*(S'; \mathbb{R}^3)$
- $n_0 : \mathcal{M}' \rightarrow T_0^*((-1, 1); \mathbb{R}^3)$, orientation of the normal fiber in the deformed configuration
- $n : \mathcal{M}' \rightarrow \mathbb{R}^3$, orientation of the normal fiber in the deformed configuration (only the vectorial part)
- n' , normal to the midsection S' of the reference configuration
- D^2W_2 , second derivative of W_2 on \mathcal{M}
- $a[\cdot]^2 = a(\cdot, \cdot)$, where a is a bilinear form
- $Q_F : T^*(S'; \mathbb{R}^3) \times \mathbb{R}^3 \rightarrow \mathbb{R}$, quadratic form associated to the flexural limit energy, where F is a section of \mathcal{M}'
- $Q_F^0 := \inf_v Q_F(\cdot, v) : T^*(S'; \mathbb{R}^3) \rightarrow \mathbb{R}$
- $I_{flex}(\psi)$, flexural limit energy of a deformation ψ
- dx' , 2-dimensional Hausdorff measure restricted to S'
- $\text{Tr}(A)$, the trace of the matrix A
- $\tilde{\psi}^\varepsilon, \tilde{W}^\varepsilon, \tilde{W}_k, \tilde{\mathcal{M}}, \tilde{L}_\varepsilon, \tilde{f}_\varepsilon, \dots$ overtilded variables are defined on the geometric configuration
- $\text{SO}(TM)$, rotations of TM , that is the fiber bundle made of all rotations of $T_x M$ (with $x \in M$), where M is a Riemann manifold
- $\text{SO}(n)$, rotations of \mathbb{R}^n
- $\text{SL}(TM)$, special group of TM , that is the fiber bundle made of all linear diffeomorphism of $T_x M$ of determinant equal to one (with $x \in M$), where M is a Riemann manifold
- \mathcal{E}_M , set of symmetric bilinear forms on the tangent space of M
- \mathcal{O} , fiber bundle over S' whose fibers are the maps from $T_{x'} S'$ into itself of zero trace
- \mathcal{P} , set of polynomials of degree lower than or equal to two on $\mathcal{E}_{S'}$

REFERENCES

- D. D. Fox, A. Raoult, and J. C. Simo. A justification of nonlinear properly invariant plate theories. *Arch. Rational Mech. Anal.*, 124(2):157–199, 1993. ISSN 0003-9527. doi: 10.1007/BF00375134. URL <http://dx.doi.org/10.1007/BF00375134>.
- Hervé Le Dret and Annie Raoult. The nonlinear membrane model as variational limit of nonlinear three-dimensional elasticity. *J. Math. Pures Appl. (9)*, 74(6): 549–578, 1995. ISSN 0021-7824.
- H. Le Dret and A. Raoult. The membrane shell model in nonlinear elasticity: a variational asymptotic derivation. *J. Nonlinear Sci.*, 6(1):59–84, 1996. ISSN 0938-8974. doi: 10.1007/s003329900003. URL <http://dx.doi.org/10.1007/s003329900003>.
- Olivier Pantz. On the justification of the nonlinear inextensional plate model. *Arch. Ration. Mech. Anal.*, 167(3):179–209, 2003. ISSN 0003-9527. doi: 10.1007/s00205-002-0238-1. URL <http://dx.doi.org/10.1007/s00205-002-0238-1>.
- Gero Friesecke, Richard D. James, and Stefan Müller. A theorem on geometric rigidity and the derivation of nonlinear plate theory from three-dimensional elasticity. *Comm. Pure Appl. Math.*, 55(11):1461–1506, 2002. ISSN 0010-3640. doi: 10.1002/cpa.10048. URL <http://dx.doi.org/10.1002/cpa.10048>.
- Gero Friesecke, Richard D. James, Maria Giovanna Mora, and Stefan Müller. Derivation of nonlinear bending theory for shells from three-dimensional nonlinear elasticity by Gamma-convergence. *C. R. Math. Acad. Sci. Paris*, 336(8): 697–702, 2003. ISSN 1631-073X. doi: 10.1016/S1631-073X(03)00028-1. URL [http://dx.doi.org/10.1016/S1631-073X\(03\)00028-1](http://dx.doi.org/10.1016/S1631-073X(03)00028-1).
- Sergio Conti, Francesco Maggi, and Stefan Müller. Rigorous derivation of Föppl’s theory for clamped elastic membranes leads to relaxation. *SIAM J. Math. Anal.*, 38(2):657–680, 2006. ISSN 0036-1410. doi: 10.1137/050632567. URL <http://dx.doi.org/10.1137/050632567>.
- Gero Friesecke, Richard D. James, and Stefan Müller. A hierarchy of plate models derived from nonlinear elasticity by gamma-convergence. *Arch. Ration. Mech. Anal.*, 180(2):183–236, 2006. ISSN 0003-9527. doi: 10.1007/s00205-005-0400-7. URL <http://dx.doi.org/10.1007/s00205-005-0400-7>.
- Sergio Conti and Francesco Maggi. Confining thin elastic sheets and folding paper. *Arch. Ration. Mech. Anal.*, 187(1):1–48, 2008. ISSN 0003-9527. doi: 10.1007/s00205-007-0076-2. URL <http://dx.doi.org/10.1007/s00205-007-0076-2>.
- P.B. Canham. The minimum energy of bending as a possible explanation of the biconcave shape of the human red blood cell. *Journal of Theoretical Biology*, 26(1):61 – 81, 1970. ISSN 0022-5193. doi: 10.1016/S0022-5193(70)80032-7. URL <http://www.sciencedirect.com/science/article/pii/S0022519370800327>.
- Wolfgang Helfrich. Elastic properties of lipid bilayers: theory and possible experiments. *Naturforsch C*, 28(11):693–703, 1973. ISSN 0341-0471.
- E.A. Evans. Bending resistance and chemically induced moments in membrane bilayers. *Biophysical Journal*, 14(12):923 – 931, 1974. ISSN 0006-3495. doi: 10.1016/S0006-3495(74)85959-X. URL <http://www.sciencedirect.com/science/article/pii/S000634957485959X>.
- James T. Jenkins. The equations of mechanical equilibrium of a model membrane. *SIAM J. Appl. Math.*, 32(4):755–764, 1977a. ISSN 0036-1399.

- S. J. Singer and Garth L. Nicolson. The fluid mosaic model of the structure of cell membranes. *Science*, 175(4023):720–731, 1972. doi: 10.1126/science.175.4023.720. URL <http://www.sciencemag.org/content/175/4023/720.abstract>.
- D. J. Steigmann. Fluid films with curvature elasticity. *Archive for Rational Mechanics and Analysis*, 150:127–152, 1999. ISSN 0003-9527. URL <http://dx.doi.org/10.1007/s002050050183>. 10.1007/s002050050183.
- S. Svetina and B. ek. Membrane bending energy and shape determination of phospholipid vesicles and red blood cells. *European Biophysics Journal*, 17:101–111, 1989. ISSN 0175-7571. URL <http://dx.doi.org/10.1007/BF00257107>. 10.1007/BF00257107.
- Udo Seifert, Karin Berndl, and Reinhard Lipowsky. Shape transformations of vesicles: Phase diagram for spontaneous- curvature and bilayer-coupling models. *Phys. Rev. A*, 44(2):1182–1202, Jul 1991. doi: 10.1103/PhysRevA.44.1182. URL <http://link.aps.org/doi/10.1103/PhysRevA.44.1182>.
- Ling Miao, Udo Seifert, Michael Wortis, and Hans-Günther Döbereiner. Budding transitions of fluid-bilayer vesicles: The effect of area-difference elasticity. *Phys. Rev. E*, 49(6):5389–5407, Jun 1994. doi: 10.1103/PhysRevE.49.5389. URL <http://link.aps.org/doi/10.1103/PhysRevE.49.5389>.
- Shankar Krishnaswamy. A cosserat-type model for the red blood cell wall. *International Journal of Engineering Science*, 34(8):873 – 899, 1996. ISSN 0020-7225. doi: 10.1016/0020-7225(95)00139-5. URL <http://www.sciencedirect.com/science/article/pii/0020722595001395>.
- Deuling, H.J. and Helfrich, W. The curvature elasticity of fluid membranes : A catalogue of vesicle shapes. *J. Phys. France*, 37(11):1335–1345, 1976. doi: 10.1051/jphys:0197600370110133500. URL <http://dx.doi.org/10.1051/jphys:0197600370110133500>.
- J. T. Jenkins. Static equilibrium configurations of a model red blood cell. *Journal of Mathematical Biology*, 4:149–169, 1977b. ISSN 0303-6812. URL <http://dx.doi.org/10.1007/BF00275981>. 10.1007/BF00275981.
- Jon C. Luke. A method for the calculation of vesicle shapes. *SIAM Journal on Applied Mathematics*, 42(2):pp. 333–345, 1982. ISSN 00361399. URL <http://www.jstor.org/stable/2101216>.
- Jon C. Luke and Jerome I. Kaplan. On the theoretical shapes of bilipid vesicles under conditions of increasing membrane area. *Biophysical Journal*, 25(1):107 – 111, 1979. ISSN 0006-3495. doi: 10.1016/S0006-3495(79)85280-7. URL <http://www.sciencedirect.com/science/article/pii/S0006349579852807>.
- Ashutosh Agrawal and David Steigmann. Boundary-value problems in the theory of lipid membranes. *Continuum Mechanics and Thermodynamics*, 21:57–82, 2009. ISSN 0935-1175. URL <http://dx.doi.org/10.1007/s00161-009-0102-8>. 10.1007/s00161-009-0102-8.
- Feng Feng and William S. Klug. Finite element modeling of lipid bilayer membranes. *Journal of Computational Physics*, 220(1):394 – 408, 2006. ISSN 0021-9991. doi: 10.1016/j.jcp.2006.05.023. URL <http://www.sciencedirect.com/science/article/pii/S0021999106002476>.
- Andrea Bonito, Ricardo H. Nochetto, and M. Sebastian Pauletti. Parametric FEM for geometric biomembranes. *J. Comput. Phys.*, 229(9):3171–3188, 2010. ISSN 0021-9991. doi: 10.1016/j.jcp.2009.12.036. URL <http://dx.doi.org/10.1016/j.jcp.2009.12.036>.

- A. Bonito, R. H. Nochetto, and M. S. Pauletti. Dynamics of biomembranes: effect of the bulk fluid. *Math. Model. Nat. Phenom.*, 6(5):25–43, 2011. ISSN 0973-5348. doi: 10.1051/mmnp/20116502. URL <http://dx.doi.org/10.1051/mmnp/20116502>.
- Gerhard Dziuk. Computational parametric Willmore flow. *Numer. Math.*, 111(1): 55–80, 2008. ISSN 0029-599X. doi: 10.1007/s00211-008-0179-1. URL <http://dx.doi.org/10.1007/s00211-008-0179-1>.
- Zhangli Peng, Xuejin Li, Igor V. Pivkin, Ming Dao, George E. Karniadakis, and Subra Suresh. Lipid bilayer and cytoskeletal interactions in a red blood cell. *Proceedings of the National Academy of Sciences*, 2013. doi: 10.1073/pnas.1311827110. URL <http://www.pnas.org/content/early/2013/07/24/1311827110.abstract>.
- Qiang Du, Chun Liu, and Xiaoqiang Wang. Simulating the deformation of vesicle membranes under elastic bending energy in three dimensions. *Journal of Computational Physics*, 212(2):757 – 777, 2006. ISSN 0021-9991. doi: 10.1016/j.jcp.2005.07.020. URL <http://www.sciencedirect.com/science/article/pii/S0021999105003566>.
- Qiang Du, Chun Liu, and Xiaoqiang Wang. A phase field approach in the numerical study of the elastic bending energy for vesicle membranes. *Journal of Computational Physics*, 198(2):450 – 468, 2004. ISSN 0021-9991. doi: 10.1016/j.jcp.2004.01.029. URL <http://www.sciencedirect.com/science/article/pii/S0021999104000373>.
- Qiang Du and Jian Zhang. Adaptive finite element method for a phase field bending elasticity model of vesicle membrane deformations. *SIAM J. Sci. Comput.*, 30(3):1634–1657, 2008. ISSN 1064-8275. doi: 10.1137/060656449. URL <http://dx.doi.org/10.1137/060656449>.
- Shravan K. Veerapaneni, Denis Gueyffier, Denis Zorin, and George Biros. A boundary integral method for simulating the dynamics of inextensible vesicles suspended in a viscous fluid in 2D. *J. Comput. Phys.*, 228(7):2334–2353, 2009. ISSN 0021-9991. doi: 10.1016/j.jcp.2008.11.036. URL <http://dx.doi.org/10.1016/j.jcp.2008.11.036>.
- Jin Sun Sohn, Yu-Hau Tseng, Shuwang Li, Axel Voigt, and John S. Lowengrub. Dynamics of multicomponent vesicles in a viscous fluid. *J. Comput. Phys.*, 229(1):119–144, 2010. ISSN 0021-9991. doi: 10.1016/j.jcp.2009.09.017. URL <http://dx.doi.org/10.1016/j.jcp.2009.09.017>.
- Yongsam Kim and Ming-Chih Lai. Simulating the dynamics of inextensible vesicles by the penalty immersed boundary method. *J. Comput. Phys.*, 229(12):4840–4853, 2010. ISSN 0021-9991. doi: 10.1016/j.jcp.2010.03.020. URL <http://dx.doi.org/10.1016/j.jcp.2010.03.020>.
- Yaling Liu, Lucy Zhang, Xiaodong Wang, and Wing Kam Liu. Coupling of Navier-Stokes equations with protein molecular dynamics and its application to hemodynamics. *Internat. J. Numer. Methods Fluids*, 46(12):1237–1252, 2004. ISSN 0271-2091. doi: 10.1002/flid.798. URL <http://dx.doi.org/10.1002/flid.798>.
- Yaling Liu and Wing Kam Liu. Rheology of red blood cell aggregation by computer simulation. *J. Comput. Phys.*, 220(1):139–154, 2006. ISSN 0021-9991. doi: 10.1016/j.jcp.2006.05.010. URL <http://dx.doi.org/10.1016/j.jcp.2006.05.010>.

- Lindsay M. Crowl and Aaron L. Fogelson. Computational model of whole blood exhibiting lateral platelet motion induced by red blood cells. *Int. J. Numer. Methods Biomed. Eng.*, 26(3-4):471–487, 2010. ISSN 2040-7939. doi: 10.1002/cnm.1274. URL <http://dx.doi.org/10.1002/cnm.1274>.
- D. Salac and M. Miksis. A level set projection model of lipid vesicles in general flows. *Journal of Computational Physics*, 230(22):8192 – 8215, 2011. ISSN 0021-9991. doi: 10.1016/j.jcp.2011.07.019. URL <http://www.sciencedirect.com/science/article/pii/S0021999111004384>.
- Emmanuel Maitre, Thomas Milcent, Georges-Henri Cottet, Annie Raoult, and Yves Usson. Applications of level set methods in computational biophysics. *Mathematical and Computer Modelling*, 49(11-12):2161 – 2169, 2009. ISSN 0895-7177. doi: 10.1016/j.mcm.2008.07.026. URL <http://www.sciencedirect.com/science/article/pii/S0895717708002422>.
- V. Doyeux, Y. Guyot, V. Chabannes, C. Prud’homme, and M. Ismail. Simulation of two-fluid flows using a finite element/level set method. Application to bubbles and vesicle dynamics. *J. Comput. Appl. Math.*, 246:251–259, 2013. ISSN 0377-0427. doi: 10.1016/j.cam.2012.05.004. URL <http://dx.doi.org/10.1016/j.cam.2012.05.004>.
- Benoit Merlet. A highly anisotropic nonlinear elasticity model for vesicles I. Eulerian formulation, rigidity estimates and vanishing energy limit. 23 pages, preprint, http://hal.archives-ouvertes.fr/hal-00848547/PDF/Merlet_VesiclesPartI.pdf, July 2013a. URL <http://hal.archives-ouvertes.fr/hal-00848547>.
- Benoit Merlet. A highly anisotropic nonlinear elasticity model for vesicles. II. Derivation of the thin bilayer bending theory. 57 pages, preprint, http://hal.archives-ouvertes.fr/hal-00848552/PDF/Merlet_VesiclesPartII.pdf, July 2013b. URL <http://hal.archives-ouvertes.fr/hal-00848552>.

CMAP, UMR CNRS 7641, ECOLE POLYTECHNIQUE, 91128 PALAISEAU CEDEX, FRANCE

E-mail address: olivier.pantz@polytechnique.org

URL: <http://www.cmap.polytechnique.fr>

DRI, INSTITUT POLYTECHNIQUE DES SCIENCES AVANCÉES, 5-9 RUE MAURICE GRANDCOING, 94200 IVRY-SUR-SEINE, FRANCE

E-mail address: karim.trabelsi@polytechnique.edu

URL: <http://www.ipsa.fr>