



HAL
open science

Consistent time-step optimization in the lattice Boltzmann method

Tobias Horstmann, Hatem Touil, Lucien Vienne, Denis Ricot, Emmanuel Lévêque

► **To cite this version:**

Tobias Horstmann, Hatem Touil, Lucien Vienne, Denis Ricot, Emmanuel Lévêque. Consistent time-step optimization in the lattice Boltzmann method. 2021. hal-03387908v1

HAL Id: hal-03387908

<https://hal.science/hal-03387908v1>

Preprint submitted on 20 Oct 2021 (v1), last revised 24 Aug 2022 (v2)

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.

Consistent time-step optimization in the lattice Boltzmann method

Tobias Horstmann^{1,2,*}, Hatem Touil², Lucien Vienne³, Denis Ricot⁴, and Emmanuel Lévêque³

¹German Aerospace Center, Institute of Propulsion Technology, Department of Engine Acoustics, Bismarckstraße 101, D-10625 Berlin, Germany 5

²CS Group, 22, av. Galilée, F-92350, Le Plessis Robinson, France

³Univ Lyon, Ecole Centrale de Lyon, CNRS, Univ Claude Bernard Lyon 1, INSA Lyon, LMFA, UMR 5509, F-69130, Ecully, France

⁴Renault, Technocentre, 1 av. du Golf, F-78280 Guyancourt, France 10

*Corresponding author. Email address: tobias.horstmann@dlr.de

October 18, 2021

Abstract

Owing to its efficiency and aptitude for a massive parallelization, the lattice Boltzmann method generally outperforms conventional solvers in terms of execution time in weakly-compressible flows. However, in the incompressible limit, the authorized time-step (being inversely proportional to the speed of sound) becomes prohibitively small so that the performance advantage over continuum-based solvers vanishes. A remedy to optimize the time-step is provided by tailoring an artificial speed of sound, which can be fixed or variable throughout the simulation, the latter case referring to an *adaptive time-stepping*. While achieving considerable speed-ups in certain flow configurations, *adaptive time-stepping* comes with the flaw that the continuities of mass density and pressure cannot be fulfilled conjointly when the speed of sound is varied. Therefore, a trade-off is needed. By leaving the mass density unchanged, the conservation of mass is preserved but the pressure presents a discontinuity in the momentum equation. In contrast, a manipulation of the mass density allows us to ensure the continuity of the pressure term in the momentum equation (per unit mass) but the conservation of mass is locally sacrificed. This algorithm requires an additional scaling operation and will therefore be called “*adaptive time-stepping with correction*” in the article. Interestingly, we found that this second trade-off is generally preferable. 15 20 25 30

In the case of a thermal plume, whose movement is governed by the balance of buoyancy and drag forces, the correction of the mass density (to ensure the continuity of the pressure force) has a beneficial impact on the resolved velocity field. In a pulsatile channel flow (Womersley flow) driven by an external body force, no difference was observed between the two versions of *adaptive time-stepping*. 35

On the other hand, if the pulsatile flow is established by inlet and outlet pressure conditions, the results obtained with a continuous pressure force agree much better with the analytical solution. Finally, by using *adaptive time-stepping* in a channel entrance flow, it was shown that the correction is compulsory for the Poiseuille flow to develop. The expected compressibility error due to the discontinuity in the mass density remains small to negligible, and the convergence rate is not affected compared to a simulation with a constant time step.

Keywords— Lattice Boltzmann method, adaptive time-stepping, time-step optimization

1 Introduction

The lattice Boltzmann (LB) method has gained prominence as a statistical approach used to simulate continuum fluid dynamics [Suc15; Krü+17]. Nonetheless, to make it a competitive alternative to conventional solvers, a series of simplifications are necessary that eventually narrow its scope of application. Concretely, the macroscopic equivalent to the standard stream-and-collide LB algorithm reduces to an isothermal, weakly compressible Navier-Stokes model [Del01].

Due to the low symmetry of standard lattices, the LB method can only be used with confidence in the range of Mach number

$$\text{Ma} = \frac{|\mathbf{u}|_{\max}}{c_s} \leq 0.3 \quad (1)$$

where compressible effects may be considered weak to negligible; $|\mathbf{u}|_{\max}$ denotes the maximum flow velocity and c_s is the speed of sound in the fluid. The restriction to isothermal fluids is another consequence of the low symmetry of the lattice, which fails to properly express the conservation of internal energy with a varying temperature [Del01]. Furthermore, the mere philosophy of the LB method conflicts with the concept of an incompressible fluid, devised to omit the computation of compressible effects provided they are negligible. Based on the idea that a gas is composed of particles that are taken into account as local phase densities, incompressibility would imply that these particles travel at an infinite speed, which - even as a mathematical concept - is difficult to reconcile.

A variety of techniques exist to lift the restrictions to low Mach numbers and isothermal fluids in the LB approach [Qia93; CK06; FST15; Li+07]. On the contrary, improvements to address incompressible fluids are rather limited. Incompressible LB models have been postulated in the past but they only decrease the order of compressibility errors in steady flows [HL97; Zou+95; Del03].

In continuum fluid mechanics, the principal motivation behind an incompressible description stems from the numerical benefit of an increased time-step. In general, the maximum time-step is expressed as

$$\Delta t_{\max} = \frac{\text{CFL} \Delta x}{c} \quad (2)$$

where the Courant-Friedrichs-Lewy (CFL) number may be viewed as the normalized maximum velocity at which flow variations can be robustly propagated by the numerical scheme [CFL28]. For compressible Navier-Stokes solvers, the reference

velocity c is the speed of sound relative to the maximum flow velocity, *i.e.*

$$c = c_s + |\mathbf{u}|_{\max} \quad \text{and} \quad \Delta t_{\max} = \frac{\text{CFL } \Delta x}{c_s(1 + \text{Ma})} \quad (3)$$

If one supposes a CFL number of unity and $\Delta x \approx 10^{-3}$ m then Δt_{\max} for air at standard conditions ($c_s \simeq 343$ m/s) is approximately $3 \cdot 10^{-6}/(1 + \text{Ma})$ s yielding very small time-steps. The assumption of a (truly) incompressible fluid removes the speed of sound from the reference velocity c . Therefore, the maximum time-step becomes

$$\Delta t_{\max} = \frac{\text{CFL } \Delta x}{|\mathbf{u}|_{\max}} \quad (4)$$

so that the same CFL number yields considerably larger values at (very) small Mach number. Typically,

$$\Delta t_{\max}^{\text{incomp.}} \simeq \frac{\Delta t_{\max}^{\text{comp.}}}{\text{Ma}}. \quad (5)$$

Conceptually, the sound speed is viewed as infinite in the incompressible model implying that pressure and velocity fields adapt instantaneously. Mathematically, an additional Poisson equation arises (by taking the divergence of the momentum equation) to describe the evolution of pressure (and in turn to project velocity in divergence-free/solenoidal space). One possibility, which is known to belong to the class of SIMPLE algorithms, solves the Poisson equation while the pressure and velocity fields are updated in a (semi-implicit) iterative process [PS72]. Another possibility is the artificial compressibility method (ACM) where a pressure term is substituted into the continuity equation via an isothermal equation of state. This circumvents the iterative stepwise update of the pressure and velocity fields, but reintroduces an artificial speed of sound [Cho67]. In that situation, pseudo-waves propagating with a finite speed are introduced to distribute the pressure.

In the LB method, the distribution functions travel from one lattice node to another during exactly one time-step, *i.e.* with a characteristic speed $c = \Delta x/\Delta t$. On the other hand, the propagation of sound is related to the effective transport of mass-density variations via the distribution functions. Therefore, the speed of sound and the speed of microscopic propagation are physically related (but not strictly equal because the distribution functions also undergo collisions). For a standard isothermal lattice, this relation is

$$c = c_0 \sqrt{3} \quad (6)$$

where $c_0 = \sqrt{p'/\rho'}$ represents an isothermal speed of sound linking the variations of pressure and mass density. As evidenced by Eq. (6), the speed of sound cannot be removed from the statistical approach. Thus, a commonly used technique to accelerate a LB simulation is to artificially decrease c_0 , or equivalently, to increase the compressibility of the fluid. This is the same technique used to maximize the time-step in ACM. In fact the two approaches share a striking similarity [HDC02; Asi+12]. Obviously, the Mach number should not exceed a critical value ($\text{Ma} = 0.3$) to remain in the domain of validity where compressibility effects are weak.

In practice, to determine an appropriate artificial speed of sound, the maximum expected flow velocity is usually overestimated to allow for a certain room of maneuver and to prevent the Mach number from exceeding its maximal allowed value during the simulation. Moreover, in unsteady simulations the maximum velocity

115 may vary by orders of magnitude. As a consequence the time-step of these simula-
 tions is often unnecessarily small. A solution is given by using an adaptive time-step,
 which can be changed throughout a simulation as a function of the current max-
 imum flow velocity. This technique allows for an optimization of the time-step,
 however, it has a side effect on the distribution functions. Being a probability mea-
 120 sure to find a particle with a microscopic velocity c , this probability changes when
 the speed of sound is modified according to Eq. (6). In literature, very few studies
 exist that describe how to handle a change of the speed of sound. Under the name
 of *Mach number annealing*, it is demonstrated in [AHS03] that the Mach number
 may be changed artificially during a simulation while preserving the dynamics of
 125 the flow characterized by a Reynolds number (Re) and a Strouhal number (St). Un-
 fortunately, the reader is deprived of the actual algorithm. This is not the case in
 [Thü+06], where a comprehensive description of the adaptive time-stepping rescal-
 ing operations is provided for the simulation of a gradient-driven free-surface flow.
 In particular, the off-equilibrium part of the distribution functions is reconstructed
 130 from the rescaled Maxwellian distribution in conjunction with a rescaling operation
 that is used in mesh refinement algorithms [FH98]. In [Lat07], the author introduces
adaptive time-stepping for the regularized LB method. Here the populations are re-
 constructed entirely from the rescaled macroscopic variables. The readjustment of
 distribution functions $f(\mathbf{x}, t, \mathbf{c})$ due to a change in \mathbf{c} is common to the aforemen-
 135 tioned studies, whether in dimensional or non-dimensional space. However, the
 continuity of the pressure is not preserved in these algorithms. In the present study
 we will comment on this impact and propose a correction to optimize the use of
adaptive time-stepping for unsteady weakly compressible flows.

140 The paper is organized as follows. Sec. 2 contains a description of the physical
 impact of a sudden change in the speed of sound on the fluid dynamics. Sec. 4
 presents the different algorithms that have been tested in this study. Sec. 5 shows
 results of each algorithm for three different test cases, while concluding remarks are
 given in Sec. 6.

2 Physical aspects of the adaptive time-stepping

145 The following considerations apply to an isothermal fluid satisfying the equation
 of state $rT_0 = c_0^2$ with r being the specific gas constant and c_0 an isothermal
 speed of sound. In this case, the governing gas kinetic equation under the Bhatna-
 gar–Gross–Krook (BGK) approximation [BGK54] reads for the distribution function
 $f(\mathbf{x}, \mathbf{c}, t)$ as

$$\frac{\partial f}{\partial t} + (\mathbf{c} \cdot \nabla) f + \left(\frac{\mathbf{f}_{ext}}{\rho} \cdot \nabla_{\mathbf{c}} \right) f = -\frac{1}{\tau} (f - f^{eq}) \quad (7)$$

150 where the equilibrium distribution f^{eq} is represented by the *Maxwellian distribution*

$$f^{eq}(\mathbf{x}, \mathbf{c}, t) = \frac{\rho(\mathbf{x}, t)}{(2\pi c_0^2)^{3/2}} \exp\left(-\frac{(\mathbf{c} - \mathbf{u}(\mathbf{x}, t))^2}{2c_0^2}\right). \quad (8)$$

\mathbf{f}_{ext} represents a possible external force (per unit volume) acting on the fluid. Fur-
 thermore, we make the assumption that a particular flow is uniquely defined by its
 Reynolds number

$$\text{Re} = \frac{UD}{\nu}$$

and, possibly, its Strouhal number

$$\text{St} = \frac{fD}{U}$$

where U and D are respectively a reference velocity and length scale, f is a frequency of oscillation of the flow and ν is the kinematic viscosity of the fluid. The invariance of the Reynolds and Strouhal numbers implies that U and hence ν must remain unaffected when the speed of sound is changed. In addition, the theory provides a direct relation between the kinematic viscosity and the speed of sound, namely

$$\nu = \tau c_0^2 \quad (9)$$

where τ is the relaxation time that directly intervenes in the BGK collision operator of Eq. (7). It follows that changing the speed of sound requires modifying the relaxation time inversely in order to keep the viscosity constant. The relaxation towards equilibrium thus occurs at a larger time-scale when c_0 is decreased. The physical explanation is straightforward. Under the equation of state $rT_0 = c_0^2$, the speed of sound is directly related to the temperature. As a consequence a sudden drop in c_0 means that molecular activity is abruptly lowered. With the particles acting in a more sedate fashion, the collisional time-scale is increased.

On the other hand, the speed of sound expresses as

$$c_0^2 = \frac{p'}{\rho'} \quad (10)$$

which implies that adaptive time-stepping will lead to discontinuities in either the pressure or the density field. In order to determine the physical consequences of this relation, it is informative to consider (through a Chapman-Enskog multiple-time-scale analysis) the macroscopic equivalent to Eq. (7) up to second-order in the Knudsen and Mach numbers [Suc15; Krü+17], *i.e.*

$$\begin{aligned} \frac{\partial \rho'}{\partial t} + \nabla \cdot (\rho \mathbf{u}) &= 0 \\ \frac{\partial \mathbf{u}}{\partial t} + (\mathbf{u} \cdot \nabla) \mathbf{u} &= -\frac{1}{\rho} \nabla (\rho' c_0^2) + \nu \Delta \mathbf{u} + \mathbf{a} \end{aligned} \quad (11)$$

with $\rho(\mathbf{x}, t) = \rho_{\text{ref}} + \rho'(\mathbf{x}, t)$ and $\mathbf{a} \equiv \mathbf{f}_{\text{ext}}/\rho$. By default, the above Navier-Stokes model is continuous in $\rho(\mathbf{x}, t)$. As a consequence, the pressure force (per unit mass) $-\nabla(\rho' c_0^2)/\rho$ in the momentum equation will be affected by an abrupt change of c_0 . It is possible to derive a pressure-continuous Navier Stokes model, but in this case the density field has to be adapted. More precisely, if the speed of sound is changed from c_0 to

$$c_0^* = \lambda c_0 \quad (12)$$

the continuity of the pressure force per unit mass requires that

$$-\frac{c_0^2}{\rho} \nabla \rho' = -\frac{(\lambda c_0)^2}{\rho^*} \nabla \rho'^* \quad (13)$$

where – here and in the following – the superscript $*$ denotes a quantity rescaled in response to a modified speed of sound. After some calculus (detailed in the Appendix A) we obtain the correction for the density

$$\rho^* = \rho_{\text{ref}} \left(\frac{\rho}{\rho_{\text{ref}}} \right)^{\frac{1}{\lambda^2}}. \quad (14)$$

180 In summary, we note the following changes as a consequence of an artificially
 modified speed of sound: i) Decreasing the speed of sound with the purpose to
 increase the time-step also increases τ , which has a decelerating effect on the relax-
 ation process. ii) By default, the Navier-Stokes model with *adaptive time-stepping*
 is density continuous. Nevertheless, a modification of the density field (Eq. (14))
 185 allows us to render it pressure continuous.

3 The lattice Boltzmann model

Besides the physical consequences of *adaptive time-stepping* on the continuous Boltz-
 mann equation, the discrete model requires some additional rescaling operations. For
 the sake of clarity, we shall start from the dimensional discrete-velocity Boltzmann
 190 equation with external force term

$$\frac{\partial f_\alpha}{\partial t} + (\mathbf{c}_\alpha \cdot \nabla) f_\alpha = -\frac{1}{\tau} (f_\alpha - f_\alpha^{eq}) + F_\alpha, \quad (15)$$

where $\alpha = 1 \cdots q - 1$ spans a discrete set of microscopic velocities (for a $DdQq$
 lattice) and

$$f_\alpha^{eq} = \omega_\alpha \rho \left(1 + \frac{\mathbf{c}_\alpha \cdot \mathbf{u}}{c_0^2} + \frac{(\mathbf{c}_\alpha \cdot \mathbf{u})^2}{2c_0^4} - \frac{u^2}{2c_0^2} + \mathcal{O}(\text{Ma}^3) \right) \quad (16)$$

is the discrete second-order (in Mach number) Maxwellian distribution that is as-
 sociated with isothermal lattice molecules; ω_α are lattice weights. According to
 195 [GZS02] the additional term that arises from the external force (per unit volume)
 \mathbf{f}_{ext} may be written as

$$F_\alpha = \omega_\alpha \left(\frac{\mathbf{c}_\alpha - \mathbf{u}}{c_0^2} + \frac{\mathbf{c}_\alpha (\mathbf{c}_\alpha \cdot \mathbf{u})}{c_0^4} \right) \cdot \mathbf{f}_{ext}. \quad (17)$$

The velocity moments are expressed as discrete sums such that

$$\begin{aligned} \sum_{\alpha=0}^{q-1} f_\alpha &= \sum_{\alpha=0}^{q-1} f_\alpha^{eq} = \rho \\ \sum_{\alpha=0}^{q-1} \mathbf{c}_\alpha f_\alpha &= \sum_{\alpha=0}^{q-1} \mathbf{c}_\alpha f_\alpha^{eq} = \rho \mathbf{u} \\ \sum_{\alpha=0}^{q-1} \mathbf{c}_\alpha \mathbf{c}_\alpha f_\alpha &= \mathbf{P} + \mathcal{O}(\text{Ma}^3) \end{aligned}$$

where \mathbf{P} is the momentum flux tensor. Similarly, the zeroth and first-order velocity
 moments of the forcing term are given by

$$\sum_{\alpha=0}^{q-1} F_\alpha = 0 \quad \sum_{\alpha=0}^{q-1} \mathbf{c}_\alpha F_\alpha = \mathbf{f}_{ext}. \quad (18)$$

200 Concerning the discretization in space and time, the left-hand-side of Eq. (15) is
 viewed as a substantial derivative, which is integrated (for each α) over a time-step

Δt along the characteristics $d\mathbf{x}_\alpha(t) = \mathbf{c}_\alpha dt$. The collision operator and the force term are then approximated by a trapezoidal rule, which leads to a semi-implicit scheme. In order to render this scheme explicit, the following change of variable is applied

$$g_\alpha = f_\alpha + \frac{\Delta t}{2\tau}(f_\alpha - f_\alpha^{eq}) - \frac{\Delta t}{2}F_\alpha. \quad (19)$$

It can be noted that F_α is at least first-order in the Chapman-Enskog multiple-timescale expansion. Therefore, it could formally be added to the non-equilibrium part of the distribution functions. Here, we refrain from this manipulation and treat separately the three contributions to g_α as

$$g_\alpha = g_\alpha^{eq} + g_\alpha^{neq} - \frac{\Delta t}{2}F_\alpha \quad (20)$$

with $g_\alpha^{eq} = f_\alpha^{eq}$ and $g_\alpha^{neq} \equiv (1 + \Delta t/2\tau)(f_\alpha - f_\alpha^{eq})$. In consequence, the moments of g_α^{eq} and g_α^{neq} remain unaffected by the external force, whereas

$$\sum_{\alpha=0}^{q-1} g_\alpha = \rho \quad \text{and} \quad \sum_{\alpha=0}^{q-1} \mathbf{c}_\alpha g_\alpha = \rho \mathbf{u} - \frac{\Delta t}{2} \mathbf{f}_{ext} \quad (21)$$

according to Eq. (18).

The resulting discrete Boltzmann equation then reads

$$\begin{aligned} g_\alpha(\mathbf{x}, t + \Delta t) &= g_\alpha(\mathbf{x} - \mathbf{c}_\alpha \Delta t, t) - \frac{\Delta t}{\tau_g} [g_\alpha(\mathbf{x} - \mathbf{c}_\alpha \Delta t, t) - g_\alpha^{eq}(\mathbf{x} - \mathbf{c}_\alpha \Delta t, t)] \\ &+ \left(1 - \frac{\Delta t}{2\tau_g}\right) \Delta t F_\alpha(\mathbf{x} - \mathbf{c}_\alpha \Delta t, t) \end{aligned}$$

which can here be simplified to

$$\begin{aligned} g_\alpha(\mathbf{x}, t + \Delta t) &= g_\alpha(\mathbf{x} - \mathbf{c}_\alpha \Delta t, t) - \frac{\Delta t}{\tau_g} g_\alpha^{neq}(\mathbf{x} - \mathbf{c}_\alpha \Delta t, t) \\ &+ \Delta t F_\alpha(\mathbf{x} - \mathbf{c}_\alpha \Delta t, t) \end{aligned}$$

with $\tau_g = \tau + \Delta t/2$ and g_α^{neq} defined by Eq. (20).

4 The algorithm

The algorithm that is presented in the following applies to the classical stream-and-collide formulation on an isothermal lattice, *e.g.* D2Q9, D3Q19, D3Q27. In this case, the time-step is given by

$$\Delta t = \frac{\Delta x}{\sqrt{3}c_0}. \quad (22)$$

In the *adaptive time-stepping* algorithm, the speed of sound is tailored in order to maintain a constant target Mach number Ma_t so that

$$c_0^*(t) = \frac{u_{\max}(t)}{\text{Ma}_t} \quad (23)$$

with the maximum velocity $u_{\max}(t) \equiv \sup_{\mathbf{x} \in \mathcal{D}} \|\mathbf{u}(\mathbf{x}, t)\|$ in the simulation domain. In unsteady simulations, this (artificial) numerical speed of sound is therefore a

225 function of time. Nevertheless, at a particular instant, it is constant throughout the fluid and proportional to u_{\max} . This should not be confused with the (real-world) physical speed of sound $c_s(\mathbf{x}, t)$ that may vary in time and space as a result of a non-constant temperature (due to compressibility effects).

The time-step (obtained after tailoring the speed of sound) is then maximal for each iteration according to

$$\Delta t^*(t) = \frac{\text{Ma}_t \Delta x}{\sqrt{3} u_{\max}(t)}. \quad (24)$$

230 In other words, the time-step is dynamically adapted so that the maximal Mach number of the flow reaches a desired target value. By considering the multiplicative factor λ (*cf.* Eq. (12)), we therefore obtain that the time-step is modified inversely as

$$\Delta t^* = \frac{1}{\lambda} \Delta t. \quad (25)$$

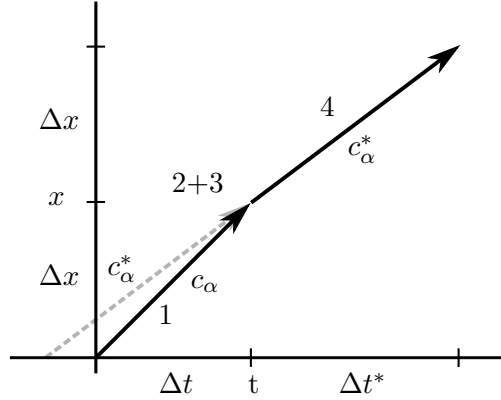


Figure 1: Sketch of the stream-and-collide algorithm with *adaptive time-stepping*. The different stages of the algorithm are (1) streaming (prior rescaling), (2+3) rescaling and collision – (4) streaming (after rescaling).

235 An illustration of the *adaptive time-stepping* algorithm is shown in Fig. 1. The successive stages of the algorithm are

1. Standard streaming of the post-collision state \widehat{g}_α with the original time-step Δt

$$g_\alpha(\mathbf{x}, t) = \widehat{g}_\alpha(\mathbf{x} - \mathbf{c}_\alpha \Delta t, t - \Delta t) \quad (26)$$

2. Rescaling of the distribution functions as if they were streamed with velocity c_α^* (*cf.* the dashed arrow in Fig. 1) plus rescaling of the relaxation time

- 240 3. Collision of the rescaled distributions

$$\widehat{g}_\alpha^*(\mathbf{x}, t) = g_\alpha^*(\mathbf{x}, t) - \frac{\Delta t^*}{\tau_g^*} g_\alpha^{*neq}(\mathbf{x}, t) + \Delta t^* F_\alpha^*(\mathbf{x}, t) \quad (27)$$

4. Streaming of the distribution functions over the adapted time-step Δt^*

$$g_\alpha^*(\mathbf{x} + \mathbf{c}_\alpha^* \Delta t^*, t + \Delta t^*) = \widehat{g}_\alpha^*(\mathbf{x}, t). \quad (28)$$

This sequence can *a priori* be repeated at each or more iterations leading to an optimized time step with respect to a target Mach number.

The remainder of this section will detail the algorithm. It should be noted that the above algorithm uses quantities in dimensional units, whereas the stream-and-collide algorithm is usually solved in non-dimensional lattice units with $\Delta\tilde{x} = \Delta\tilde{t} = 1$. Somewhat against our intuition, the speed of sound in this framework remains constant and equal to $\tilde{c}_0 = 1/\sqrt{3}$. By simple reasoning it thus follows that the fluid velocity $\tilde{\mathbf{u}}$ (entering in the definition of the equilibrium distribution) requires rescaling to keep $\tilde{c}_0 \cdot \text{Ma}_t$ constant when the time-step is adapted. We recall that Ma_t is the target Mach number that shall not be exceeded during a simulation. In other words, the maximum fluid velocity \tilde{u}_{max} (in lattice units) must remain constant.

4.1 The rescaling

The rescaling of g_α is achieved by separately adapting g_α^{eq} , F_α and g_α^{neq} . In lattice units the second-order equilibrium function of the lattice Boltzmann equation reads

$$g_\alpha^{eq} = \omega_\alpha \rho \left(1 + 3(\mathbf{e}_\alpha \cdot \tilde{\mathbf{u}}) + \frac{9}{2}(\mathbf{e}_\alpha \cdot \tilde{\mathbf{u}})^2 - \frac{3}{2}\tilde{u}^2 + \mathcal{O}(\text{Ma}^3) \right) \quad (29)$$

with $\tilde{\mathbf{u}} = \mathbf{u}\Delta t/\Delta x$ and $\mathbf{e}_\alpha = \tilde{\mathbf{c}}_\alpha$ represents the lattice link in the direction α . Changing the time-step from Δt to Δt^* thus changes the non-dimensional fluid velocity. The rescaling of g_α^{eq} is therefore achieved in three steps:

- (i) The computation of the fluid velocity from Eq. (21)

$$\tilde{\mathbf{u}} = \frac{1}{\sum_{\alpha=0}^{q-1} g_\alpha} \left(\sum_{\alpha=0}^{q-1} g_\alpha \mathbf{e}_\alpha + \frac{\tilde{\mathbf{f}}_{ext}}{2} \right)$$

- (ii) The rescaling of the fluid velocity $\tilde{\mathbf{u}}^* = \tilde{\mathbf{u}}/\lambda$ with λ defined by $c_0^* = \lambda c_0$
 (iii) The computation of the rescaled equilibrium function $g_\alpha^{*eq}(\rho^*, \tilde{\mathbf{u}}^*)$. At this point it is possible to consider either that the mass density is not affected, *i.e.* $\rho^* = \rho$, or that the mass density is modified according to Eq. (13) to preserve the continuity of the pressure force (per unit mass).

The contribution of an external force is taken into account similarly. In lattice units, the forcing term expresses as

$$\tilde{F}_\alpha = \omega_\alpha \left(\frac{3}{2}(\mathbf{e}_\alpha - \tilde{\mathbf{u}}) + \frac{9}{2}(\mathbf{e}_\alpha(\mathbf{e}_\alpha \cdot \tilde{\mathbf{u}})) \right) \cdot \tilde{\mathbf{f}}_{ext} \quad (30)$$

In addition to updating the fluid velocity, the external force needs to be modified in the previous equation. This is done by considering

$$\tilde{\mathbf{f}}_{ext}^* = \mathbf{f}_{ext} \frac{(\Delta t^*)^2}{\Delta x} = \frac{1}{\lambda^2} \tilde{\mathbf{f}}_{ext}.$$

The rescaling of g_α^{neq} is not straightforward since its projection onto moment space includes non-hydrodynamic moments whose rescaling is not intuitive. A possible solution would be to reconstruct a regularized non-equilibrium distribution, as suggested in [LC06; Lat07]. Here an alternative to the rescaling by regularization is

270 favored by exploiting the continuity of the viscous stress. Let us mention that the
 rescaling by regularization relies on the continuity of $\sum_{\alpha} g_{\alpha}^{neq} \mathbf{c}_{\alpha} \mathbf{c}_{\alpha}$, which does not
 lead to any particular properties at the macroscopic level. In addition, regulariza-
 tion should be applied to each (local) time step in order not to mix regularized and
 275 non-regularized distributions. This can induce a significant computational overload
 compared to direct rescaling, either when *adaptive time-stepping* is not applied at
 every iteration or when it comes to composite multi-resolution lattices [TRL14].

Before detailing this method, we first establish the rescaling of the non-dimensional
 relaxation time $\tilde{\tau}_g = \tau/\Delta t + 1/2$. Clearly one possibility would be to replace τ by
 Eq. (9) and substitute for Δt^* and c_0^* . However, in order to avoid the return to
 dimensional space it may be preferable to compute $\tilde{\tau}_g^*$ directly from $\tilde{\tau}_g$. Therefore,
 considering that the viscosity must remain unaltered with

$$\nu = \left(\tilde{\tau}_g - \frac{1}{2} \right) \frac{c_0 \Delta x}{\sqrt{3}},$$

and changing the speed of sound from c_0 to λc_0 directly leads to

$$\tilde{\tau}_g^* = \frac{1}{\lambda} \left(\tilde{\tau}_g - \frac{1}{2} \right) + \frac{1}{2}. \quad (31)$$

The rescaling of g_{α}^{neq} can now be processed as followed. A Chapman-Enskog analysis
 establishes that

$$\sum_{\alpha=0}^{q-1} g_{\alpha}^{neq} \mathbf{c}_{\alpha} \mathbf{c}_{\alpha} = -2\rho_{\text{ref}} \tau_g c_0^2 \mathbf{S} + \mathcal{O}(\text{Ma}^3) \quad (32)$$

280 where \mathbf{S} represent the rate-of-strain tensor. Let us remark that the density is here
 approximated by its reference value ρ_{ref} . Since $\rho'/\rho_{\text{ref}} = \mathcal{O}(\text{Ma}^2)$, the resulting
 discrepancy is $\mathcal{O}(\text{Ma}^3)$, which is of the same order as the LB method itself. The
 continuity of $\rho_{\text{ref}} \nu \mathbf{S}$ then gives

$$\frac{\sum_{\alpha=0}^{q-1} g_{\alpha}^{neq} \mathbf{c}_{\alpha} \mathbf{c}_{\alpha}}{\tau_g c_0^2} = \frac{\sum_{\alpha=0}^{q-1} g_{\alpha}^{*neq} \mathbf{c}_{\alpha}^* \mathbf{c}_{\alpha}^*}{\tau_g^* c_0^{*2}} \quad (33)$$

which is equivalent to

$$\frac{3 \sum_{\alpha=0}^{q-1} g_{\alpha}^{neq} \mathbf{e}_{\alpha} \mathbf{e}_{\alpha}}{\tilde{\tau}_g \Delta t} = \frac{3 \sum_{\alpha=0}^{q-1} g_{\alpha}^{*neq} \mathbf{e}_{\alpha} \mathbf{e}_{\alpha}}{\tilde{\tau}_g^* \Delta t^*} \quad (34)$$

285 and eventually yields

$$g_{\alpha}^{*neq} = \frac{1}{\lambda} \frac{\tilde{\tau}_g^*}{\tilde{\tau}_g} g_{\alpha}^{neq} = \frac{1}{\lambda} \left(\frac{1}{\lambda} \left(1 - \frac{1}{2\tilde{\tau}_g} \right) + \frac{1}{2\tilde{\tau}_g} \right) g_{\alpha}^{neq} \quad (35)$$

by assuming that the g_{α}^{neq} 's rescale isotropically.

4.2 Summary of the algorithm

As seen previously, the implementation of *adaptive time-stepping* can be achieved
 in several ways that differ by the physical arguments used. An important issue is

the entanglement of p and ρ via the speed of sound c_0 , which unavoidably introduces a (small) discontinuity in either the mass or momentum transport equation. While previous studies are based on the default setting, which ensures the continuity of ρ but not of p , it is questionable whether this is always a good choice. The main objective of this study is to specifically address this question in regards of numerical accuracy and stability issues. In the following, ATS will refer to *adaptive time-stepping* with the continuity of the density field whereas “ATS with correction” will indicate the continuity of the pressure force by accounting for an additional step (*cf.* Eq. (14)) in the algorithm. An overview of the algorithm employed here is provided in Table 1.

290

295

Table 1: Stream-and-collide algorithm with *adaptive time-stepping*. The different operations are presented in lattice units (indicated by a tilde).

- step (1)** Update g_α via streaming by using Eq. (26)
- step (2)** Compute ρ , $\tilde{\mathbf{f}}_{ext}$, $\tilde{\mathbf{u}}$, and g_α^{eq} to obtain g_α^{neq} by using Eq. (20)
- step (3)** Compute \tilde{u}_{max} and $\lambda = \tilde{u}_{max} / \tilde{c}_0 Ma_t$
- step (4)** Compute ρ^* by using Eq. (14) for ATS with correction
- step (5)** Compute $\tilde{\mathbf{u}}^*$ and $\tilde{\mathbf{f}}_{ext}^*$
- step (6)** Compute g_α^{*eq} and \tilde{F}_α^* by applying Eqs. (29) and (30) with the rescaled variables
- step (7)** Compute $\tilde{\tau}_g^*$ by using Eq. (31)
- step (8)** Compute g_α^{*neq} using Eq. (35) together with g_α^{neq} from step (2)
- step (9)** Compute \tilde{g}_α^* by using Eq. (27)

5 Validation

300

Our first two test cases are chosen to highlight the various benefits of the ATS algorithm. Let us recall that the algorithm presented in this study relies on a fixed Mach number that is defined prior to each simulation. This target Mach number is a trade-off between performance gain and stability. The classical BGK scheme is said to remain stable for Mach numbers up to 0.4 [SC96]. Nevertheless, the time-step does not necessarily have to be at the stability limit for this method to be highly efficient. The strength of the algorithm is to adapt to the current maximum flow velocity. This is of interested for simulations where the maximum velocity is either unknown or varies significantly. Explicit examples are natural convection phenomena in thermal flows and oscillating flows, respectively. Therefore, we present a thermal plume and a Womersley channel flow in two dimensions for both ATS and “ATS with correction” algorithms. The former test case shall provide the reader with an intuitive understanding of this numerical technique, while the latter will examine in more detail the influence of the additional correction step for different simulation setups. In the absence of an oscillating pressure gradient, the Womersley channel flow becomes a Poiseuille channel flow, which constitutes the third test case. Designed as a channel entrance flow with a uniform inlet velocity, it features the standard boundary conditions used for many industrial applications and is worth being assessed in the context of an adaptive time-step. The correction of the density

305

310

315

is found essential in that situation.

5.1 Natural convection in thermal flow

Although an isothermal lattice Boltzmann model is used in this study, it is possible to introduce a temperature-driven buoyancy force into the modelling via the Boussinesq hypothesis

$$\mathbf{f}_b(\mathbf{x}, t) = \rho(\mathbf{x}, t)\mathbf{g}\beta(T(\mathbf{x}, t) - T_0) \quad (36)$$

where β is the coefficient of thermal expansion of the fluid, \mathbf{g} is the gravitational acceleration and T_0 is the temperature at rest. We use a simple two-dimensional Gaussian distribution to introduce a variation in the temperature field, *i.e.*

$$T(\mathbf{x}, t_0) = T_0 + \exp\left(-\frac{x^2 + y^2}{R^2}\right)\Delta T$$

where R is the radius of the hot spot and ΔT is the maximum temperature difference with respect to the background temperature T_0 .

The interest of this test case lies in the time-dependent variation of the flow velocity. Initialized in a fluid at rest, *i.e.* $\mathbf{u}(\mathbf{x}, 0) = \mathbf{0}$, the thermal plume will accelerate until buoyancy and drag forces are balanced. The velocity field is taken into account by the Boltzmann equation with the force \mathbf{f}_b , while the development of the temperature field is governed by a simple advection-diffusion equation

$$\frac{\partial T}{\partial t} + (\mathbf{u} \cdot \nabla) T = \kappa \nabla^2 T \quad (37)$$

where κ is the thermal diffusivity of the fluid (air). The above equation is discretized using a finite-difference approach on the same spatio-temporal grid as the discrete Boltzmann scheme. Commonly referred to as an hybrid LB method, it is here easily brought in line with the *adaptive time-stepping* algorithm. In non-dimensional velocity space, a discrete form of Eq. (37) writes as

$$T(\mathbf{x}, t + 1) = T(\mathbf{x}, t) - (\tilde{\mathbf{u}} \cdot \nabla_h) T + \tilde{\kappa} \Delta_h T,$$

where ∇_h and Δ_h stand for finite-difference gradient and Laplacian operators. Given that the lattice Boltzmann scheme already provides the rescaled bulk velocity $\tilde{\mathbf{u}}^*$, the only modification required is the rescaling of thermal diffusivity, *i.e.* $\tilde{\kappa}^* = \nu \Delta t_n / Pr \Delta x^2$ with the Prandtl number $Pr = \nu / \kappa$. The complete algorithm is illustrated in Fig. 2.

The hot spot (or plume) is initialized in the center of a 100 m \times 200 m periodic domain. The spatial resolution is $\Delta x = 0.4$ m and $R = 50\Delta x$. The initial speed of sound is 343.2 ms⁻¹ and the target Mach number was fixed at $Ma_t = 0.15$. This target Mach number was reached gradually in order to avoid abrupt changes in the speed of sound at the beginning of the simulation. Specifically, we have used a capped decrease of c_0 according to

$$c_0^* = \max\left(\frac{u_{\max}}{Ma_t}, c_0^{0.9}\right). \quad (38)$$

At first, the results of the two *adaptive time-stepping* algorithms are compared to a reference simulation with a constant time-step and for an initial temperature

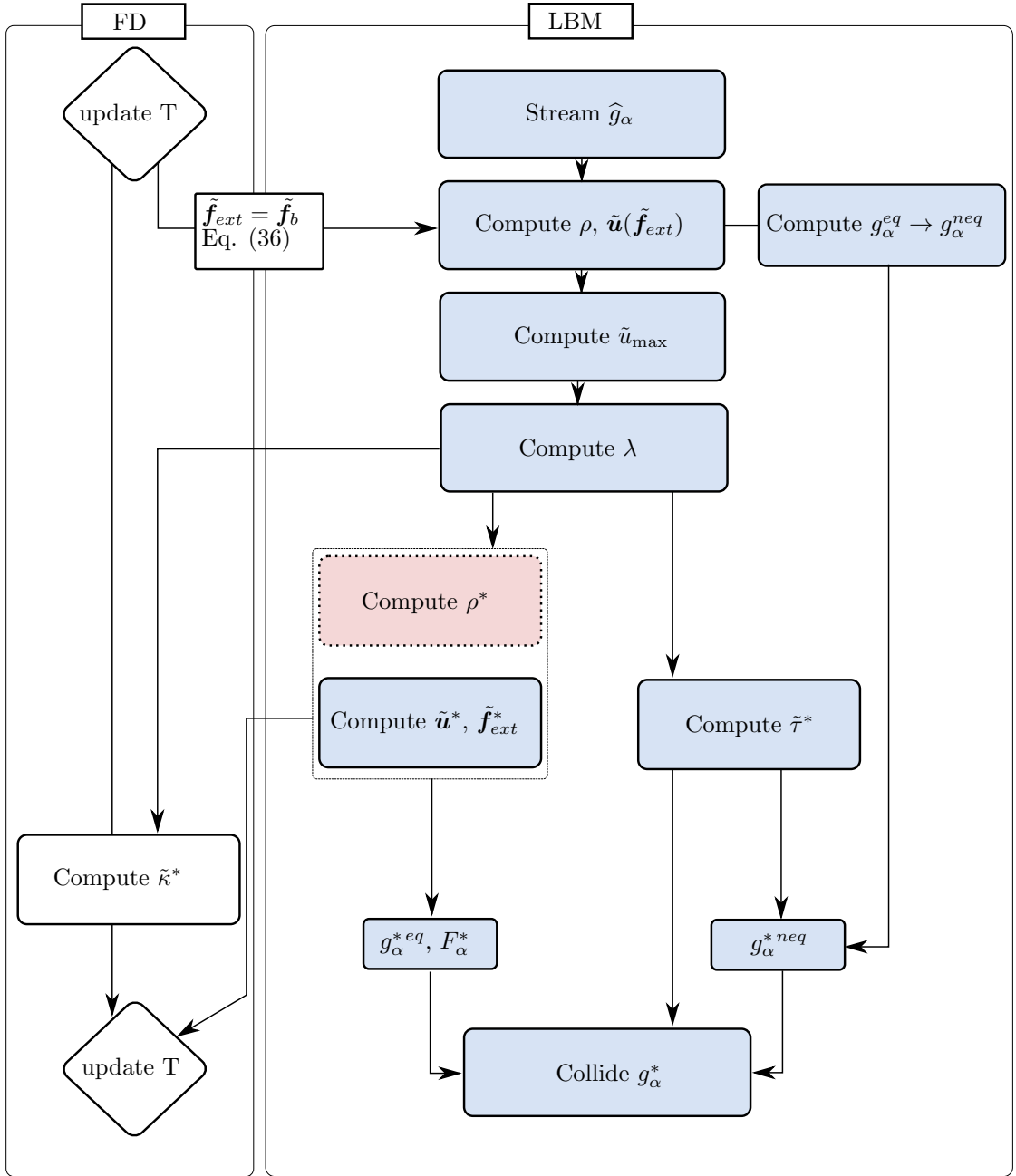


Figure 2: Diagram of the *adaptive time-stepping* (ATS) algorithm in lattice units. Blue boxes coincide with the original algorithm found in [Thü+06] that is extended here by accounting for an external force. White boxes indicate the extension to thermal dynamics using finite differences to solve a advection-diffusion equation for the temperature. Finally, the red box indicates the operation required on the density field (to ensure the continuity of the pressure force) that leads to the “ATS with correction” algorithm.

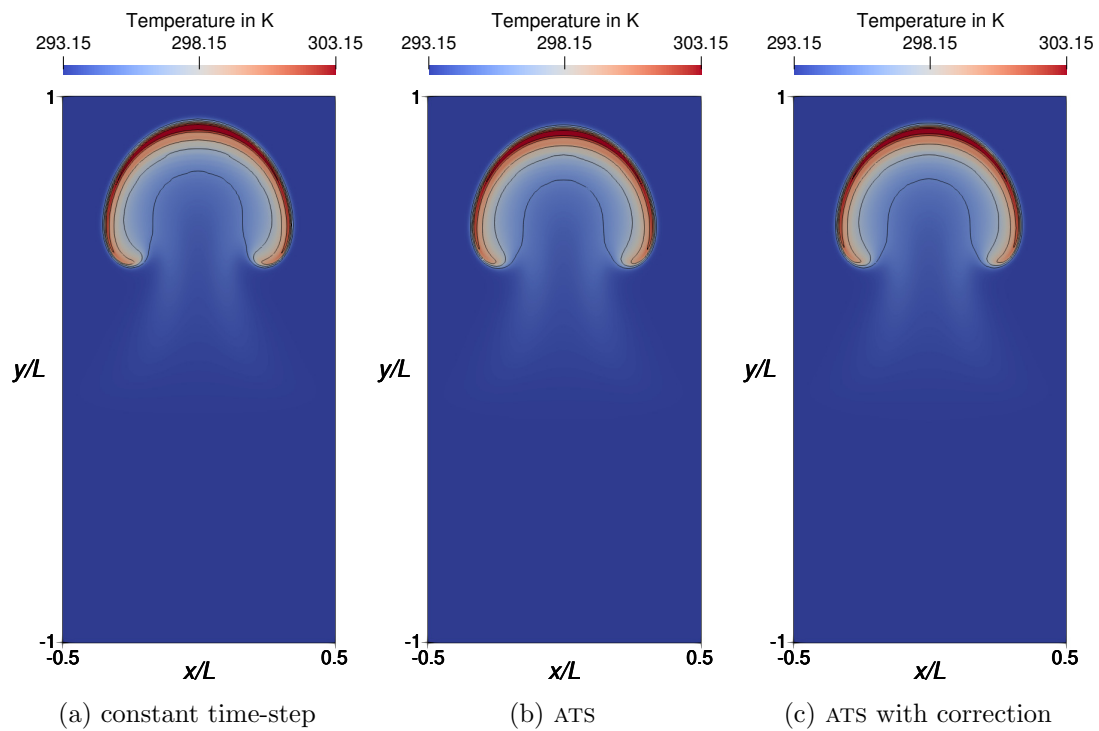


Figure 3: Natural convection of a Gaussian spot of temperature elevation with $\Delta T = 10$ K after 45.3 s. $L = 100$ m and $\Delta x = 0.4$ m. *Left*: Temperature field after 67310 iterations by using a constant time-step of $\Delta t = 6.729 \times 10^{-4}$ s. *Center*: After 3690 iterations by using ATS. *Right*: After 3690 iterations by using ATS with correction. The target Mach number is $\text{Ma}_t = 0.15$.

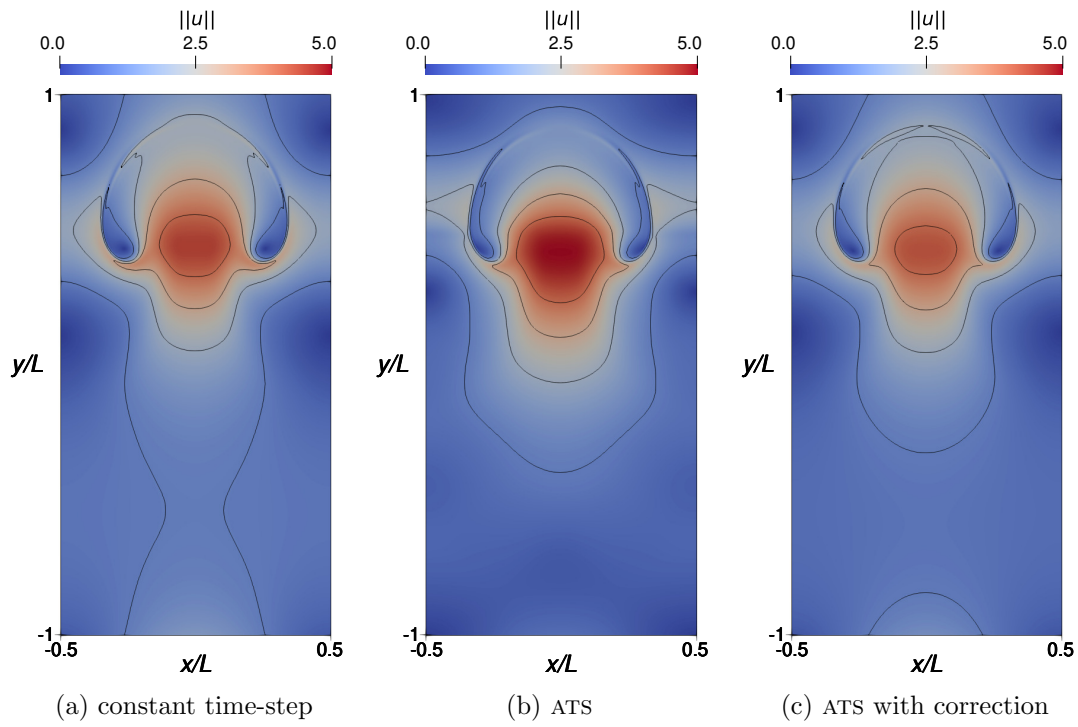


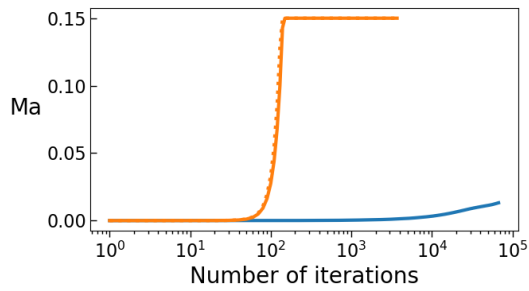
Figure 4: The field of velocity magnitude $\|\mathbf{u}\| = \sqrt{u_x^2 + u_y^2}$ in ms^{-1} is displayed at the same instant as in Fig. 3

345 difference of $\Delta T = 10$ K at a reference temperature of $T_0 = 293.15$ K. The time-step
 is reevaluated every 10 iterations. Due to the variable time-step, the simulations
 are out of sync with each other. At 45.2 seconds, the simulation times match by
 four decimals corresponding to 3690 and 67310 iterations for the adaptive and constant
 350 time-stepping algorithms, respectively. The corresponding temperature fields
 are shown in Fig. 3. Qualitatively, the solutions agree with each other. The plume
 has reached the same elevation in the three simulations despite the (very) different
 time-steps. The contour lines of the temperature field indicate that the deformation
 of the rising plume is suitably captured as well. Let us mention that some
 unavoidable slight differences stem from the artificial compressibility of the acceler-
 355 ated simulations. However, these differences remain here very small at $Ma_t = 0.15$.
 The velocity norms are compared in Fig 4. The maximum velocity is slightly higher
 in case of ATS, whereas the results from ATS with correction better agree with the
 reference (as evidenced later). Since the plume has covered the same distance in the
 three cases, it can only follow that the velocity field oscillate around the reference.

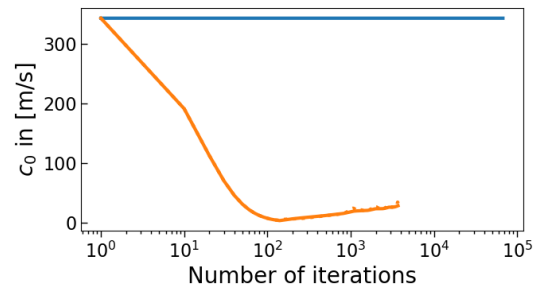
360 The evolution of the parameters related to the time-stepping is shown in Fig. 5.
 In the simulations with ATS, the target Mach number is reached after a transient
 of only about 100 iterations (Fig. 5a), during which the artificial speed of sound
 declines gradually according to the ramp defined in Eq. (38). After this transient,
 the speed of sound starts being controlled by $u_{\max}(t)$ and increases progressively
 365 as the plume rises (Fig. 5b). As a corollary, the time-step adapts itself during the
 transient to reach the target Mach number, then decreases to keep the Mach number
 constant as the flow develops (Fig. 5c). On the contrary, in the simulation with a
 constant time-step, the Mach number remains small reaching a maximum value of
 $Ma = 0.013$ at the end of the run. In Fig. 5c, the difference between the adapted and
 370 reference time-steps directly illustrate the significant gain obtained on the execution
 time for a given physical duration. Finally, let us notice that the ATS algorithm gives
 rise to some oscillatory behavior towards the end of the runs, particularly visible in
 the evolution of the time step. This behavior is much less pronounced when using
 ATS with correction.

375 The robustness of ATS is examined in a second experiment by varying the initial
 temperature difference ΔT between the thermal plume and the surrounding fluid.
 In addition to $\Delta T = 10$ K, $\Delta T = 1$ K and $\Delta T = 100$ K are also considered. Each
 simulation was run for 3690 iterations as previously. The results are expected to
 be very similar with respect to the number of iterations. Qualitatively, a lower
 380 temperature difference leads to a slower rise of the plume, which is compensated by
 a larger time-step. Conversely, a stronger difference has the opposite effect. More
 quantitatively, the dynamics results from the balance between the buoyancy and
 the drag, which justifies that $T - T_0$ varies typically as u . Therefore, the time-
 step adapts itself inversely to $T - T_0$, and the displacement of the plume remains
 385 (almost) constant with respect to the number of iterations. This behavior is correctly
 observed in Fig. 6, where the final temperature fields are displayed.

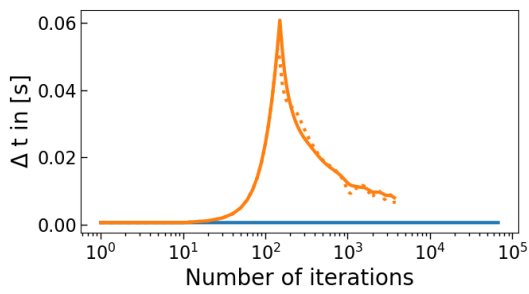
The adaptation of the speed of sound and time-step remain consistent for the
 different dynamics, as evidenced in Fig. 7. Finally, let us notice that the results
 obtained by ATS (without correction) contain once again a wavy perturbation, which
 390 is particularly pronounced in case of a high temperature difference.



(a)



(b)



(c)

Figure 5: Evolution of the parameters related to the time-stepping during the rise of the thermal plume over a duration (physical time) of 45.3 s. (—): constant time-step; (.....): ATS; (—): ATS with correction. Let us notice that the number of iterations is presented on a logarithmic scale.

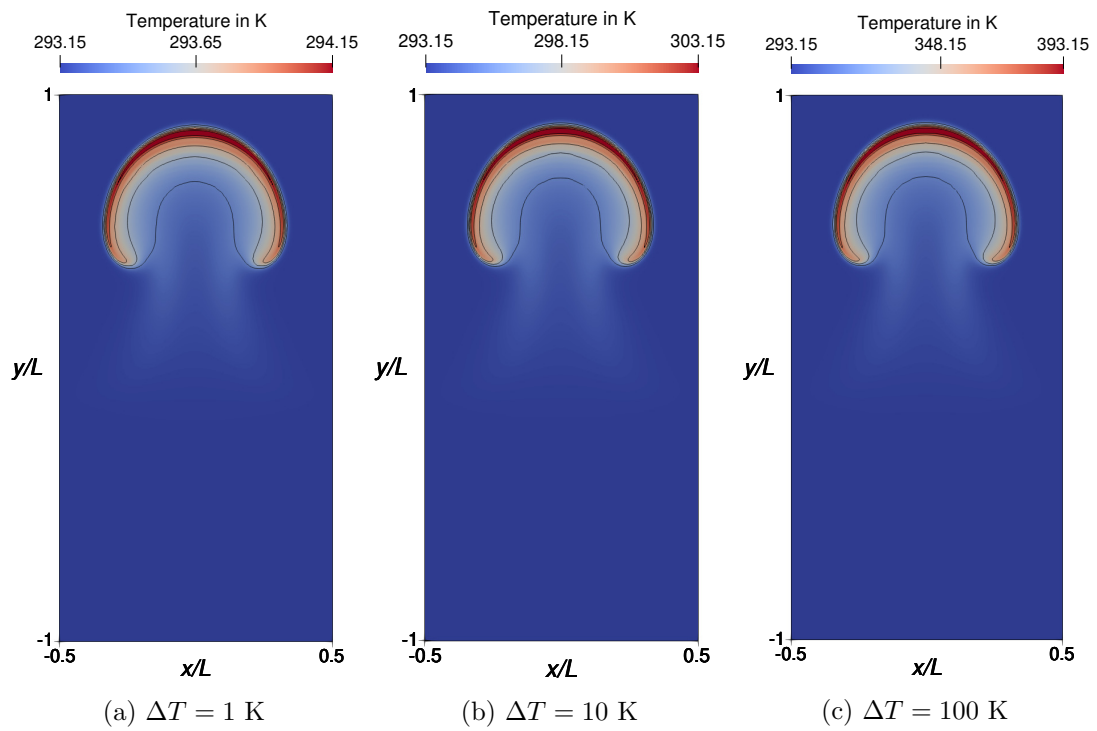


Figure 6: Natural convection of three Gaussian hot spots using ATS with initial temperature differences $\Delta T = 0$ K, 10 K and 100 K, after 3690 iterations. Iso-countours correspond for each figure to $T_0 + j\Delta T$ with $j = 0.2, 0.4, 0.6, 0.8$ and 1.

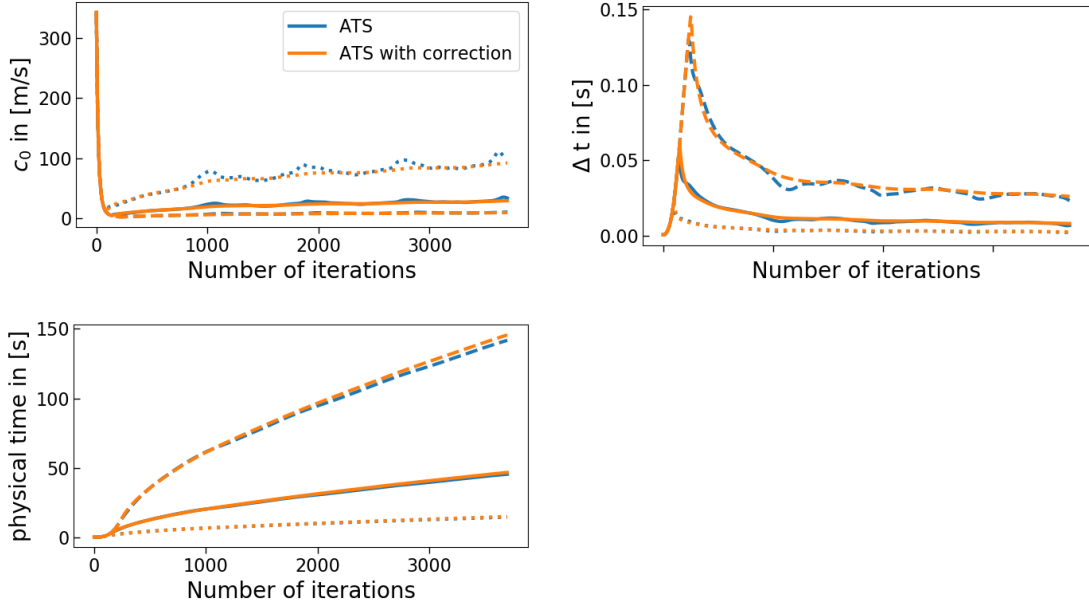


Figure 7: Evolution of the parameters related to the adaptive time-stepping for various initial temperature differences of the thermal plume: (---): $\Delta T = 1$ K; (—): $\Delta T = 10$ K; (.....): $\Delta T = 100$ K.

5.2 Womersley flow

The Womersley flow is a pulsating flow in a 2D channel. In contrast to a Poiseuille flow, the pressure gradient oscillates according to

$$\frac{\partial P}{\partial x} = A \cos(\omega t) \quad (39)$$

where A is the amplitude and ω the pulsation. The latter is related to the pulsating period T as

$$\omega = 2\pi f = \frac{2\pi}{T}. \quad (40)$$

This problem has an exact solution in the laminar regime [IGC13] as the flow stays parallel to the walls and only depends on time and on the wall normal coordinate. Precisely,

$$\begin{aligned} u_x(y, t) &= \Re \left\{ i \frac{A}{\rho \omega} \left(1 - \frac{\cos(\lambda(\frac{2y}{L_y} - 1))}{\cos(\lambda)} \right) e^{i\omega t} \right\} \\ u_y &= 0 \end{aligned} \quad (41)$$

where L_y is the channel height and λ is related to the so-called Womersley number α with

$$\lambda^2 = -i\alpha^2 \quad \text{and} \quad \alpha^2 = \frac{L_y^2 \omega}{4\nu}.$$

400

Similar to the Reynolds number, α relates inertial to viscous forces. Velocity profiles for different Womersley numbers α at constant amplitude A and pulsation ω are shown in Fig. 8. The velocity is normalized with respect to the maximum velocity U_0 achieved in a Poiseuille flow, *i.e.* for $\omega = 0$. For low values of α , the velocity

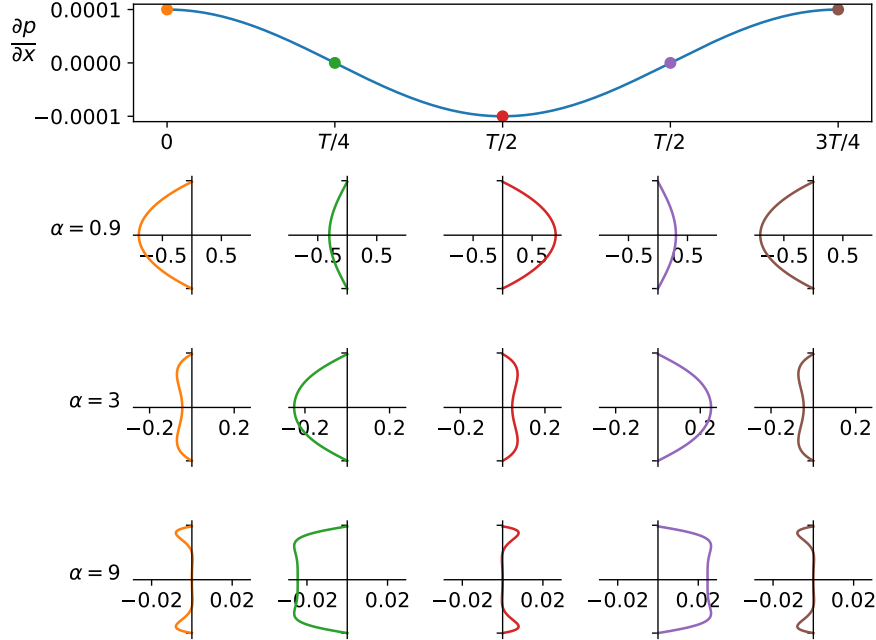


Figure 8: Temporal evolution of the pressure gradient (top) and the corresponding velocity profile $u_x(y, t)/U_0$ for different values of the Womersley number α .

405

profiles remain parabolic. The maximum velocity of the Womersley flow is achieved at the center-line of the channel and is of the same order of magnitude as U_0 . With increasing α , the reluctance of the fluid to change direction increases so that it lags behind the pressure gradient. As a result the maximum velocity decreases and its occurrence moves towards the walls of the channel.

410

415

This multi-parametric problem offers many different possibilities to be implemented. In contrast to the previous example, there exists a maximum velocity for each configuration, which can be obtained from the analytical solution of the flow. For the sake of simplicity, we shall restrict our consideration to Womersley numbers that yields an overall maximum flow velocity located at the center-line of the channel, *i.e.* $U_{\max} = \max\{u_x(L_y/2, t) : 0 \leq t < T\}$. At this point, it is important to distinguish between the overall maximum velocity U_{\max} that is achieved during a period T , and the current maximum velocity $u_{\max}(t) = \max\{u_x(y, t) : 0 \leq y \leq L_y\}$. According to Fig. 8, $u_{\max}(t)$ can become zero at small α , which would result here in an infinitely large time-step. We assume that the Womersley number is sufficiently high to avoid this pitfall.

420

Our ATS algorithm requires a target Mach number Ma_t as an input parameter. In that case, \tilde{c}_0 and consequently \tilde{u}_{\max}^* (after rescaling) remain constant. In a non-dimensional framework, the scaling factor λ can therefore directly be obtained as

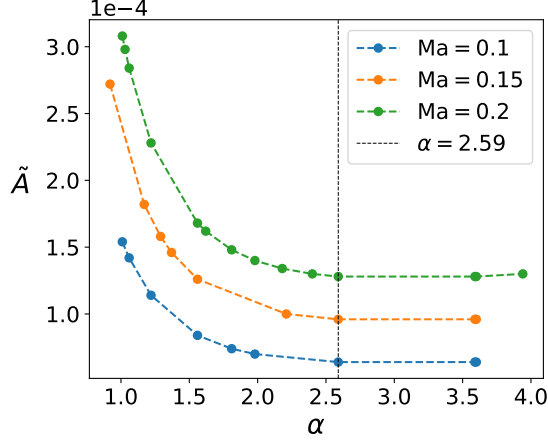


Figure 9: The non-dimensional amplitude \tilde{A} as a function of the Womersley number α for different Mach numbers.

$\lambda = \tilde{u}_{\max}/\tilde{U}_{\max}$ so that $\tilde{u}_{\max}^* = \tilde{u}_{\max}/\lambda = \tilde{U}_{\max}$.

In order to select the three target Mach numbers $\text{Ma}_t = 0.1, 0.15$ and 0.2 , the remaining parameters are now selected such that $\tilde{U}_{\max} = 0.1\tilde{c}_0, 0.15\tilde{c}_0$ and $0.2\tilde{c}_0$, respectively. In accordance with [Nas+14], $Re = 100$ and $\tilde{T} = 5000$ to mimic real life flow phenomena that can be encountered in the smaller arteries of the human body. Next, we determine \tilde{A} and α . According to Eq. (41), both parameters influence the velocity field and thus \tilde{U}_{\max} . In the $(\tilde{A} - \alpha)$ space, the isolines corresponding the three target Mach numbers (with $\rho = 1$ and $\tilde{L}_y = 48$) are plotted in Fig. 9. In the light of our preceding considerations, we picked numerically the matching pairs for $\alpha = 2.59$ with an error tolerance of 0.02% on the Mach number. Finally, the viscosity is given by

$$\nu = \frac{\omega L_y^2}{4\alpha^2}. \quad (42)$$

In the case of Wormerley flow, the pressure gradient can either be established by an external body force, *i.e.*

$$\mathbf{f}_{ext}(t) = -A \cos(\omega t) \mathbf{e}_x \quad (43)$$

or by setting the respective inlet and outlet conditions with, by definition,

$$\frac{p_{in}(t) - p_{out}(t)}{L_x} = c_0^2 \frac{\rho_{in}(t) - \rho_{out}(t)}{L_x} = -A \cos(\omega t).$$

By keeping the averaged density equal to the reference density, *i.e.*

$$\frac{\rho_{in}(t) + \rho_{out}(t)}{2} = \rho_{ref}$$

we eventually obtain the following boundary conditions for the mass density

$$\rho_{in}(t) = \rho_{ref} - \frac{AL_x}{2c_0^2(t)} \cos(\omega t) \quad \text{and} \quad \rho_{out}(t) = \rho_{ref} + \frac{AL_x}{2c_0^2(t)} \cos(\omega t). \quad (44)$$

For the sake of simplicity, the two strategies will be referred to as “body-force” and “inlet/outlet” from here onward. While the former applies directly a pressure gradient to every node of the fluid, the latter enforces the density at the inlet and outlet and induce the pressure gradient through the momentum equation. Therefore, it is expected that the correction of the pressure will have an impact on the results obtained by the inlet/outlet strategy. The simulation parameters are summarized in Table 2.

\tilde{A}	α	Re	Ma_t	$\tilde{T} = 2\pi/\tilde{\omega}$	$\tilde{T}_{\text{sound}} = \tilde{L}_x/\tilde{c}_0$
6.4×10^{-5}	2.59	100	0.1	5000	26
9.6×10^{-5}	2.59	100	0.15	5000	26
1.28×10^{-4}	2.59	100	0.2	5000	26

Table 2: Parameters and characteristics of the simulations (in lattice units). \tilde{T} is the period of oscillation of the pressure gradient. For comparison, \tilde{T}_{sound} is the time required by a local perturbation of density to propagate through the channel. The condition $\tilde{T}_{\text{sound}} \ll \tilde{T}$ ensures that the pressure field has time to equilibrate for each value of the oscillating forcing.

The domain size of the simulations is $(n_x \times n_y) = (24 \times 48)$ and the total number of iterations of each simulation is $5\tilde{T}$. The initial state is set to equilibrium with a zero velocity field and a reference density of unity. The same values are used to impose non-slip boundary conditions for the top and bottom of the domain [ZH97]. The initial speed of sound in all simulations is $U_{\text{max}}/\text{Ma}_t$. The *adaptive time-stepping* strategy is activated at $\tilde{t} = t/T = 0.75$ (after a brief transient to allow the oscillating flow to stabilize) and applied at every iteration.

Fig. 10 compares the evolution of the speed of sound and the maximum Mach number, as well as the flow velocity recorded at the center point $(L_x/2, L_y/2)$ of the domain. Only results for the intermediate Mach number ($\text{Ma}_t = 0.15$) are presented for both the body-force and the inlet/outlet strategies. The results obtained for the two other target Mach numbers were very comparable.

In the non-adaptive case, the speed of sound remains constant at $c_0 = U_{\text{max}}/\text{Ma}_t$. When *adaptive time-stepping* is used, c_0 oscillates between $U_{\text{max}}/\text{Ma}_t$ and a minimum value, which is identical (within 5%) for the two forcing strategies (see Figs. 10a and 10b). As a corollary, the Mach number remains constant when ATS is used, whereas it oscillates when the time-step is kept constant (see Figs. 10c and 10d). The lowest value that $u_{\text{max}}(t)$ attains during a cycle is approximately $U_{\text{max}}/10$. Finally, the non-dimensional velocity \tilde{u}_x at the center point of the domain normalized by \tilde{U}_{max} is shown in Figs. 10e and 10f. In the absence of adaptive time-stepping, \tilde{u}_x closely follows a sinusoidal curve. Once the Mach number is locked to optimize the time-step, $\tilde{u}_x/\tilde{U}_{\text{max}}$ behaves as a periodic step function. Therefore, the maximum gain attainable for the Womersley flow can be approximated by the ratio between the area of a step function and the half-period of a sine function, *i.e.* $\pi/2 \approx 1.57$. With the same number of iterations, the number of periods simulated with the proposed ATS algorithms is increased by a factor of 1.53, which is very close to the optimal gain.

So far, the forcing strategy (body-force or inlet/outlet conditions) and the correction of the pressure field have not revealed any notable differences in the results.

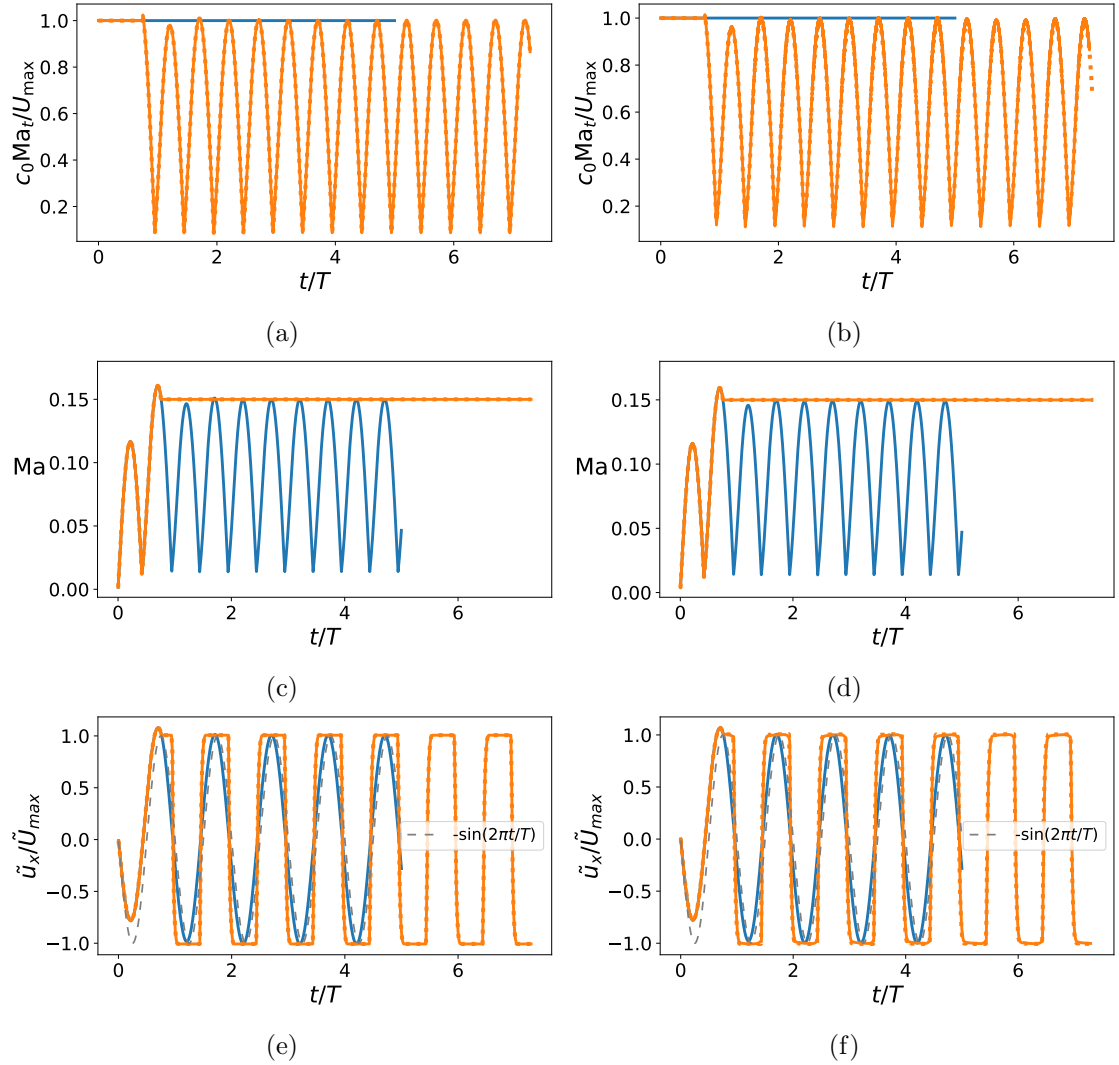


Figure 10: Temporal evolution of non-dimensional parameters related to the *adaptive time-stepping* for the Womersley flow. The target Mach number is $Ma_t = 0.15$. Recall that U_{\max} is the (constant) overall maximum velocity and $u_x(t)$ refers to the probed velocity at the center of the oscillating flow. The Mach number is $Ma = u_{\max}(t)/c_0$ when c_0 is the speed of sound. *Left column*: body-force stirring; *Right column*: inlet/outlet density conditions. (—): constant time-step; (.....): ATS; (—): ATS with correction.

In order to gain quantitative insights about the accuracy of our ATS algorithms, we measured the L_2 -norm of the (normalized) velocity error defined as

$$\varepsilon_u(t) = \frac{\sqrt{\sum_{y_j=1}^{n_y} (u_x(L_x/2, y_j, t) - \bar{u}_x(y_j, t))^2}}{U_{max}} \quad (45)$$

where \bar{u}_x is the analytical solution given by Eq. (41).

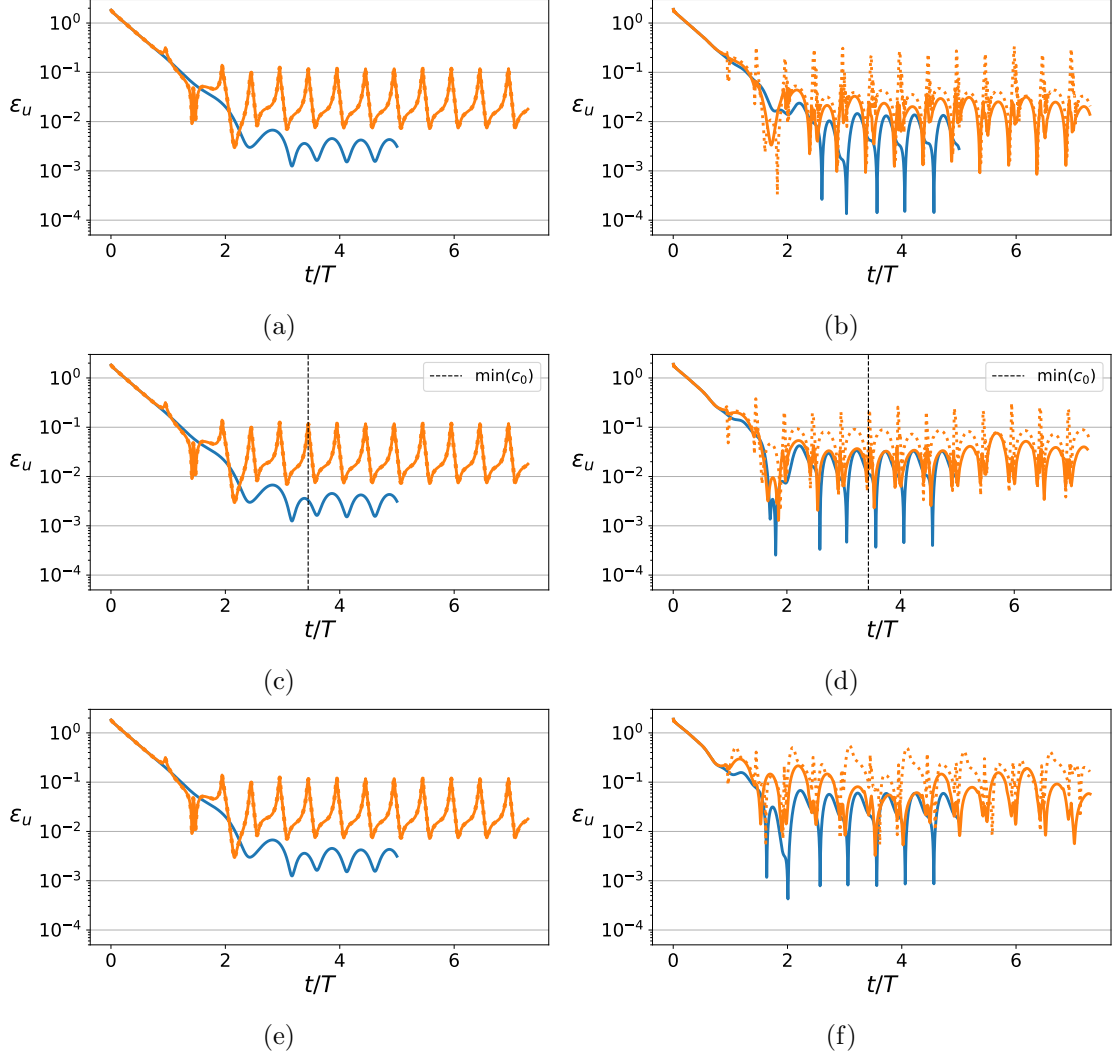


Figure 11: Temporal evolution of the L_2 -norm of the velocity error for standard and adaptive time-stepping at different Mach numbers. *Top row:* $Ma_t = 0.1$; *Middle row:* $Ma_t = 0.15$; *Bottom row:* $Ma_t = 0.2$. *Left column:* body-force stirring; *Right column:* inlet/outlet density conditions. (—): constant time-step; (·····): ATS; (—): ATS with correction.

Fig. 11 shows the error with respect to t/T for the different Mach numbers and forcing strategies. In all cases, a transient error decays rapidly over the first two cycles. When the flow is controlled via a (uniform) body force (see Figs. 11a, 11c and 11e), the error of the *adaptive time-stepping* does not depend on the Mach number and the correction of the pressure has no influence. In that case, the pressure gradient is mainly monitored by the body force and not by the density variations. The error is roughly one order of magnitude higher than the one achieved with a constant time-step, and peaks when the speed of sound is at a local minimum, *i.e.* during the flow reversal. The situation is different when considering an inlet/outlet density (or pressure) conditions. In that case, the error increases with the Mach number, which is understandable since the flow equilibrates through pressure waves. Importantly though, ATS is significantly improved when a correction of the pressure field is applied. As a result, the error remains of the same order as the error encountered in a simulation with a constant time-step, while being about 1.5 times faster. During the flow reversal, the error spikes but does not necessarily exceed the deviation from the analytical solution attained throughout the rest of the cycle.

475

480

485

490

5.3 Channel entrance flow

The previous observations led us to examine a final numerical experiment that specifically addresses the issue of the pressure correction. By using the same channel height (L_y) and Reynolds number as in the previous setup, the length of the channel was increased by a factor 20 to simulate a channel entrance flow [Dur+05]. This flow scenario is achieved with a Dirichlet boundary condition for the velocity and the pressure at the inlet and the outlet, respectively. Measurements are performed at the position $x = 0.8L_x$ where the flow reaches a Poiseuille parabolic velocity profile. In order to improve the convergence and also to make the use of *adaptive time-stepping* relevant, the following ramp (in time) was considered for the velocity at the inlet

$$\begin{aligned}
 U_{\text{in}}(t) &= \sin\left(\frac{\pi}{2} \frac{t}{T_{\text{ramp}}}\right) U_{\text{bulk}} & \text{for } t \leq T_{\text{ramp}} \\
 U_{\text{in}}(t) &= U_{\text{bulk}} & \text{for } t > T_{\text{ramp}}
 \end{aligned} \tag{46}$$

with $U_{\text{bulk}} = 0.75U_0$ and $\tilde{T}_{\text{ramp}} = 10000$ being a fifth of the total number of iterations \tilde{T}_{tot} . Let us mention that U_{bulk} was fixed such that the maximum velocity at $x = 0.8L_x$ is $U_{\text{max}} = U_0$ in the asymptotic steady regime, *i.e.* when $t \gg T_{\text{ramp}}$. Initially, the velocity in the channel is zero and the pressure is uniform at an arbitrary reference value. This reference pressure is kept as the outlet boundary condition during the simulation. Finally, to avoid an impractically small speed of sound in the ATS algorithms, the time-step was locked until the Mach number at the inlet was greater than 0.03. Fig. 12 summarizes the time evolution of the Mach number, the normalized speed of sound and the L_2 -norm of the velocity error defined by Eq. (45) with the Poiseuille profile $\bar{u}(y) = 4U_0(L_y - y)y/L_y^2$. At the beginning, the Mach number (at $x = 0.8L_x$) follows the Mach number ramp imposed at the inlet in the three simulations. When *adaptive time-stepping* is enabled, the Mach number jumps by construction to the target value $\text{Ma}_t = 0.15$. Conjointly, the parabolic Poiseuille profile begins to develop so that the Mach number in the standard algorithm starts to deviate from the Mach number at the inlet (U_{bulk}/c_0) to converge to the asymptotic value $\text{Ma}_t = U_0/c_0$. The speed of sound remains constant in the standard

495

500

505

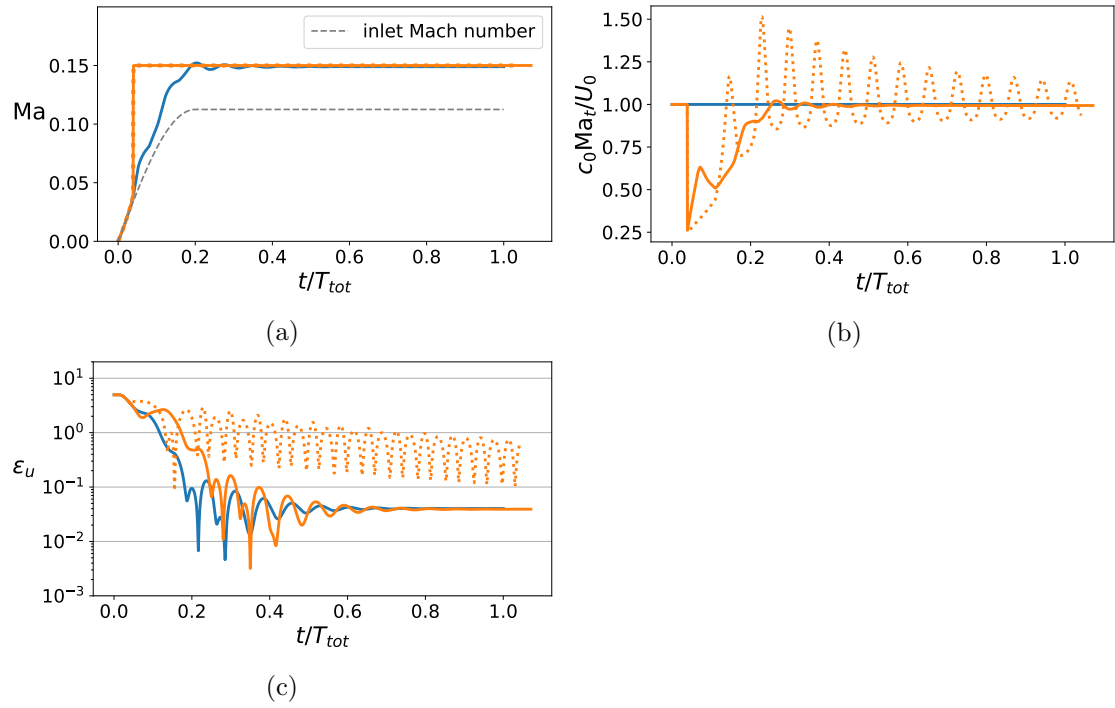


Figure 12: Time evolution of the Mach number (a), normalized speed of sound (b) and L_2 -norm of the velocity error (c) for standard and adaptive time-stepping in a channel entrance flow simulation. Signals are recorded at $x = 0.8L_x$. (—): constant time-step; (.....): ATS; (—): ATS with correction.

approach, while it drops instantly once it is unlocked in the ATS simulations. When the pressure correction is used, the speed of sound consistently recovers its initial value rapidly. On the contrary, without the correction of the pressure, the speed of sound overshoots and enters an oscillatory pattern that slowly stabilizes but remains present throughout the remainder of the simulation. Such important artefact does not occur when the pressure correction is applied. Comparing the error, it becomes apparent that the “ATS with correction” simulation remains consistent with the constant time-step simulation, while the ATS simulation without the pressure correction does not converge at all. The normalized velocity profiles at $x = 0.8L_x$ at the end of

510

515

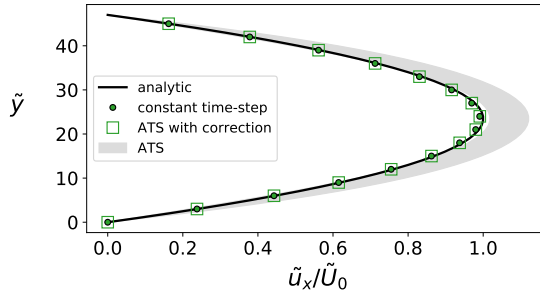


Figure 13: Normalized velocity profiles across the channel at $x = 0.8L_x$ obtained with constant and adaptive time-stepping. The grey area indicates the velocity interval, over which the solution of the ATS algorithm oscillates during the last period, *i.e.* between $t/T_{tot} = 1.013$ and $t/T_{tot} = 1.03$.

the simulation are displayed in Fig. 13). We observe that the results obtained with the standard and “ATS with correction” algorithms match the analytical Poiseuille profile. On the other hand, without the correction of the pressure, the discrepancy is much higher.

520

6 Conclusion

This study presents a physically consistent framework for *adaptive time-stepping* in the LB method, which constitutes an interesting technique to speed-up nearly-incompressible unsteady simulations. Despite being employed in a handful of studies, it is somewhat surprising that this technique has not been used to a wider extent by the LB community.

525

A critical aspect of *adaptive time-stepping* is the inevitable introduction of an error in either the mass or the momentum conservation due to an indirect entanglement of density and pressure variations through the time-step. A natural reluctance to violate the conservation of mass is probably one of the reasons why in previous studies the continuity of the mass density was preferred, and the error always passed into the momentum equation. Here, we investigate the opposite by preserving the pressure force (per unit mass) $-c_0^2 \nabla(\log \rho)$ while sacrificing the continuity of the density field. Interestingly, it was shown that even in the case of natural convection, where the two driving forces are buoyancy and drag, the results are better when the error is attributed to the mass density. In channel flows with classical velocity inlet and pressure outlet conditions, the correction of the pressure is essential for

530

535

convergence. Otherwise, some oscillatory artefacts appear in the velocity field that greatly deteriorate the accuracy of the solution.

540 A physical explanation for these observations is not straightforward. Nevertheless, an indication may be given by comparing the errors made either on the density or the pressure force when varying the time step. Let us first introduce the small parameter

$$\varepsilon = \frac{\Delta t^* - \Delta t}{\Delta t}$$

545 quantifying the (relative) change of the time-step. The continuity of the density field then yields a relative error on the pressure force that varies as

$$\left(\frac{c_0^{*2} - c_0^2}{c_0^2} \right) \sim \varepsilon.$$

On the other hand, the continuity of the pressure force leads to a relative error on the density field

$$\left(\frac{\rho^* - \rho}{\rho} \right) \sim \varepsilon \frac{\rho - \rho_{\text{ref}}}{\rho} \sim \varepsilon \text{Ma}^2.$$

550 For a small Mach number, this relative error is therefore much smaller than the one obtained by keeping the density continuous. Also, if ε is smaller than Ma (which is a valid assumption) the error introduced by “ATS with correction” remains below the error in $\mathcal{O}(\text{Ma}^3)$ of the LB method itself. This provides a plausible justification for the advantage of considering the continuity of the pressure force in the *adaptive time-stepping*.

555 To our best knowledge, this is the first study that proposes a correction of the pressure by accepting small compressibility errors in the solution. In that case, we found that the impact on the fluid dynamics remains negligible. Such algorithm should widen the range of possible applications such as biological flows, transient thermal flows and oscillating flows in general, where the maximum velocity undergoes large variations.

560 Acknowledgment

This work is part of a scientific collaboration including CS-GROUP, Renault, Airbus, Ecole Centrale de Lyon, CNRS and Aix-Marseille University, and it is related to the development of the software ProLBTM.

References

- 565 [AHS03] A. M. Artoli, A. G. Hoekstra, and P. M.A. Sloot. “Accelerated lattice BGK method for unsteady simulations through Mach number annealing”. In: *International Journal of Modern Physics C* 14.6 (2003), pp. 835–845.
- [Asi+12] Pietro Asinari et al. “Link-Wise Artificial Compressibility Method”. In: *J. Comput. Phys.* 231.15 (2012), pp. 5109–5143.
- 570

- [BGK54] P. L. Bhatnagar, E. P. Gross, and M. Krook. “A model for collision processes in gases. I. Small amplitude processes in charged and neutral one-component systems”. In: *Physical Review* 94.3 (1954), pp. 511–525.
- [CFL28] R Courant, K Friedrichs, and H Lewy. “Über die partiellen Differenzgleichungen der mathematischen Physik”. In: *Mathematische Annalen* 100.1 (1928), pp. 32–74. 575
- [Cho67] Alexandre Joel Chorin. “A numerical method for solving incompressible viscous flow problems”. In: *Journal of Computational Physics* 2.1 (1967), pp. 12–26. 580
- [CK06] Shyam S. Chikatamarla and Iliya V. Karlin. “Entropy and Galilean invariance of lattice Boltzmann theories”. In: *Physical Review Letters* 97.19 (2006), pp. 1–4.
- [Del01] Paul J Dellar. “Bulk and shear viscosities in lattice Boltzmann equations”. In: *Physical Review E* (2001). 585
- [Del03] Paul J Dellar. “Incompressible limits of lattice Boltzmann equations using multiple relaxation times”. In: *Journal of Computational Physics* 190 (2003), pp. 351–370.
- [Dur+05] F. Durst et al. “The development lengths of laminar pipe and channel flows”. In: *Journal of Fluids Engineering* 127.6 (2005), pp. 1154–1160. 590
- [FH98] O. Filippova and D. Hänel. “Grid refinement for lattice-BGK models”. In: *Journal of Computational Physics* 147 (1998), pp. 219–228.
- [FST15] Yongliang Feng, Pierre Sagaut, and Wenquan Tao. “A Three Dimensional Lattice Model for Thermal Compressible Flow on Standard Lattices”. In: *Journal of Computational Physics* (2015). 595
- [GZS02] Zhaoli Guo, Chuguang Zheng, and Baochang Shi. “Discrete lattice effects on the forcing term in the lattice Boltzmann method”. In: *Physical Review E* 65 (2002), pp. 1–6. 600
- [HDC02] Xiaoyi He, Gary D. Doolen, and T. Clark. “Comparison of the lattice Boltzmann method and the artificial compressibility method for Navier-Stokes equations”. In: *Journal of Computational Physics* 179.2 (2002), pp. 439–451.
- [HL97] Xiaoyi He and Li-shi Luo. “Lattice Boltzmann Model for Incompressible Navier-Stokes Equation”. In: *Journal of Statistical Physics* 88 (1997). 605
- [IGC13] I.G. Currie. *Fundamental Mechanics of Fluids*. 4th. CRC Press, 2013.

- 610 [Krü+17] Timm Krüger et al. *The lattice Boltzmann method*. 1st ed. Springer, Cham, 2017.
- [Lat07] Jonas Latt. “Hydrodynamic limit of lattice Boltzmann equations”. PhD thesis. Université de Genève, 2007.
- 615 [LC06] Jonas Latt and Bastien Chopard. “Lattice Boltzmann method with regularized pre-collision distribution functions”. In: *Mathematics and Computers in Simulations* 72 (2006), pp. 165–168.
- [Li+07] Q. Li et al. “Coupled double-distribution-function lattice Boltzmann method for the compressible Navier-Stokes equations”. In: *Physical Review E* 76.5 (Nov. 2007).
- 620 [Nas+14] Rupert W. Nash et al. “Choice of boundary condition for lattice-Boltzmann simulation of moderate-Reynolds-number flow in complex domains”. In: *Physical Review E - Statistical, Nonlinear, and Soft Matter Physics* 89.2 (2014), pp. 1–13.
- [PS72] S. V. Patankar and D. B. Spalding. “A calculation procedure for heat, mass and momentum transfer in three-dimensional parabolic flows”. In: *International Journal of Heat and Mass Transfer* 15.10 (1972), pp. 1787–1806.
- 625 [Qia93] Y. H. Qian. “Simulating thermohydrodynamics with lattice BGK models”. In: *Journal of Scientific Computing* 8.3 (1993), pp. 231–242.
- [SC96] James D Sterling and Shiyi Chen. “Stability Analysis of Lattice Boltzmann Methods”. In: *Journal of Computational Physics* 123 (1996), pp. 196–206.
- [Suc15] Sauro Succi. “Lattice Boltzmann 2038”. In: *EPL (Europhysics Letters)* 109.5 (2015).
- 635 [Thü+06] Nils Thürey et al. “Optimization and stabilization of LBM free surface flow simulations using adaptive parameterization”. In: *Computers and Fluids* 35.8-9 (2006), pp. 934–939.
- [TRL14] Hatem Touil, Denis Ricot, and Emmanuel Lévêque. “Direct and large-eddy simulation of turbulent flows on composite multi-resolution grids by the lattice Boltzmann method”. In: *Journal of Computational Physics* 256 (2014), pp. 220–233.
- 640 [TR14] Hatem Touil, Denis Ricot, and Emmanuel Lévêque. “Direct and large-eddy simulation of turbulent flows on composite multi-resolution grids by the lattice Boltzmann method”. In: *Journal of Computational Physics* 256 (2014), pp. 220–233.
- [ZH97] Qisu Zou and Xiaoyi He. “On pressure and velocity boundary conditions for the lattice Boltzmann BGK model”. In: *Physics of Fluids* 9.6 (1997), pp. 1591–1596.
- 645 [Zou+95] Qisu Zou et al. “A improved incompressible lattice Boltzmann model for time-independent flows”. In: *Journal of Statistical Physics* 81.1-2 (Oct. 1995), pp. 35–48.

A Continuity of the pressure force

The continuity of the pressure force per unit mass in the momentum equation expresses as

$$-\frac{c_0^2}{\rho} \nabla \rho' = -\frac{(\lambda c_0)^2}{\rho^*} \nabla \rho'^*$$

with $\rho = \rho_{\text{ref}} + \rho'$ and $\rho^* = \rho_{\text{ref}} + \rho'^*$ by definition. This equation may be rewritten as

$$\frac{1}{\rho^*} \nabla \rho^* = \frac{1}{\lambda^2} \frac{1}{\rho} \nabla \rho$$

which yields the general solution

$$\log(\rho^*) = \frac{1}{\lambda^2} \log(\rho) + C.$$

If $\rho' = \rho'^* = 0$ then $\rho = \rho^* = \rho_{\text{ref}}$. Therefore $C = \ln(\rho_{\text{ref}}) - (1/\lambda^2) \log(\rho_{\text{ref}})$, which eventually gives

$$\rho^* = \rho_{\text{ref}} \left(\frac{\rho}{\rho_{\text{ref}}} \right)^{\frac{1}{\lambda^2}}.$$

If one assumes that the Mach number is small, ρ' and $\rho'^* \ll \rho_{\text{ref}}$. Therefore, $\rho'^* \approx (1/\lambda^2)\rho'$ as a first-order approximation.

650