Viscointertial regime of immersed granular flows

The MIT Faculty has made this article openly available. Please share how this access benefits you. Your story matters.

Citation

As Published
http://dx.doi.org/10.1103/PhysRevE.96.012901

Publisher
American Physical Society

Version
Final published version

Accessed
Sat Aug 26 21:11:42 EDT 2017

Citable Link
http://hdl.handle.net/1721.1/110589

Terms of Use
Article is made available in accordance with the publisher's policy and may be subject to US copyright law. Please refer to the publisher's site for terms of use.

Detailed Terms
Viscoinertial regime of immersed granular flows

L. Amarsid,1,2,3,* J.-Y. Delenne,4,1 P. Mutabaruka,5,1 Y. Monerie,2,3,4 F. Perales,1,6,‡ and F. Radjai2,3,5,§

1IRSNS, PSN, CE Cadarache, BP3-13115 St Paul-Lez-Durance Cedex, France
2LMGC, CNRS, University of Montpellier, 163 rue Auguste Broussonnet, 34090 Montpellier, France
3Laboratoire MIST, IRSN-CNRS, University of Montpellier, France
4IATE, UMR1208 INRA-CIRAD, University of Montpellier–SupAgro, 34060 Montpellier, France
5 < MSE >2, UMI CNRS-MIT, Massachusetts Institute of Technology, 77 Massachusetts Avenue, Cambridge, Massachusetts 02139, USA
6Laboratoire MIST, IRSN-CNRS, University of Montpellier 2, France

(Received 13 February 2017; published 5 July 2017)

By means of extensive coupled molecular dynamics–lattice Boltzmann simulations, accounting for grain dynamics and subparticle resolution of the fluid phase, we analyze steady inertial granular flows sheared by a viscous fluid. We show that, for a broad range of system parameters (shear rate, confining stress, fluid viscosity, and relative fluid-grain density), the frictional strength and packing fraction can be described by a modified inertial number incorporating the fluid effect. In a dual viscous description, the effective viscosity diverges as the inverse square of the difference between the packing fraction and its jamming value, as observed in experiments. We also find that the fabric and force anisotropies extracted from the contact network are well described by the modified inertial number, thus providing clear evidence for the role of these key structural parameters in dense suspensions.

DOI: 10.1103/PhysRevE.96.012901

I. INTRODUCTION

The rheology of granular materials immersed in a viscous fluid is fundamental for modeling many natural flows and industrial applications [1–3]. The flow can be driven by particle motions, as in submarine avalanches, or by the fluid, like in sediment transport in a river. The immersed flow of non-Brownian particles is characterized by its packing fraction $\Phi$, and bulk friction coefficient $\mu = \sigma_f / \sigma_s$, where $\sigma_f$ is shear stress and $\sigma_s$ is the normal stress applied on the granular phase. The fluid affects granular flow by its viscosity $\eta_f$ in combination with the shear rate $\dot{\gamma}$ and $\sigma_s$ via the dimensionless viscous number $I_v = \eta_f \dot{\gamma} / \sigma_s$, which represents the ratio of the Stokes time $t_S = \eta_f / \sigma_s$ and flow time $t_f = \dot{\gamma}^{-1}$ [4]. Hence, the rheology of an immersed granular material can be described by two functions $\mu(I_v)$ and $\Phi(I_v)$, both tending smoothly to their critical values $\mu_0$ and $\Phi_0$ as $I_v \to 0$.

This frictional description has its dual viscous description in which the packing fraction $\Phi$ is the control parameter, and the flow is described by effective normal and shear viscosities $\eta_n$ and $\eta_t$, respectively [4], defined by $\sigma_n = \eta_n \dot{\gamma}$ and $\sigma_t = \eta_t \dot{\gamma}$, where $\sigma_n$ is the internal normal stress of the granular packing. We note that $\sigma_n$ can be different from the applied normal stress $\sigma_s$ due to dynamic effects or in the presence of a suspending fluid, as we shall see below.

The functions $\eta_n(\Phi)$ and $\eta_t(\Phi)$ are classically determined from flow experiments at imposed volume, but can also be deduced from $\mu(I_v)$, $\Phi(I_v)$, and the definition of $I_v$ by assuming $\sigma_n = \sigma_s$. This yields $\eta_n(\Phi) = \eta_f(I_v(\Phi)) / I_v(\Phi)$ and $\eta_t(\Phi) = \eta_f(I_v(\Phi)) / I_v(\Phi)$. This shows that the effective viscosities diverge when $\Phi \to \Phi_0$. The only difference between pressure-imposed and volume-imposed flows is that in the latter case the particles cease to flow at the jamming packing fraction $\Phi_0$, whereas in the former the flow continues. This mapping between the two descriptions was used by Boyer et al. to get experimentally as close as possible to the jamming packing fraction $\Phi_0$ by means of a pressure-controlled flow and to deduce the algebraic divergence of $\eta_n$ and $\eta_t$, as $[\Phi(0) - \Phi]^2$ [4].

The issue with $I_v$ as control parameter is that the limit values $\mu_0$ and $\Phi_0$ do not characterize only the material but depend on the inertial number $I = \dot{\gamma} d (\rho_s / \sigma_s)^{1/2}$, where $d$ is the mean particle diameter and $\rho_s$ is particle density [5]. This means that, in scaling effective properties of the flow, $I_v$ is not simply a replacement for $I$ in the presence of a fluid, as suggested by Boyer et al. [4], but the effective flow properties are controlled by both $I$ and $I_v$. Trulsson et al. partially addressed this issue by molecular dynamics (MD) simulations with a fluid drag force applied to all particle centers [6]. Their results seem to support the idea that a general control parameter combining $I$ and $I_v$ can account for both inertial and viscous effects in a unified framework. However, since the fluid was introduced only through its drag force effect, it is unclear whether this framework can account for other major fluid parameters such as density and volume effects, which are crucial in the dense limit, as well as for lift forces induced by shear. Indeed, the relative density $r = \rho_s / \rho_f$, where $\rho_f$ is fluid density, controls the flow regime [7]. In the same way, the fluid volume between particles carries negative or positive dynamic pore pressures, which in the dense regime should strongly affect particle dynamics. Hence, the framework presented by Trulsson et al. may simply be a consequence of drag forces directly introduced at the particle level. This means that the visco-inertial flow needs to be investigated by means of fully resolved simulations of the fluid phase.

In this paper, we analyze the effective flow behavior in the visco-inertial regime using extensive MD simulations

*Corresponding author. 
†lhassan.amarsid@umontpellier.fr
‡jean-yves.delenne@supagro.inra.fr
§pmutabar@mit.edu
¶yann.monerie@umontpellier.fr
∥frederic.perales@irsn.fr
‡franck.radjai@umontpellier.fr; fradjai@mit.edu

MSE >2, UMI CNRS-MIT, Massachusetts Institute of Technology, 77 Massachusetts Avenue, Cambridge, Massachusetts 02139, USA

Laboratoire MIST, IRSN-CNRS, University of Montpellier 2, France
In order to allow the fluid to flow between pores, we add at least one fluid node to the contact zones between each pair of particles. The node plays the same role as constrictions between pores in three dimensions (3D). The permeability of the structure is controlled by the number of fluid nodes added in contact zones. A brief description of LBM is given in Appendix A. A uniform velocity gradient \( \dot{\gamma} \) is applied to the fluid nodes at domain boundary to generate a plane Couette flow. The fluid is initially at rest and the shear rate \( \dot{\gamma} \) is increased gradually from zero to its final value. The above boundary conditions allow for a homogeneous velocity field of the particles driven by the fluid.

All simulations analyzed in this paper were performed with 1253 particles and a well-resolved fluid with lattice step \( \leq 0.03d \). The friction coefficient \( \mu_s \) is set to 0.4 between the particles and to 0 with the top wall. Gravity is set to zero. The bottom wall is made rough by sticking a layer of particles to the bottom of the box. The fluid viscosity \( \eta_f \) is varied in the range \([\eta_w, 2500\eta_w]\), with water viscosity \( \eta_w \), shear rate \( \dot{\gamma} \) in the range \([0.28, 5.6]\) s\(^{-1}\), confining pressure \( \sigma_s \) in the range \([20, 120]\) Pa, and relative density \( r \) in the range \([0.5, 3]\). The data presented below are average values in the steady state with their error bars representing the standard deviation of fluctuations. The control volume contains the whole volume below the top mobile wall excluding five particle layers adjacent to the walls. Indeed, the velocity profiles are nearly linear except in the five layers close to the walls. Video samples of the simulations can be found in Ref. [11].

### III. SCALING WITH MODIFIED INERTIAL NUMBER

Figure 2 displays \( \mu \) and \( \Phi \) as a function of both \( I \) and \( I_v \). We see that while \( I \) and \( I_v \) cover a wide range, none of them scales the data points over the whole range. The observed partial correlations of the data reflect either the variation of \( \eta_f \) alone or only that of \( \dot{\gamma} \) and \( \sigma_s \). An important effect that may influence the scaling is the presence of shear-induced lift forces [12]. This leads to a weak stress gradient across the sample so that the vertical pressure acting on the solid phase is different.
from the applied stress $\sigma_t$. This can be analyzed by calculating the stress tensor from the forces between the particles. In the data presented in Fig. 2 and in all figures presented below, this effect is accounted for by replacing $\sigma_t$ by $\sigma_t = \sigma_t + \sigma_L$, where $\sigma_L$ is the average vertical stress induced by lift forces. Up to this weak correction of the confining stress, the query is whether a single dimensionless parameter, possibly combining $I$ and $I_v$, can be found to scale the data for all system parameters.

In order to recombine all parameters in a single meaningful variable, we consider the characteristic stresses acting on particles. Note that from each characteristic stress $\Sigma$, a characteristic time $T = d(\rho_i/\Sigma)^{1/2}$ can be defined. The system involves three stresses of different origins:

1. the static stress $\sigma_s$,
2. the viscous stress $\sigma_v \sim \eta_f \dot{\gamma}$,
3. the inertial stress $\sigma_i \sim \rho_i (d \dot{\gamma})^2$.

The corresponding characteristic times are $t_s = d(\rho_i/\sigma_s)^{1/2}$, $t_v = d(\rho_i/\eta_f \dot{\gamma})^{1/2}$, and $t_i = \dot{\gamma}^{-1}$, respectively. We also have the Stokes time $ts = \eta_f/\sigma_s$, which is simply a function of the three other times: $t_0 = t_s(t_i/t_s)^2$. From these three time scales, two independent dimensionless numbers can be built: $I = t_s/t_i = (\sigma_s/\sigma_i)^2$ and $J = t_s/t_v = (\sigma_v/\sigma_s)^{1/2} = (\eta_f \dot{\gamma}/\sigma_s)^{1/2}$. The latter is simply the square root of the Stokes number. The ratio $St = (t_s/t_i)^{1/2} = \sigma_i/\sigma_v = I^2/J^2 = \rho_i d^2 \dot{\gamma}/\eta_f$ is the Stokes number.

The total shear stress $\tau$ in steady flow is the sum of static, inertial, and viscous stresses, and the visco-inertial regime is, by definition, governed by the sum of viscous and inertial stresses in comparison to the static stress. Hence, the natural combination characterizing the visco-inertial regime is $(\alpha_i \sigma_i + \alpha_v \sigma_v)/\sigma_s = \alpha_i I^2 + \alpha_v J^2 = I^2 (\alpha_i + \alpha_v/\text{St})$, where $\alpha_i$ and $\alpha_v$ are constant parameters to be determined. Taking the square root of this ratio, we get a modified inertial number

$$I_m = I \left(\alpha_i + \alpha_v/\text{St}\right)^{1/2},$$

which can be called the visco-inertial number. If this number is equivalent to the inertial number $I$ when $\text{St} \to \infty$, we get $\alpha_i = 1$. Hence, the parameter space of the visco-inertial regime can be reduced only if the functions $\Phi(I_m)$ and $\mu(I_m)$ are uniquely defined for a constant value of $\alpha_v$ that reflects the details of interactions at the particle level.

Figure 3 shows $\mu$ and $\Phi$ as a function of $I_m$ for $\alpha_v = 2.0$, which allows for a remarkable collapse of all our simulation data up to a weak statistical error. Not only do all system parameters affect the rheology only through $I_m$, but we also observe that, consistent with the laminar nature of flow, the data points for different values of the relative density $r$ fall also on the same curve. The parameter $r$ was absent from the simulations of Trulsson et al., and this observation provides strong support for a unified scaling of immersed and dry granular flows.

The functional forms that excellently fit the two plots are

$$\mu(I_m) = \mu_c + \frac{\delta \mu b}{1 + b/I_m},$$

$$\Phi(I_m) = \frac{\Phi_c}{1 + a I_m},$$

where $\mu_c = 0.280 \pm 0.002$ is the quasistatic bulk friction coefficient (in contrast to $\mu_0$, which is a function of $I$), $\Phi_c = 0.8123 \pm 0.0003$ is the steady-state packing fraction (in contrast to $\Phi_0$ which is a function of $I$), $b = 0.246 \pm 0.008$, $\delta \mu = 0.783 \pm 0.014$, and $a = 0.750 \pm 0.003$. Note that the scaling of $\mu$ and $\Phi$ with $I_m$ is strictly the same as in dry granular flows. This is a consequence of the definition of $I_m$ and the choice of $\alpha_i = 1$.

IV. EFFECTIVE VISCOSITIES

In the dual description of the flow, we can express effective viscosities as a function of $\Phi$ using the relations (2) and (3). In the presence of particle-inertial forces in addition to viscous forces, $\sigma_i$ and $\sigma_n$, at given $\Phi$ should scale with $\sigma_i + \alpha_i \sigma_n$:

$$\sigma_i = c_i (\sigma_i + \alpha_i \sigma_n) \equiv c_i \sigma_n I_m^2,$$

$$\sigma_n = c_n (\sigma_i + \alpha_i \sigma_n) \equiv c_n \sigma_n I_m^2,$$

where $c_i$ and $c_n$ play the roles of dimensionless effective viscosities in the visco-inertial regime. We see that they are given by $c_n = 1/I_m^2$ and $c_i = \mu/I_m^2$, leading to the following analytic expressions readily deduced from (2) and (3):

$$c_n = a^2 \left(\frac{\Phi}{\Phi_c - \Phi}\right)^2,$$

$$c_i = a^2 \left(\frac{\Phi}{\Phi_c - \Phi}\right)^2 \left[\mu_c + \frac{\delta \mu b}{1 + b/\Phi_c - \Phi}\right].$$
forces at contacts oriented along $\theta_{\epsilon}$ (agreement with our data, and both viscosities diverge as Figure 4 shows that these expressions are in excellent symbols and their colors represent the same groups of simulations as (b) as a function of packing fraction from the simulation data. The Figure 2. The fits are the analytical expressions (6) and (7).

Figure 4 shows that these expressions are in excellent agreement with our data, and both viscosities diverge as $(\Phi_{c} - \Phi)^{-2}$ for $\Phi \rightarrow \Phi_{c}$, consistent with results of Boyer et al. and Trullsson et al. [4,6]. Note that in the viscous limit where $\text{St} \rightarrow 0$ ($I_m \rightarrow (\alpha_c I_m)^{1/2}$), we have $c_n \rightarrow \eta_n/(\alpha_c \eta_f)$ and $c_t \rightarrow \eta_t/(\alpha_c \eta_f)$. Hence, $c_n$ and $c_t$ can be viewed as a generalization of $\eta_n$ and $\eta_t$ to the visco-inertial regime. Equivalently, $1/c_n = I_m^2$ represents the generalized fluidity of the suspension [13].

V. EFFECT OF FLUID ON GRANULAR TEXTURE

This remarkable scaling of macroscopic variables with $I_m$ reflects the joint effects of particle inertia and fluid viscous forces on the particles. However, it is crucial for a better understanding of the flow to see to what extent the underlying texture properties are controlled by $I_m$. In particular, we consider here the fabric and force anisotropies and their relationship with the visco-inertial number. The fabric anisotropy $a_{f}$ of the contact network describes the excess of the number of contacts oriented along the principal strain-rate direction $\theta_{t}$ and their lack along the perpendicular direction $\theta_{t} + \pi/2$. The normal force anisotropy $a_{n}$ reflects the larger value of the mean normal force at contacts oriented along the principal stress direction $\theta_{n}$ compared to those in the perpendicular direction $\theta_{n} + \pi/2$. The tangential force anisotropy $a_{t}$ quantifies the stronger mobilization of friction forces at contacts oriented along $\theta_{t} \pm \pi/4$ compared to those in the principal direction $\theta_{t}$; see Appendix B for more details on the definition of these variables and their connection with the stress tensor.

By assuming $\theta_{t} = \theta_{n}$, an additive partition of the stress tensor leads to [14–17]

$$\mu(I_m) \simeq \frac{1}{2}[a_{c}(I_m) + a_{n}(I_m) + a_{t}(I_m)].$$

This is an outstanding relationship as it reveals the origins of shear strength in granular flows, herein expressed as a function of $I_m$. Figure 5 shows $\mu$ as a function of $I_m$ both directly obtained from the simulation data and calculated using Eq. (8). We see that the latter holds for all values of $I_m$. The inset shows the three anisotropies as a function of $I_m$. The fabric anisotropy $a_{f}$ increases monotonically with $I_m$ whereas the tangential force anisotropy $a_{t}$ increases to a lesser extent and $a_{n}$ declines.

Since $\mu$ is the half-sum of the three anisotropies and the normal and tangential force anisotropies add up to a nearly constant value, the main contribution to the increase of $\mu$ arises from the strong increase of $a_{f}$, as observed also in dry granular flows [16]. Correlatively, the divergence of the shear and normal viscosities as $I_m \rightarrow 0$ is accompanied by a decrease of fabric anisotropy. This is consistent with the fact that the packing fraction increases as the static limit is approached, and the anisotropy tends to decrease as packing fraction increases [18]. More importantly, Fig. 5 shows that the visco-inertial number provides a unique control parameter not only for the macroscopic flow parameters but also for the anisotropy parameters.

The scaling of anisotropy parameters with the visco-inertial number is not an evident property. It reflects the fact that the fluid affects the behavior by allowing for larger contact network anisotropy and stronger force chains. Not all texture parameters are expected to follow this scaling. For example, the coordination number $Z$ does not follow this scaling, as shown in Fig. 6. This fact does not, however, affect the shear strength $\mu$ since the coordination number $Z$ does not enter the expression (8). We may attribute this absence of scaling for $Z$ by the effect of lubricification forces between near-neighboring
particles. These forces depend on the shear rate and viscosity but not on the confining pressure. For the shear stress, the important parameter is not the number of contacts but rather how their orientations are distributed in different directions.

VI. CONCLUSIONS

In summary, our extensive simulations of non-Brownian particles in a sheared suspending fluid provide strong evidence that the parameter space can be reduced to a single visco-inertial number, which can also be interpreted as a generalized fluidity parameter of the suspension tending to the inverse of effective viscosity in the viscous limit. The robustness of the visco-inertial number was shown for frictional and inertial stresses. The rheology is mainly governed by the contact network anisotropy, reflecting the joint effects of viscous and inertial stresses. Our results obtained by means of subparticle computational fluid dynamics simulations considerably extend the scope of investigations to 3D and turbulent flows. Despite computational effort, the extension of this investigation to 3D and turbulent flows is desirable.

ACKNOWLEDGMENT

The authors would like to thank the IRSN (French authority of nuclear safety) for financial support.

APPENDIX A: NUMERICAL METHOD

In LBM, the fluid is described by the time-dependent distribution function $f(r,v,t)$ of particle positions $r$ and velocities $v$. The spatiotemporal evolution of $f$ is governed by the Boltzmann equation:

$$\left( \frac{\partial}{\partial t} + \mathbf{v} \cdot \nabla \right) f(r,v,t) = \Omega_{\text{coll}},$$

(A1)

where $m$ is the particle mass, $\mathbf{F}(r)$ summarizes external forces, and $\Omega_{\text{coll}}$ represents the collision operator describing the dynamics of collisions between fluid particles. The simplest collision model is the BGK (Bhatnagar-Gross-Krook) operator [19]:

$$\Omega_{\text{BGK}} = -\frac{1}{\tau} (f - f^0)$$

(A2)

where $\tau$ is the relaxation time and $f^0$ is the Maxwell-Boltzmann distribution. It is shown that the Boltzmann equation with the BGK collision operator yields to the Navier-Stokes equations [20]. The central idea of LBM [21–23] is to discretize the velocity vector space in a finite number of directions. We used the so-called D2Q9 (two dimensions with nine velocity directions), as shown in Fig. 7. A distinct distribution function $f_i$ is associated to each velocity direction $\mathbf{e}_i$.

The discretized equations along different directions are solved in two steps:

Collision : $f_i^{\text{out}}(r,t) = f_i(r,t) + \Omega_i$,  
Stream : $f_i(r + \Delta t \mathbf{e}_i,t + \Delta t) = f_i^{\text{out}}(r,t)$,  

(A3)

where $\Delta t$ is the time step and the $f_i^{\text{out}}(r,t)$ are the distribution functions after the collision step. The density $\rho(r,t)$ and fluid velocity $\mathbf{u}(r,t)$ are obtained as follows:

$$\rho(r,t) = \sum_i f_i(r,t),$$

$$\rho(r,t) \mathbf{u}(r,t) = \sum_i f_i(r,t) \mathbf{e}_i.$$  

(A4)

The BGK operator is simple but leads to fluctuating velocity fields. In our simulations, we used instead a multirelaxation time (MRT) collision approach [24,25]. It consists in associating nine moments to every fluid node, corresponding to the nine distribution functions, through a matrix $\mathbf{M}$ such that

$$\mathbf{m} = \mathbf{M f}$$  

(A5)
where \( \mathbf{m} = (m_0, m_1, \ldots, m_8)^T \) is the moment vector, \( \mathbf{f} = (f_0, f_1, \ldots, f_8)^T \) and \( \mathbf{M} \) is given by

\[
\mathbf{M} = \begin{pmatrix}
1 & 1 & 1 & 1 & 1 & 1 & 1 & 1 & 1 \\
-4 & -1 & 2 & -1 & 2 & -1 & 2 & -1 & 2 \\
4 & -2 & -1 & 1 & -2 & 1 & -2 & 1 & 1 \\
0 & 1 & 1 & 0 & -1 & -1 & 0 & 1 & 0 \\
0 & -2 & 1 & 0 & -1 & 2 & -1 & 0 & 1 \\
0 & 0 & 1 & 1 & 1 & 0 & -1 & -1 & -1 \\
0 & 0 & 1 & -2 & 1 & 0 & -1 & 2 & -1 \\
0 & 1 & 0 & -1 & 0 & 1 & 0 & -1 & 0 \\
0 & 0 & 1 & 0 & -1 & 0 & 1 & 0 & 1 - 1
\end{pmatrix}.
\]

Hence, the collision step is applied in the moment space, each moment \( m_i \) being relaxed to its equilibrium state \( m_i^{eq} \) with a relaxation time \( s_i \). The moments corresponding to the density \( \rho(\mathbf{r}, t) \) and the flux \( \mathbf{j}(\mathbf{r}, t) = \mathbf{f}(\mathbf{r}, t) \mathbf{u}(\mathbf{r}, t) \) are conserved. The moment vector \( \mathbf{m}^{out} \) resulting from collision can be written as follows:

\[
\mathbf{m}^{out} = \mathbf{S}(\mathbf{m} - \mathbf{m}^{eq}),
\]

where \( \mathbf{S} = \text{diag}(0, s_1, s_2, 0, s_4, 0, s_6, s_7, s_8) \) is a diagonal \( 9 \times 9 \) matrix. All relaxation times are proportional to \( \tau^{-1} \) [26]. The equilibrium moment vector \( \mathbf{m}^{eq} \) is given by

\[
\mathbf{m}^{eq} = \begin{pmatrix}
\rho \\
-2\rho + 3(j_x^2 + j_y^2)/\rho \\
-3(j_x^2 + j_y^2)/\rho \\
j_x \\
-j_x \\
j_y \\
-j_y \\
(j_x^2 - j_y^2)/\rho \\
j_x j_y/\rho
\end{pmatrix}.
\]

The distribution functions \( f_i^{out}(\mathbf{r}, t) \) resulting from the collision step are given by \( \mathbf{f}^{out} = \mathbf{M}^{-1} \mathbf{m}^{out} \). Finally, the streaming step is applied in the velocity space.

The no-slip boundary conditions are implemented through the bounce-back rule which consists of reflecting back the original fluid nodes in the opposite direction. The bounce-back rule, which consists of reflecting back the fluid nodes, is applied in the velocity space.

The no-slip boundary conditions are implemented through the bounce-back rule which consists of reflecting back the fluid nodes, is applied in the velocity space.

The no-slip boundary conditions are implemented through the bounce-back rule which consists of reflecting back the fluid nodes in the opposite direction. The bounce-back rule, which consists of reflecting back the fluid nodes, is applied in the velocity space.

The no-slip boundary conditions are implemented through the bounce-back rule which consists of reflecting back the fluid nodes, is applied in the velocity space.

The no-slip boundary conditions are implemented through the bounce-back rule which consists of reflecting back the fluid nodes in the opposite direction. The bounce-back rule, which consists of reflecting back the fluid nodes, is applied in the velocity space.

The no-slip boundary conditions are implemented through the bounce-back rule which consists of reflecting back the fluid nodes, is applied in the velocity space.

The no-slip boundary conditions are implemented through the bounce-back rule which consists of reflecting back the fluid nodes in the opposite direction. The bounce-back rule, which consists of reflecting back the fluid nodes, is applied in the velocity space.

The no-slip boundary conditions are implemented through the bounce-back rule which consists of reflecting back the fluid nodes, is applied in the velocity space.

The no-slip boundary conditions are implemented through the bounce-back rule which consists of reflecting back the fluid nodes in the opposite direction. The bounce-back rule, which consists of reflecting back the fluid nodes, is applied in the velocity space.

The no-slip boundary conditions are implemented through the bounce-back rule which consists of reflecting back the fluid nodes, is applied in the velocity space.

The no-slip boundary conditions are implemented through the bounce-back rule which consists of reflecting back the fluid nodes in the opposite direction. The bounce-back rule, which consists of reflecting back the fluid nodes, is applied in the velocity space.

The no-slip boundary conditions are implemented through the bounce-back rule which consists of reflecting back the fluid nodes, is applied in the velocity space.

The no-slip boundary conditions are implemented through the bounce-back rule which consists of reflecting back the fluid nodes in the opposite direction. The bounce-back rule, which consists of reflecting back the fluid nodes, is applied in the velocity space.

The no-slip boundary conditions are implemented through the bounce-back rule which consists of reflecting back the fluid nodes, is applied in the velocity space.

The no-slip boundary conditions are implemented through the bounce-back rule which consists of reflecting back the fluid nodes in the opposite direction. The bounce-back rule, which consists of reflecting back the fluid nodes, is applied in the velocity space.

The no-slip boundary conditions are implemented through the bounce-back rule which consists of reflecting back the fluid nodes, is applied in the velocity space.

The no-slip boundary conditions are implemented through the bounce-back rule which consists of reflecting back the fluid nodes in the opposite direction. The bounce-back rule, which consists of reflecting back the fluid nodes, is applied in the velocity space.

The no-slip boundary conditions are implemented through the bounce-back rule which consists of reflecting back the fluid nodes, is applied in the velocity space.