

Apparent slip mechanism between two spheres based on solvent rheology: Theory and implication for the shear thinning of non-Brownian suspensions

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Analytical results for the apparent slip between two spheres in a simple biviscous model of a shear-thinning fluid are presented. Velocity profiles and apparent slip lengths along the surfaces are analyzed in order to characterize the physical mechanism. It is shown that in this non-Newtonian model, the effect of shear thinning limited to high-shear rates in the interstitial regions between close spheres can be alternatively interpreted as the onset of an apparent shear-rate-dependent slippage effect. The results of the theory compare well with experiments from the literature showing the presence of surface slip on a particle approaching a planar wall. In terms of implications for suspensions rheology, the present results bridge the “hidden” solvent shear-thinning theory [Vázquez-Quesada *et al.*, *Phys. Rev. Lett.* **117**, 108001 (2016)] with slip-based models presented recently [Kroupa *et al.*, *Phys. Chem. Chem. Phys.* **19**, 5979 (2017)] as a possible explanation for the mechanism behind the shear thinning in hard-sphere non-Brownian suspensions.

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I. INTRODUCTION

The shear-thinning behavior observed in a concentrated non-Brownian suspension of spheres represents a long-standing issue in rheology [1–4]. In fact, in the absence of Brownian and inertial effects, and if a simple Newtonian solvent is considered, the resulting bulk suspension viscosity should be constant [5], which contradicts experimental results.

Three main hypotheses have been advanced in order to explain the observed shear thinning of non-Brownian suspensions (STNBSs): an effective size of the particles that is reduced with

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increasing shear, shear thinning of the solvent itself at the high shear rates occurring at interstitial regions, and slip at the surface spheres.

Advocating the first hypothesis, some authors have recently associated the STNBSs to the fact that the system behaves essentially as a suspension of soft particles at low shear stresses, with an apparent size that includes both the hard sphere and a part of the surrounding soft repulsive potential [6]. In simulations performed by Mari *et al.* [7], for example, the value of the repulsive length in the interparticle potential is chosen on the order of 5% of the particle radius a (i.e., the so-called Debye length $\kappa^{-1} = 0.05a$) corresponding to typical nonzero force contributions on a spatial range of order of $0.1 \mu\text{m}$ for particles a few μm in size. As discussed recently by Chatte *et al.* [6], under these conditions the “apparent” size of the particles (including hard-core and soft repulsive range) will decrease with increasing shear rate, leading to a suspension viscosity which decreases as a function of the shear rate, i.e., shear thinning.

Previous results [8] had also shown that excluded volume effects can induce a shear-thinning contribution coming from interparticle repulsive forces; however, this was significant only for soft particles. In the hard-sphere limit, the interparticle contribution to the suspension viscosity was found irrelevant. This was also confirmed by independent non-Brownian simulations using nearly hard-core potentials (i.e., small effective repulsive ranges $\kappa^{-1} = 0.001a$) where Newtonian or opposite mild shear-thickening behavior was always reported [9–11]. It is worth noticing that in most of the noncolloidal systems studied in the literature where shear thinning has been observed [1–4], the particle size exceeds typically $40 \mu\text{m}$, and shear thinning was observed for surface roughness down to $30\text{--}90 \text{ nm}$ [12], i.e., in the hard-sphere regime discussed above. Therefore, for hard-sphere noncolloidal systems shear thinning is unlikely to be due to an excluded volume effect.

The second hypothesis, advanced in a recent work [13], suggests that in the interstitial regions between spherical particles very high shear rates can occur, which can trigger a local shear thinning of a “nominally” Newtonian matrix; this effect, in turn, can lead to shear thinning at the average (macroscopic) shear rate. Since the matrix shear thinning is negligible at the mean shear rate and the local thinning occurs well beyond, the matrix was considered as “nominally” Newtonian and the mechanism was termed “hidden.” Good agreement between experiments and simulations for suspensions of noncolloidal particles in several silicon oils was obtained. It should be noticed that in the case of simpler liquids (e.g., glycerol), although shear thinning does take place at extremely large critical bulk shear rates, hydrodynamic conditions under strong confinement can alter the liquid property itself. In the work by Jabbarzadeh *et al.* [14], for example, molecular simulations have shown that extreme confinement (on the order of tens of molecular sizes), similar to that encountered in the squeezing flow between nearly touching surfaces, can anticipate significantly the onset of matrix shear thinning with respect to the bulk predictions, and therefore this effect might play a role even for liquids which are more Newtonian than low-molecular-weight oils. This was also shown in experiments with other simple lubricating fluids, such as octamethylcyclotetrasiloxane (OMCTS), where a massive increase in the viscosity was observed as a result of liquid-solid [15–17] or continuous glasslike transitions [18,19]. In all these cases, the effective liquid shear thinning is the result of an increased viscosity with respect to the bulk value linked to an alteration of its microstructure in response to strong confinement and generally disappear for liquid film thickness beyond 10 molecular dimensions [20,21].

Finally, very recently a new mechanism has been proposed for noncolloidal shear thinning based on the connection between the slip on the surface of the particles and the suspension viscosity [22]. Results based on a phenomenological slip model showed that surface slip might be crucial for the rheology and flow in crowded systems and can be an alternative explanation for STNBSs. Traditionally, when a deviation of the no-slip behavior is detected, a constant slip length is generally used to characterize the slip on the experiment. Such a length is calculated *a posteriori* from the results of the experiment, so it is not used to do predictions but rather to calibrate the system behavior [23]. In Ref. [22] a phenomenological relation between slip length and local shear rate in the gap between two particles was assumed.

In fact, slip between fluids and solids has been observed in many experiments [24] and reported, for example, in concentrated polymer suspensions and polymer melts [25,26]. Different explanations and theories have been proposed related to the microstructural changes of the polymer configuration of the fluid close to the solid interface [27]. Shear-dependent wall slip has been observed even in Newtonian (low-molecular-weight) liquids [20,21,28], and it can be associated to several factors, including surface properties (wetting, rugosity, dissolved gas layer, cavitation) [29–33], as well as bulk liquid effects including apparent shear thinning at large shear rates or viscous heating [34,35]. It should be remarked that the effective shear thinning of the liquid related to a slippage effect is intrinsically different from the confinement-induced effect discussed above. In fact, in that case the effect disappears for wider lubricating films, whereas effective shear thinning (e.g., reduced lubrication force) is still observed in experiments [20,21].

In this work, we take the point of view that an apparent slip at the surface of the particles may be due to shear-banding effects, i.e., a decrease in viscosity in a very thin near-to-wall layer due to either a change in the structure of liquid, enhanced nucleation of a dissolved gas [23,28,31], or a shear thinning of the solvent itself [36]. We reformulate some of the conclusions made in the “hidden” shear-thinning approach of Ref. [13] in terms of an apparent slip length occurring in the interstitial squeezing flow between particles and show that this interpretation is consistent with the arguments made by Kroupa *et al.* [22].

Specifically, we present a study of the velocity profiles of the squeezing flow of a shear thinning fluid between two spheres in the lubrication regime, when the gap between the spheres is much smaller than their radii. To allow for an analytical treatment, the solvent is modeled as a simple biviscous fluid. The model developed in Ref. [37] allows us to calculate the resulting “apparent” slip length, which is well defined, exhibits a natural dependence on the local shear rate, and has a one-to-one correspondence to the underlying liquid rheology.

The results show that “hidden” shear-thinning models for nominally Newtonian matrices such as those presented in Ref. [13] can be alternatively interpreted in terms of an “apparent” surface slip. This leads to an equally possible explanation of the shear-thinning mechanism on noncolloidal systems, in agreement with the one put forward recently by Kroupa *et al.* [22].

II. BIVISCOUS MODEL: SQUEEZING FLOW

In this section, we present a summary of the relevant analytical results for the velocity profile, pressure gradient, and force in the squeezing flow between two very close spheres embedded in a biviscous fluid. Further details can be found in Ref. [37]. The fluid is assumed to obey the following constitutive equation for the shear viscosity:

$$\eta = \begin{cases} \eta_0, & \text{if } |\dot{\gamma}| < \dot{\gamma}_c \\ \eta_1, & \text{if } |\dot{\gamma}| \geq \dot{\gamma}_c \end{cases} \quad \text{with } \eta_0 > \eta_1. \quad (1)$$

That is, for shear rates below the critical shear rate $\dot{\gamma}_c$ the fluid has a large viscosity η_0 , while it has a small viscosity η_1 for larger shear rates. This is a very schematic rheological model that nevertheless captures the basic phenomenology of a shear-thinning fluid. Because of the simplicity of the model, analytical solutions can be found.

We consider the flow problem in which two spheres of radii a_1 and a_2 separated by a minimum gap h_0 approach each other with a normal relative velocity V . The flow problem in which the two spheres move with parallel and opposite velocities thus generating a shear flow is not considered here because it gives much smaller velocity profiles. The system of reference and the different parameters of the problem are shown in Fig. 1. When the velocity V is small enough and the gap is small, we may assume the validity of the lubrication regime, where the fluid squeezes unidirectionally along the r coordinate. Under these conditions, the Stokes equations *with no-slip boundary conditions* plus continuity of stress along the gap can be solved. In the case of a biviscous fluid, the model can display a position r -dependent shear banding; that is, depending on critical conditions the fluid displays two regions denoted as $\mathcal{R}_m$, $\mathcal{R}_b$ along the horizontal direction. The flow

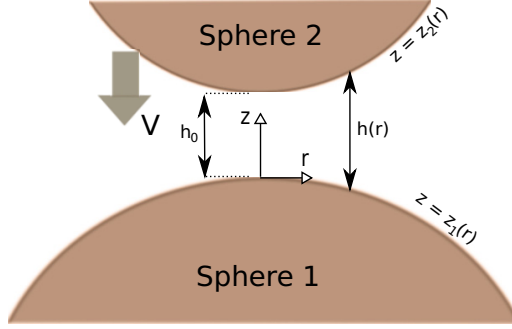


FIG. 1. Scheme of two approaching spheres of radii a_1, a_2 in the lubrication regime ($h_0 \ll a_1, a_2$).

field has no discontinuities in the shear rate in $\mathcal{R}_m$ (monoviscouslike region), whereas in region $\mathcal{R}_b$ (biviscouslike region) the fluid does shear band. As discussed in Ref. [37], these regions $\mathcal{R}_m$ and $\mathcal{R}_b$ are defined as

$$r \in \mathcal{R}_m \text{ if } r \leq r_1 \text{ or } r \geq r_2, \quad r \in \mathcal{R}_b \text{ if } r_1 < r < r_2, \quad (2)$$

where r_1 and r_2 are given by

$$r_1 = Q - \sqrt{-3 - Q^2 + 4\left(\frac{\dot{\gamma}_s^{\max}}{\dot{\gamma}_c}\right)\frac{1}{Q}}, \quad r_2 = Q + \sqrt{-3 - Q^2 + 4\left(\frac{\dot{\gamma}_s^{\max}}{\dot{\gamma}_c}\right)\frac{1}{Q}} \quad (3)$$

with

$$Q = \sqrt{-1 + \frac{1}{\mathcal{P}} + \mathcal{P}}, \quad \mathcal{P} = \sqrt[3]{-1 + 2\left(\frac{\dot{\gamma}_s^{\max}}{\dot{\gamma}_c}\right)^2 + 2\left(\frac{\dot{\gamma}_s^{\max}}{\dot{\gamma}_c}\right)\sqrt{\left(\frac{\dot{\gamma}_s^{\max}}{\dot{\gamma}_c}\right)^2 - 1}}, \quad (4)$$

and

$$\dot{\gamma}_s^{\max} = \frac{9}{8} \frac{V}{h_0} \sqrt{\frac{3}{2} \frac{a}{h_0}}, \quad (5)$$

where $a = (1/a_1 + 1/a_2)^{-1}$ is the reduced radius. The fluid shear bands, $\mathcal{R}_b \neq \emptyset$, when the following condition for the interparticle separation distance is satisfied:

$$h_0 < h_0^{\lim}, \quad (6)$$

where

$$h_0^{\lim} \equiv \frac{3}{4} \sqrt[3]{\frac{9V^2a}{2\dot{\gamma}_c^2}}. \quad (7)$$

This condition is equivalent to

$$\dot{\gamma}_c < \dot{\gamma}_s^{\max}. \quad (8)$$

In other words, given a bead radius a and a fluid rheology $\dot{\gamma}_c$, there is a critical distance $h_0^{\lim}$, determined by the approaching velocity V such that, for smaller separations h_0 , the flow shows two distinct regions along the r direction, a shear-banded one $\mathcal{R}_b$ with discontinuities separating two fluids with different viscosities, and another one $\mathcal{R}_m$ without velocity-gradient discontinuities.

The corresponding flow field is given by the following piecewise two-dimensional velocity profile:

$$u(r, z) = \begin{cases} u(r, z)|_m, & \text{if } r \in \mathcal{R}_m \\ u(r, z)|_b, & \text{if } r \in \mathcal{R}_b \end{cases}. \quad (9)$$

The functional form of $u(r, z)|_m$ and $u(r, z)|_b$ is

$$\begin{aligned} u(r, z)|_m &= \frac{1}{2\eta_0} \frac{\partial p}{\partial r}(r) [z - z_1(r)][z - z_2(r)], \\ u(r, z)|_b &= \begin{cases} \frac{1}{2\eta_0} \frac{\partial p}{\partial r}(r) \{ [z - c(r)]^2 - z_c^2(r) \} + \frac{1}{2\eta_1} \frac{\partial p}{\partial r}(r) [z_c^2(r) - h^2(r)/4], & \text{if } |z - c(r)| < z_c(r) \\ \frac{1}{2\eta_1} \frac{\partial p}{\partial r}(r) [z - z_1(r)][z - z_2(r)], & \text{if } |z - c(r)| \geq z_c(r) \end{cases}. \end{aligned} \quad (10)$$

Here $c(r) = [z_1(r) + z_2(r)]/2$ is the loci of points equidistant from both surfaces. At a given value of $r \in \mathcal{R}_b$, the velocity profile has two different regions: in $|z - c(r)| \leq z_c(r)$ the effective viscosity is η_0 , while in $|z - c(r)| > z_c(r)$ it is η_1 . The distance $z_c(r)$ is given, when $r \in \mathcal{R}_b$, as

$$\left[\frac{2z_c(r)}{h(r)} \right] = S_1 + S_2, \quad (11)$$

$$\begin{aligned} S_1 &= \sqrt[3]{\frac{\eta_0}{2(\eta_0 - \eta_1)} \left\{ 1 + \sqrt{1 + \frac{4}{27} \frac{\eta_1^2}{\eta_0^2} \frac{\eta_1}{\eta_0 - \eta_1} \left[\frac{\dot{\gamma}_s(r)}{\dot{\gamma}_c} \right]^3} \right\}}, \\ S_2 &= \sqrt[3]{\frac{\eta_0}{2(\eta_0 - \eta_1)} \left\{ 1 - \sqrt{1 + \frac{4}{27} \frac{\eta_1^2}{\eta_0^2} \frac{\eta_1}{\eta_0 - \eta_1} \left[\frac{\dot{\gamma}_s(r)}{\dot{\gamma}_c} \right]^3} \right\}}, \end{aligned} \quad (12)$$

where $h(r)$ is the gap size at the coordinate r and $\dot{\gamma}_s(r) = 3rV/h^2(r)$ is the maximum shear rate that a Newtonian fluid would have at the given position r . This maximum would occur at the surface of the spheres.

The value $\dot{\gamma}_s^{\max}$ in Eq. (5) is the maximum of the function $\dot{\gamma}_s(r)$. Note that $z_c(r)$ is a distance which cannot be greater than $h(r)/2$ when $r \in \mathcal{R}_b$. When the gap between the two spheres is much smaller than the radii of the spheres, the distance between the two surfaces is given approximately by $h(r) = h_0(1 + \frac{r^2}{2ah_0})$.

The pressure gradient for $r \in \mathcal{R}_b$ is

$$\left. \frac{\partial p}{\partial r}(r) \right|_b = -\frac{\dot{\gamma}_c \eta_0}{z_c(r)}. \quad (13)$$

In the region $\mathcal{R}_m$ the pressure gradient is given by the usual expression for a Newtonian fluid:

$$\left. \frac{\partial p}{\partial r}(r) \right|_m = -\frac{6\eta_0 r V}{h^3(r)}. \quad (14)$$

Finally, an approximated analytical expression of the force between both spheres is given by [37]

$$F = \begin{cases} F_1 & \text{if } h_0 \geq h_0^{\lim} \\ F_2 & \text{if } h_0 < h_0^{\lim} \end{cases}. \quad (15)$$

where

$$\begin{aligned}
 F_1 &= 6\pi\eta_0 V a^2 h_0^{-1}, \\
 F_2 &= 2\pi \left(3a^2\eta_0 V h_0^{-1} + 3a^2(\eta_0 - \eta_1) V [h^{-1}(r_2) - h^{-1}(r_1)] \{2 - h_0[h^{-1}(r_1) + h^{-1}(r_2)]\} \right. \\
 &\quad \left. + 2a\dot{\gamma}_c(\eta_0 - \eta_1) \left\{ (r_2 - r_1) + [h(r_1) - 2h_0] \sqrt{\frac{2a}{h_0}} \left[\arctan\left(\frac{r_2}{\sqrt{2ah_0}}\right) - \arctan\left(\frac{r_1}{\sqrt{2ah_0}}\right) \right] \right\} \right).
 \end{aligned} \tag{16}$$

III. SLIP MECHANISM

A. Conditions for slip between spheres

Although from a theoretical perspective slip and shear banding are different phenomena, from an experimental point of view, a shear-banding fluid may produce a behavior that is indistinguishable from slip of a usual Newtonian fluid. In this section, we reinterpret the results of a shear-banding biviscous fluid in terms of an apparent slip of a Newtonian fluid with a viscosity given by η_0 .

Note that the expressions for r_1 and r_2 in Eq. (3) do not depend on the viscosity ratio η_1/η_0 . This implies that the region $\mathcal{R}_b \neq \emptyset$ as long as the condition (8) holds. In this case, a boundary layer of viscosity η_1 on the surfaces of the spheres between the coordinates r_1 and r_2 is always present. From Eq. (11) the following limit can be calculated:

$$\lim_{\eta_1/\eta_0 \rightarrow 0} z_c(r) = \frac{h(r)}{2}, \tag{17}$$

indicating that such a layer becomes infinitesimal in such a limit. Therefore, in the limit $\eta_1/\eta_0 \rightarrow 0$, if (8) is satisfied, there will be a real slip in the region $\mathcal{R}_b$.

The other parameter controlling $z_c(r)$ is $\dot{\gamma}_c$; in the limit $\dot{\gamma}_c \rightarrow 0$ one obtains that $r_1 \rightarrow 0$, $r_2 \rightarrow \infty$, and $h_0^{\lim} \rightarrow \infty$, $z_c(r) \rightarrow 0$ and the behavior of a solvent of viscosity η_1 is recovered. Therefore, for small $\dot{\gamma}_c$, the layer of the fluid of viscosity η_1 close to the spheres cannot be considered as a thin layer, given that the thickness of such a layer is comparable to the gap. In conclusion, an apparent slip can be observed if (1) the condition (8) is satisfied, (2) the critical shear rate $\dot{\gamma}_c$ is high enough, i.e., $\dot{\gamma}_c/\dot{\gamma}_s(r_{\max})$ is not too small, and (3) the viscosity ratio η_1/η_0 is low enough. The last two conditions can be reduced to $z_c(r) \sim h(r)/2$ at some $r \in [r_1, r_2]$, which is the condition of having a thin layer.

Of course, the proposed conditions are only necessary to detect slip if the velocity profiles are observed directly. If the slip is inferred from indirect methods, any deviation from the Newtonian behavior could be interpreted as slip, making the first condition as the only necessary one.

B. Calculation of the slip length

Under the above conditions, a slip behavior will be present between both spheres in the region $\mathcal{R}_b$; this is in $r_1 < r < r_2$. The apparent slip length λ is defined as the distance beyond the interface at which the liquid velocity extrapolates to zero, and it can be calculated in that region from the analytical velocity profile (10).

If we consider that the region $|z - c(r)| \geq z_c(r)$ is very narrow, the apparent velocity u_{ap} in the surface of the spheres will be the extension of the velocity profile in the region $|z - c(r)| < z_c(r)$ to the coordinates $z = z_1(r)$ or $z = z_2(r)$. The expression is

$$u_{ap} = \frac{1}{2} \frac{\partial p}{\partial r}(r) \Big|_b \left(\frac{1}{\eta_1} - \frac{1}{\eta_0} \right) [z_c^2(r) - h^2(r)/4]. \tag{18}$$

Within the lubrication regime, given that $a \gg r$, curvature effects do not enter into the slip length, which can be calculated as

$$\lambda(r) = \frac{u_{ap}(r)}{\dot{\gamma}^{ap}(r)}, \quad (19)$$

where the apparent shear rate $\dot{\gamma}^{ap}(r)$ on the coordinates $z = z_1(r)$ or $z = z_2(r)$ is given by

$$\dot{\gamma}^{ap}(r) = -\frac{1}{\eta_0} \frac{\partial p}{\partial r}(r) \bigg|_b \frac{h(r)}{2} = \dot{\gamma}_c \left[\frac{h(r)}{2z_c(r)} \right]. \quad (20)$$

The final result is

$$\lambda(r) = \frac{h(r)}{4} \left(\frac{\eta_0}{\eta_1} - 1 \right) \left\{ 1 - \left[\frac{\dot{\gamma}_c}{\dot{\gamma}^{ap}(r)} \right]^2 \right\}, \quad (21)$$

where the relation of the slip length with $\dot{\gamma}^{ap}(r)$ has been explicitly written. Note that the slip length in the present model depends on the coordinate r , at variance with usual approximations in which it is taken as a constant [28,38,39].

Note also that when the limit $\eta_1 = \eta_0$ is considered, although the length λ cannot be considered properly as a slip length because the condition $\eta_1/\eta_0 \ll 1$ is not held, the no-slip condition is recovered because $\lim_{\eta_1 \rightarrow \eta_0} \lambda(r) = 0$. In the opposite limit, when $\eta_1/\eta_0 \rightarrow 0$, the analytical expression for the slip length tends to

$$\lim_{\eta_1/\eta_0 \rightarrow 0} \lambda(r) = \frac{h(r)}{6} \left[\frac{\dot{\gamma}_s(r)}{\dot{\gamma}_c} - 1 \right]; \quad (22)$$

i.e., it is proportional to the wall shear rate leading to higher apparent fluid slippage on regions of higher local wall shear rates. Note that Eqs. (21) and (22) are valid only in the biviscous regions. In the monoviscous regions, i.e., for $\dot{\gamma}_s(r) \leq \dot{\gamma}_c$, there is not apparent slip, and the slip length is 0.

C. Slip length: Model parameters sensitivity

In order to illustrate the above theoretical results in more practical terms, we consider the typical case of two noncolloidal spherical particles of radius $R = 50 \mu\text{m}$ approaching each other along the line joining their centers at a relative velocity $V = a\dot{\gamma}$. We investigate the effect of changing the ratio η_1/η_0 within the range 0.01–0.9 while keeping the following parameters constant: $\dot{\gamma} = 1 \text{ s}^{-1}$, $\dot{\gamma}_{\text{cut}} = 1000 \text{ s}^{-1}$, and $h_0 = 4 \times 10^{-3} R$.

In Fig. 2 (top) we plot the velocity profiles at the coordinate $r = (r_1 + r_2)/2$. As already mentioned, r_1 and r_2 do not depend on the ratio η_1/η_0 , so they remain the same for all velocity profiles. When the ratio η_1/η_0 is decreased the profile tends to a slip behavior in the region $\mathcal{R}_b$, determined by the limit (22) of the slip length. Even at the largest ratios one can calculate an “apparent” slip length for all the profiles via Eq. (21). Such slip lengths are shown, for the whole region $\mathcal{R}_b$, in Fig. 2 (bottom). Note that, although r_1 can be very small, as in the figure, it is not zero, so in this model there will be always a region around $r = 0$ free of apparent slip.

The same system is considered now at a fixed viscosity ratio $\eta_1/\eta_0 = 10^{-2}$ with a varying dimensionless gap between the spheres $h_0^* = h_0/h_0^{\text{lim}}$ within the range 0.01–0.9. Given that now the values of r_1 and r_2 are changing with h_0 , the velocity profiles are drawn at the position $r = 2 \mu\text{m} \approx 0.15 h_0^{\text{lim}}$, which is included in the region $\mathcal{R}_b$ for all the gaps considered. In Fig. 3 (top) the velocity profiles of the system at different gaps are shown, while in Fig. 3 (bottom) we show the corresponding slip lengths. The slip length in the limit $h_0 \rightarrow 0$ is also shown.

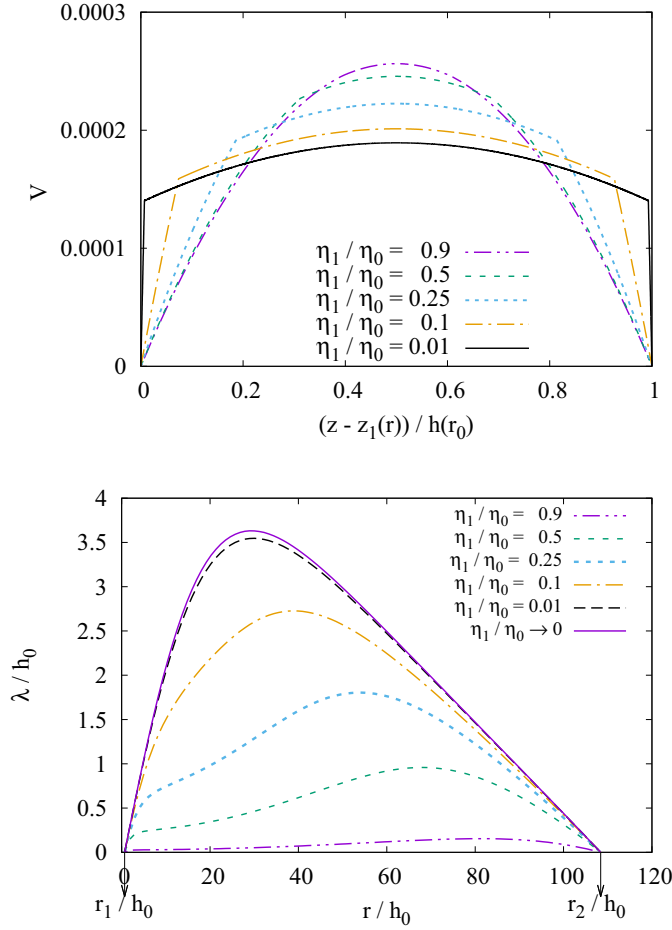


FIG. 2. Top: velocity profiles for different ratios η_1/η_0 at $r = (r_1 + r_2)/2$. Bottom: slip lengths of the same systems at region $r_1 < r < r_2$ calculated via Eq. (21). The limiting curve (22) for the $\eta_1/\eta_0 \rightarrow 0$ case has also been drawn.

IV. COMPARISON WITH EXPERIMENTS

A. Hydrodynamic force acting on a sphere approaching a flat surface

Direct observations of slip from velocity profiles of simple fluids are scarce, with most of them, to the best of our knowledge, being estimated in microchannels or between planar walls [24]. In this section, the analytical expression of the force between spheres in Eqs. (15) and (16) will be compared against the experimental results in Ref. [28], where slip has been inferred indirectly from drag on a sphere. The experiment consists of a sphere of radius $a_1 = 10.4 \mu\text{m}$ immersed in an aqueous sucrose solution and approaching a plane surface ($a_2 \rightarrow \infty$). The control parameter is the approaching velocity V , whereas the force on the sphere is measured as a function of the separation distance. For an approaching velocity $V = 2400 \text{ nm/s}$ the experiments show a good agreement with the analytical theory for a Newtonian (Stokes) fluid with no-slip boundary conditions enforced (Fig. 4, blue circles for experiments and blue line for theory). On the other hand, for $V = 21\,600 \text{ nm/s}$ deviations from the analytical result with no-slip conditions are observed (Fig. 4, red circles).

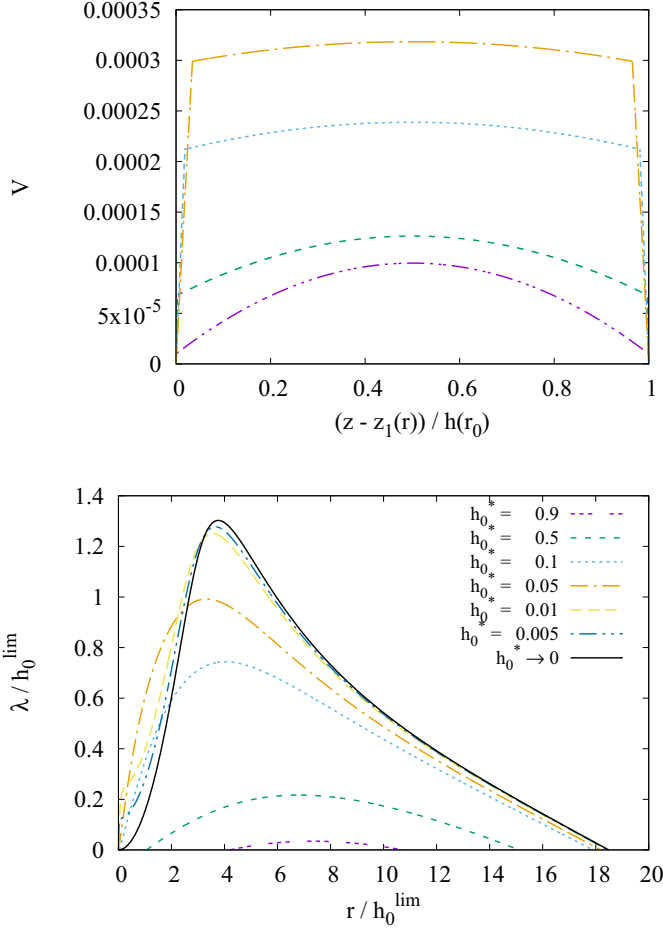


FIG. 3. Top: Velocity profiles for different values of $h_0^* = h_0 / h_0^{\lim}$ at $r = 2 \times 10^{-6}$ m $\approx 0.15 h_0^{\lim}$. The key is the same than the bottom one. Bottom: The corresponding slip lengths calculated via Eq. (21). The limit of such an equation when $h_0 \rightarrow 0$ has been also drawn.

Such a deviation can be explained by relaxing the Newtonian approximation and letting the fluid show shear thinning at very high shear rates only. The biviscous model offers a possible route for the explanation of experimental data. From Eq. (7) the critical shear rate can be calculated as

$$\dot{\gamma}_c = V \sqrt{\frac{9a}{2} \left(\frac{3}{4h_0^{\lim}} \right)^3}. \quad (23)$$

From the experimental data with $V = 2400$ nm/s, given that the fluid behaves as Newtonian at least until $1/h = 3 \times 10^7$ m $^{-1}$ one can deduce that $h_0^{\lim} < 3 \times 10^{-8}$ m leading to $\dot{\gamma}_c > 1800$ s $^{-1}$. On the other hand, from the experimental curve at $V = 21\,600$ nm/s, we consider that deviations from the Newtonian behavior occur around $1/h \sim 10^7$ m $^{-1}$, which would give a critical shear rate of $\dot{\gamma}_c \sim 3000$ s $^{-1}$. Note that this value for the critical shear rate is also close to that given in Ref. [22] for the onset of slippage effects in a Newtonian fluid. By assuming, for example, $\dot{\gamma}_c = 3500$ s $^{-1}$, the present biviscous model can adjust quite well experimental data by choosing $\eta_1/\eta_0 = 0.6$ (Fig. 4, red line). In conclusion, the present biviscous model exhibits a high-shear-rate shear-thinning behavior, which can be alternatively interpreted as an apparent shear-rate-dependent surface slip occurring between

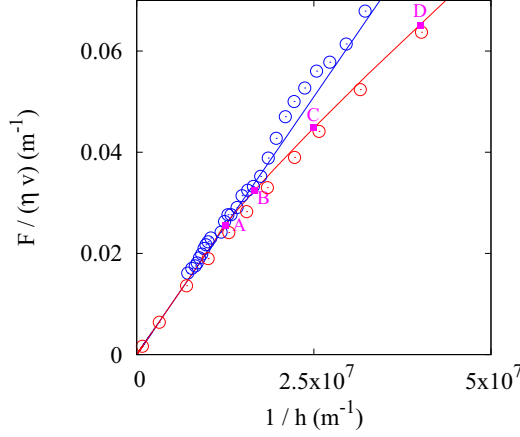


FIG. 4. Force on the approaching sphere as a function of the inverse of the separation between the sphere and the plane surface. The experimental results by Neto *et al.* [28] have been drawn with circles. For the case of an approaching velocity 2400 nm/s the fluid displays Newtonian behavior (blue circles) for the whole range of distances. For the case of an approaching velocity of 21 600 nm/s, the Newtonian behavior is observed only at large separations. As the sphere approaches the planar surface, deviations from the Newtonian behavior are observed. The continuum red line is obtained from the theoretical result (16) with the following fitting parameters: $\dot{\gamma}_c = 3500 \text{ s}^{-1}$ and $\eta_1/\eta_0 = 0.6$. The pink squares A–D correspond to the frames shown at Fig. 5.

two approaching spheres. The results are consistent with the experimental data provided by Neto *et al.* [28]. Shear-rate dependence of the slip length was also reported for other Newtonian liquids and was shown to correlate with surface wettability and roughness [20,21,40]. It should be noticed that in the case of smoothed hydrophobic surfaces small apparent slip lengths were observed with no shear-rate dependence up to 5000 s^{-1} [41].

In Fig. 5 the local topology of the interface separating fluid regions with viscosity η_1 (low-viscosity regions: light gray) from those where the viscosity is η_0 (high-viscosity regions: dark gray) is shown. Low-viscosity regions can be interpreted as those where apparent slip occurs. When

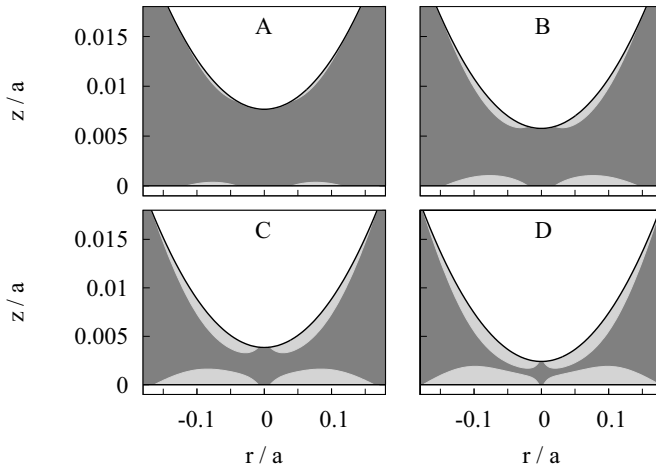


FIG. 5. Low-viscosity regions for the squeezing flow between a sphere of radius $a_1 = 10.4 \mu\text{m}$ and a flat surface for different gaps h_0 . Snapshots A–D correspond to the points in the force highlighted in Fig. 4.

the sphere gets closer to the surface, the low-viscosity regions become wider, spreading into the bulk in the limit $h_0 \rightarrow 0$. Snapshots A–D correspond to the points in the force highlighted in Fig. 4. Bigger deviations arise in correspondence with increased near-wall low-viscosity layers. Deviations from the no-slip Stokes result for approaching particles are expected to become significant for gaps $h_0 < h_0^{\text{lim}}$ at the given macroscopic shear rate.

B. Hydrodynamic force between two spheres: Biviscous versus slip models

In Ref. [22], a phenomenological model has been proposed for the slip length. Even though the model is phenomenological, we compare our results with the result of this model because the latter accounts for experimental results. In this way, we are comparing indirectly with experimental results [3,4]. The slip-length model of Ref. [22] is given by the following phenomenological expression:

$$\lambda(\dot{\gamma}(r)) = \lambda_{\text{max}} - \frac{\lambda_{\text{max}} - \lambda_{\text{min}}}{1 + \frac{\dot{\gamma}(r)}{\dot{\gamma}_c}}, \quad (24)$$

and they have performed an inverse-engineering study, i.e., they have adjusted the model parameters to match the STNBSs observed in experiments [3,4]. In the case of Ref. [3], they found that $\lambda_{\text{min}}/a = 0$, $\lambda_{\text{max}}/a = 0.04$, $\dot{\gamma}_c = 1000 \text{ s}^{-1}$ was reproducing accurately experimental data for the suspension viscosity. In this section we compute the corresponding lubrication force acting between two spheres assuming that a slip governed by Eq. (24) occurred on the particle surface. In order to do it, we start from the expression for the pressure gradient $\partial p(r)/\partial r$ presented in Ref. [22] [see their Eqs. (9)–(12)] and double integrate it numerically along the radial coordinate r to derive the force:

$$p(r) = - \int_r^\infty \frac{\partial p}{\partial r}(r) dr \quad (25)$$

and

$$F = \int_0^\infty p(r) 2\pi r dr. \quad (26)$$

In the solution, we consider an average bulk shear rate $\dot{\gamma} = 10 \text{ s}^{-1}$ and particle radii $a = 43 \text{ }\mu\text{m}$. Figure 6 shows the resulting interparticle lubrication force coming from the Kroupa analysis with experimentally derived model parameters and the lubrication force obtained analytically for the biviscous model, Eq. (16). In the biviscous model viscosity η_0 is taken to be that of the glycerol solution of Ref. [3]. In addition the critical shear rate $\dot{\gamma}_c = 1000 \text{ s}^{-1}$ is chosen to be the same as as in the slip model of Ref. [22]. The effect of the only left adjusting parameter in the biviscous model, i.e., the viscosity ratio η_1/η_0 , is also shown in the figure. It is clear that the biviscous model captures well the transition between the far-field behavior and the short-range force plateau at the given distance ($\approx 0.1a$). For $\eta_1/\eta_0 = 0.001$ the agreement with the Kroupa model (matching the experiment) is excellent.

The result shows that, in terms of interparticle lubrication forces, the behavior of a solvent matrix characterized by “hidden shear thinning” behavior at large shear rate [13] is consistent with that one of a perfectly Newtonian matrix with finite slippage at the particle surface. As a consequence, the origin of the STNBSs can be interpreted in both ways.

As a last remark, it should be noticed that the critical shear rate $\dot{\gamma}_c$ in the biviscous model needed to obtain quantitative agreement with Kroupa’s results (and therefore with the experimental data of Ref. [3]) is orders of magnitudes smaller than realistic values for which shear thinning of the solvent practically occurs; i.e., for glycerol matrices $\dot{\gamma}_c$ is expected to be in the order 10^4 – 10^5 s^{-1} or even larger [36]. The present low value of $\dot{\gamma}_c$ is consistent with experimentally adjusted wall-slip models and shows that an effective shear-thinning rheology (rather than a true bulk rheology) of the solvent in the narrow gaps between approaching particles can ultimately lead to the observed overall STNBSs.

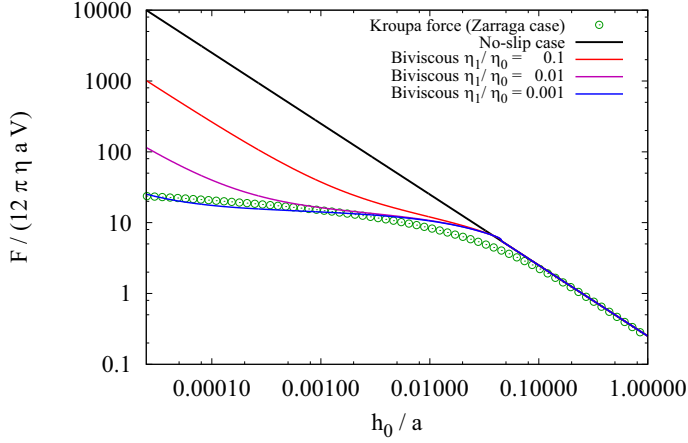


FIG. 6. Comparison of interparticle lubrication forces between two spheres of radius $a = 43 \mu\text{m}$ in a shear flow with average $\dot{\gamma} = 10 \text{ s}^{-1}$. Black line: Newtonian lubrication. Blue circles: force derived by Kroupa's slip model [22] with experimentally adjusted parameters ($\lambda_{\min}/a = 0$, $\lambda_{\max}/a = 0.04$, $\dot{\gamma}_c = 1000 \text{ s}^{-1}$). Red, purple, and green lines: biviscous model; $\dot{\gamma}_c = 1000 \text{ s}^{-1}$.

V. CONCLUSIONS

In this article the biviscous model has been used to determine the apparent slip between spheres caused by a “hidden” shear thinning of a nominally Newtonian fluid, i.e., occurring only at high shear rates or as a result of the interaction of the lubricating liquid with the external geometries. We propose to use a critical shear rate $\dot{\gamma}_c$ (for the onset of shear-thinning effects) and a high-to-low shear viscosity ratio η_1/η_0 as the relevant parameters to determine the “apparent” slip. Differently to current slip-length models, where slippage velocity is considered as a boundary condition and used to calibrate experimental results, here the slip length can be uniquely inferred from the rheological parameters of the solvent ($\dot{\gamma}_c$ and η_1/η_0) which could be estimated through independent rheological or AFM experiments. The biviscous model may describe two physical situations, either a *true* rheological property of the solvent (as in silicon oils [36]) or an *effective* viscosity decrease in simpler fluids, e.g., due to wall roughness or near-wall liquid inhomogeneities such as microstructural changes or presence of a dissolved phase [22,23,28]. Moreover, we demonstrate that in this biviscous model, the slip length is not constant along the surface but contains a natural shear-rate dependence in agreement with experimental findings [20,21,28] and previous suggested theories [32,33]. Experimental data for the hydrodynamic drag experienced by a sphere approaching a planar wall are excellently reproduced by the model. In terms of implications for suspension rheology, these results bridge therefore the “hidden” solvent shear-thinning theory [13] with the slip-based models presented recently in Ref. [22] as a possible explanation of the shear thinning in hard-sphere non-Brownian suspensions.

The present biviscous model allows also for a natural incorporation of noncontinuum nanorheology effects for lubricating fluids under very confined situations and can pave the way towards more realistic short-range hydrodynamic interaction models between particles in real suspensions.

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