Rheological properties of the soft-disk model of two-dimensional foams

Langlois, Vincent; Hutzler, Stefan; Weaire, Denis

Published in:
Physical Review E

Link to article, DOI:
10.1103/PhysRevE.78.021401

Publication date:
2008

Document Version
Publisher's PDF, also known as Version of record

Link back to DTU Orbit

Citation (APA):
Langlois, V., Hutzler, S., & Weaire, D. (2008). Rheological properties of the soft-disk model of two-dimensional foams. Physical Review E, 78(2), 021401. https://doi.org/10.1103/PhysRevE.78.021401
Rheological properties of the soft-disk model of two-dimensional foams

Vincent J. Langlois*
School of Physics, Trinity College Dublin, Dublin 2, Ireland and Department of Physics, Technical University of Denmark, 2800 Kgs. Lyngby, Denmark

Stefan Hutzler and Denis Weaire
School of Physics, Trinity College Dublin, Dublin 2, Ireland

(Received 29 April 2008; published 1 August 2008)

The soft-disk model previously developed and applied by Durian [D. J. Durian, Phys. Rev. Lett. 75, 4780 (1995)] is brought to bear on problems of foam rheology of longstanding and current interest, using two-dimensional systems. The questions at issue include the origin of the Herschel-Bulkley relation, normal stress effects (dilatancy), and localization in the presence of wall drag. We show that even a model that incorporates only linear viscous effects at the local level gives rise to nonlinear (power-law) dependence of the limit stress on strain rate. With wall drag, shear localization is found. Its nonexponential form and the variation of localization length with boundary velocity are well described by a continuum model in the spirit of Janiaud et al. [Phys. Rev. Lett. 97, 038302 (2006)]. Other results satisfactorily link localization to model parameters, and hence tie together continuum and local descriptions.

DOI: 10.1103/PhysRevE.78.021401 PACS number(s): 83.80.Lz, 47.57.Bc

I. INTRODUCTION

A. Foam rheology

While the deformation and flow properties of foams are broadly understood in terms of shear elastic modulus, yield stress, etc. [1–3], many details remain perplexing. A fuller understanding must address both the local forces that operate at the level of the individual films, and the way in which these forces combine to determine the overall response to strain. This paper will be entirely devoted to the second question. It uses a particularly simple representation of bubbles and their mutual forces, as previously developed by Durian [4–7]. This model may be unrealistic in some respects, but its simplicity and computational tractability makes it attractive at a time when more precise descriptions of dynamic properties are lacking.

Our immediate goal is to thoroughly analyze the properties of this model, per se. Since it represents bubbles as soft disks [in the two-dimensional (2D) case] it has some relevance to granular materials as well. As is often the case in foam physics, this study will remain for the time being in two dimensions, which has obvious advantages. The main experimental literature of foam rheology is concerned with ordinary three-dimensional (3D) foams, but in recent years, considerable attention has been focused on their 2D counterparts.

It has turned out that the obvious 2D experimental sample, consisting of foam trapped between two plates [see Fig. 1(c)], has shear properties that are significantly affected by viscous drag forces exerted by the confining plates. We will therefore be concerned with two quite different but related cases: with and without such forces. The latter case can be realized experimentally in 2D as a Bragg raft [Fig. 1(a)] and it is roughly analogous to a 3D foam, because of the absence of confinement-induced forces. Having defined the model, we will deal with this case first, and proceed to introduce the wall forces at a later stage. In summary, the main goals of this work are (1) to extract the parameters of a continuum (Herschel-Bulkley) formalism, by means of simulation, and (2) to use these in a continuum model for 2D shear localization. Both of these goals are satisfactorily realized within the scope of the present calculations.

B. Questions raised by experiments

It is necessary first to review some of the history of foam rheology. The more traditional 3D experiments, for example, those of Khan et al. [8], have used various types of rheometers to explore the relation between stress, strain, and strain rate. This is most straightforward when strain rate does not change sign; that is, increments of shear are always in the same sense, and we shall adopt that restriction here also. Awkward complications arise from hysteretic effects in the more general case [9]. The conclusion has generally been that the data is well accounted for by the Herschel-Bulkley equation. This adds to the quasistatic stress-strain relation a second term proportional to the strain rate $\dot{\gamma}$, raised to the power $a$,

$$\sigma = \sigma_y + c_v \dot{\gamma}^a,$$

where $\sigma$ and $\sigma_y$ denote, respectively, stress and yield stress, $c_v$ is the viscosity component of stress (also called consistency), and $a$ is the Herschel-Bulkley exponent. However, the exact value of $a$ is still subject to debate. Most experiments agree on the shear-thinning behavior of foams, whereas for $a=1$ (the Bingham fluid), the effective viscosity (that is, stress divided by strain rate), tends to constant at high rate of strain, it tends to zero if $a<1$. While Schwartz and Princen [10] predicted theoretically that $a=2/3$, various values between 0.25 and 1 have been measured experimentally (see, for instance, [8,11–16]). Besides, Denkov et al.

*vincent.langlois@fysik.dtu.dk
[14] and Katgert et al. [15] showed, respectively, that \( a \) depends on the properties of the surfactant and on the polydispersity of the foam. The obvious question arises: what determines the value of \( a \)? There has been little understanding of this up until now, and it is one of our stated objectives to explore the question within the context of a simple model.

A second question concerns normal stresses or the related phenomenon of dilatancy. While often mentioned in the context of granular materials, these effects have not been much considered for foams, apart from the work of Weaire and Hutzler [17,18], confined to the quasistatic case. As they noted, there is, in principle, a second and possibly more important dynamic contribution to these effects, well known to rheologists \([19,20]\). So one may ask: what are the dynamic effects in a simple model?

As in much of foam physics, recourse may be made to two-dimensional systems, for the sake of simplicity and transparency. We will not attempt a full review of the various rheological experiments on 2D foams that have been performed over the last five years \([13,15,21,22]\). Equally, we will give no details of the quasistatic calculations \([23,24]\) that have been added to account for them, or the continuum theory \([25–28]\) which has offered an alternative viewpoint. We will argue that the present paper strongly supports the continuum model. Some results of the quasistatic calculations suggest that corrections are needed to account for the role of polydispersity, and we do not rule that out.

**C. Soft sphere or disk model**

In the model developed by Durian around 1995, 2D bubbles are represented by circular disks. When overlapping (and only then) they interact via a simple spring force, the displacement of the spring being the radial overlap (see Fig. 2). The elastic repulsive force \( \mathbf{F}_n \) acting on bubble \( i \) due to bubble \( j \) is given by

\[
\mathbf{F}_n = \kappa \frac{2R_0}{R_i + R_j} \Delta_{ij} \mathbf{n}_{ij},
\]

where \( \kappa \) is the coefficient of elasticity, \( \mathbf{n}_{ij} \) is the normal vector between bubbles \( i \) and \( j \),

\[
\mathbf{n}_{ij} = \frac{\mathbf{r}_i - \mathbf{r}_j}{|\mathbf{r}_i - \mathbf{r}_j|},
\]

and the overlap \( \Delta_{ij} \) (see Fig. 2) is given by

\[
\Delta_{ij} = \begin{cases} 
    (R_i + R_j) - |\mathbf{r}_i - \mathbf{r}_j| & \text{if } (R_i + R_j) < |\mathbf{r}_i - \mathbf{r}_j| \\
    0 & \text{otherwise}.
\end{cases}
\]

\( R_i \) and \( R_j \) are the radii of bubbles \( i \) and \( j \), centered at \( \mathbf{r}_i \) and \( \mathbf{r}_j \), respectively, and \( R_0 \) is the average bubble radius of the entire bubble packing. The ratio \( \frac{2R_0}{R_i + R_j} \) in Eq. (2) takes into account that larger bubbles are easier to deform than smaller ones. In such a model one may define an effective liquid fraction \( \phi \) (which ignores the overlap of disks) as \( \phi = 1 - N(R_i^2 \pi)/A \), where \( A \) is the area of the confinement of the disks and \( N \) is the total number of disks. Since a packing of polydisperse disks loses its mechanical rigidity for \( \phi > 0.16 \), for higher values of \( \phi \) it no longer represents a two-dimensional foam. In all the following, the liquid fraction will be chosen as \( \phi = 0.05 \).

A real flowing foam dissipates energy by viscous friction in the films and Plateau borders separating the bubbles. The films are not explicitly represented in our model. The simplest expression, as used by Durian \([5]\) and adopted here, represents the viscous force \( \mathbf{F}_d \) on bubble \( i \) associated with a neighboring bubble \( j \) as

\[
\mathbf{F}_d = -c_v (\mathbf{v}_i - \mathbf{v}_j),
\]

where \( c_v \) is the dissipation constant for the bubble-bubble interaction and \( \mathbf{v}_i \) and \( \mathbf{v}_j \) are the respective bubble velocities.

![FIG. 2. Overlap \( \Delta_{ij} \) between two contacting bubbles of radii \( R_i \) and \( R_j \), located at \( \mathbf{r}_i \) and \( \mathbf{r}_j \), respectively.](image-url)
defined above. In practice we used the Verlet algorithm, with a bubble mass \( m_b \) small enough to assume that inertial effects are negligible (the ratio \( km_b/c_b^2 \) was set to 0.01). We present results here only for the eventual steady state after long times. The results obtained are significantly different from those of Durian. This may be due to the strong approximation that he used, as indicated above. This approximation had been removed later by Tewari et al. [6], but they were not interested in the rheological properties and concluded that both versions of the model were equivalent.

II. RESPONSE TO SIMPLE SHEAR IN THE ABSENCE OF WALL DRAG

A. Herschel-Bulkley power-law exponent

In this section we deal with the Durian model, with no additional viscous drag force from confining plates. Our first computations concerned the evaluation of its flow properties under simple shear, in particular, the value of the Herschel-Bulkley power-law exponent \( a \) in Eq. (1). To this purpose we generated assemblies of 1000–10 000 bubbles in a rectangular confinement, as shown in Fig. 3, using periodic boundary conditions in the horizontal direction. The bubbles at the upper boundary, which are treated as attached to this boundary, are given a constant velocity \( V \). This corresponds to the application of strain at a constant rate \( \gamma = V/H \) for a sample of width \( H \) (see Fig. 3). We will only consider polydisperse foams, the bubble radii having a uniform distribution within the range \( R = R_0 (1 \pm 0.15) \).

At this stage it is convenient to introduce a dimensionless Deborah number \( \text{De} = \gamma c_b/\kappa \), which is defined as the ratio of the characteristic time of the material that is sheared to the characteristic time of the deformation process [29]. We identify the latter with \( \gamma^{-1} \), and the material time scale we relate to the competition of energy storage and dissipation at the level of bubble-bubble interactions. This then results in the definition \( \text{De} = \gamma c_b/\kappa \).

Under an imposed boundary shear, after a transient, the system of disks reaches a steady average state. It is characterized by a linear average velocity profile of the disks with regard to their vertical position in the sample, so there is no localization in the present case; we may proceed to extract the constitutive law in a straightforward way. We determine the stress at the moving boundary as a function of \( V \). In Fig. 4 we display this variation as a function of the dimensionless Deborah number \( \text{De} \). A least-square fit of the data with the Herschel-Bulkley form results in

\[
\sigma/\kappa = 0.0043 + 0.26\text{De}^{0.54 \pm 0.01}.
\]

The model foam thus exhibits a strongly nonlinear, shear-thinning rheological behavior, despite the linearity of all local forces. This can be attributed to the importance of disorder in such a polydisperse jammed system: the velocity fluctuations cannot be neglected and make the trajectories of the bubbles strongly nonaffine. Therefore the simple image of bubble layers sliding over each other is misleading. In the initial version of the model applied by Durian [4,5], the mean-field approximation effectively subtracts these disordered motions, which results in a Bingham rheology \((a = 1)\).

Let us note that the nonlinear behavior does not affect the average velocity profile in the simple shear geometry we adopted (as opposed to what happens in a cylindrical Couette geometry [13]). The shear localization that is discussed in later sections is not seen here.

The value of the Herschel-Bulkley exponent we obtain [Eq. (1)] \( a = 0.54 \pm 0.01 \) is roughly consistent with the various experimental measurements already mentioned, that showed a nonlinear, shear-thinning behavior [8,11–16].
increase its liquid fraction under shear has been qualitatively reported in experiments [31], we do not know of any quantitative experimental measurements of the dynamic variation of the osmotic pressure, which is particularly difficult to obtain in 2D. Rheologists sometimes cite the old work of Bag-nold [19] as indicating a quadratic dependence, very different from the power law in Eq. (7). We have not proceeded any further with the analysis of the normal stress.

III. EFFECT OF WALL DRAG

A. Adding wall drag

As we have already mentioned, certain experiments with a 2D foam between two glass plates exhibit a new type of flow; when one boundary is moved to produce shear, the resulting shear is exponentially localized at that boundary [22]. According to the continuum model of Janiaud et al. [25–27] this is to be understood as a direct effect of the wall drag, which we will now introduce into the numerical model. This continuum model in its original form assumed a constitutive equation of the Bingham type (\(\alpha=1\)), and added in the drag force \(F_w\), as a body force proportional to local velocity. This predicted an exponential localization of the flow along the moving boundary, the width of the shear band being independent of the driving velocity.

To mimic that theoretical model in our simulations, we now add to the forces acting on bubble \(i\), a wall drag force \(F_w\), given in the most simple form by

\[
F_w = -c_w |\mathbf{v}_i| \frac{\mathbf{v}_i}{|\mathbf{v}_i|},
\]

where \(c_w\) is a drag constant. According to Bretherton [32] and Denkov et al. [14], for surfactants with low surface viscosities, the friction of the foam due to the wall is characterized by \(b=2/3\). However, to keep all the ingredients of our model linear, we will adopt \(b=1\). The resulting equation of motion is again solved numerically, using the second order Verlet method. Note that we have just established in Sec. II A that the appropriate constitutive law is not that of Bingham, so we will have cause to return to this point again. As we shall see, the results for the steady shear at long times exhibit localization, which conforms well to the prediction of the continuum model.

B. Flow localization

Shear simulations with added wall drag were performed on an assembly of 10 000 soft disks, for constant values of \(\chi=c_b/c_w\), the ratio of dissipation and drag constant. Once a steady state was reached, we determined the velocity profile between static and moving boundary by performing time averages over the horizontal velocity components of the bubbles. While in the absence of wall friction \((c_w=0\), i.e., \(\chi \rightarrow \infty\)), the velocity profile is roughly linear throughout the sample, this is no longer the case for finite values of \(\chi\). Instead, flow is localized near the moving boundary, as shown in Fig. 6. A localization length \(\lambda_{1/10}\) can be arbitrarily defined by measuring the distance from the moving boundary at which the velocity reaches 1/10th of its value at the
The velocity field is enhanced as wall friction is increased, i.e., with decreasing $\chi$. The solid lines, which agree closely with the data, were obtained numerically from the continuum model, using the Herschel-Bulkley index $a=0.54$. Such a power-law scaling can also be found theoretically by modifying the continuum model of Janiaud et al. [25] to incorporate the Herschel-Bulkley rheology we exhibited here: the velocity field $v(y)$ has to be such that the internal dissipation is exactly balanced by the friction along the confining plates; that is,

$$c_d \frac{dv}{dy} = c_d(v(y))^b. \tag{10}$$

Here $c_d$ is the wall drag constant $c_w$ of Eq. (8) per unit area, i.e., $c_d = c_w/(\pi R_0^2)$ and $b=1$ for the case of the simple linear form for the drag force considered here. The exponent $a$ is the one from the Herschel-Bulkley relationship, i.e., $a = 0.54$ in our case [see Eq. (6)], and $c_v$ is the viscosity component of stress.

Figure 8 shows the evolution of $\lambda_{1/10}$ with the Deborah number $De$, when varying the driving velocity $V$. A least-square fit results in

$$\lambda_{1/10}/R_0 = 0.87De^{-0.30}. \tag{9}$$

Such a power-law scaling can also be found theoretically by modifying the continuum model of Janiaud et al. [25] to incorporate the Herschel-Bulkley rheology we exhibited here: the velocity field $v(y)$ has to be such that the internal dissipation is exactly balanced by the friction along the confining plates; that is,

$$c_d \frac{dv}{dy} = c_d(v(y))^b. \tag{10}$$

Here $c_d$ is the wall drag constant $c_w$ of Eq. (8) per unit area, i.e., $c_d = c_w/(\pi R_0^2)$ and $b=1$ for the case of the simple linear form for the drag force considered here. The exponent $a$ is the one from the Herschel-Bulkley relationship, i.e., $a = 0.54$ in our case [see Eq. (6)], and $c_v$ is the viscosity component of stress. This equation can be solved to predict the localization of the flow, as shown theoretically in [28]: the (exponentially defined) localization length $l$ is given by

$$l \approx \left(\frac{c_v}{c_d}\right)^{1+a} \nu^{-b/1+a}, \tag{11}$$

provided that $l$ is much less than the sample size (the two definitions of the localization length are related by $\lambda_{1/10} = \ln(10)l$; see the Appendix). This immediately gives the scaling $\lambda_{1/10} \propto c_w^{0.65}$, in excellent agreement with the data shown in Fig. 7. However, the success of the continuum model is not confined to such scaling relationships, but is fully quantitative. Rewriting Eq. (11) in terms of the Deborah number gives

$$l \approx \left(\frac{c_v}{c_d}\right)^{1+a} \nu^{-b/1+a}, \tag{12}$$

Inserting our simulation input parameters for $c_{d,fr} H$, $c_{fr} R_0$, and $b=1$, together with our numerical results for the values of $c_v$ and $a=0.54$, we obtain

$$l/R_0 = 0.36De^{-0.30}. \tag{13}$$

This corresponds to $\lambda_{1/10}/R_0 = 0.83De^{-0.30}$, which is also in excellent agreement with the numerical results (see Fig. 8). We have concentrated on the localization length here, but the full velocity profile, which turns out to be only approximately exponential, can be numerically calculated from it. As Fig. 6 shows, the result is in close agreement with the data.
IV. CONCLUSIONS

By means of numerical simulations based on the soft-disk model, we have shown that the Herschel-Bulkley rheology of a 2D foam can be derived from a discrete model with good agreement with experiments. Our model also predicts a dynamic dilatancy, that is, a tendency of the foam to increase its volume when sheared. Finally, we added to the classical bubble model the viscous friction experienced by a foam under confinement, which resulted in the formation of shear bands, as had been observed in experiments and derived with a continuum model. Adapting this rheological model to our parameters led to a consistent picture of the link between the local and global descriptions. This must be regarded as a strong indication for the theory, but it will be important to examine the limits of that conclusion.

Previous quasistatic simulations [24] suggest that disorder (polydispersity) plays a role. So far, our interpretation is that wall drag is responsible for localization, as prescribed by the continuum model, but that the eventual localization length contains an additional contribution from polydispersity, not contained within the continuum description.

In further studies, we will explore this by concentrating on the limit in which \( \lambda \to 0 \) for the continuum model, and seek to identify the effect of polydispersity, which is expected to be significant in that limit. It will also be informative to repeat simulations with different local forces, to try to establish the precise relation (if any) of the Herschel-Bulkley parameters to these forces. Finally, if our conclusions are sustained for this and similar models, the core theoretical problem of the origin of the Herschel-Bulkley nonlinearity will remain.

ACKNOWLEDGMENTS

This work was supported by the European Space Agency (MAP AO-99-108:C14914/02/NL/SH and MAP AO-99-075:C14308/00/NL/SH) and Science Foundation Ireland (RFP05/RFP/PHY0016). We also acknowledge the European Union COST P21 action on “The Physics of Droplets.” V.J.L. would like to acknowledge the Grenoble Foam Mechanics Workshop and to thank G. Katgert and M. van Hecke for fruitful discussions, and for providing some of their results before publication.

APPENDIX: DEFINITION OF LOCALIZATION LENGTH

The precise definition of localization length is, in general, arbitrary. For the analysis of our numerical results we chose to use

\[ \nu(\lambda_{1/10}) = V/10, \]

(A1)

where \( \nu(y) \) is the local average bubble velocity, measured a distance \( y \) away from the boundary, which moves at velocity \( V \).

In the continuum theory an alternative definition was used [28]

\[ \nu(l) = V/e. \]

(A2)

There is no general relation between these two, but since localization is, in general, approximately exponential, we may use

\[ \lambda_{1/10} = \ln 10l \approx 2.30l \]

(A3)

to adjust the theoretical prediction. We have adopted this procedure in Sec. III B.

[1] A. M. Krainik, Annu. Rev. Fluid Mech. 20, 325 (1988).
[2] D. Weaire and M. A. Fortes, Adv. Phys. 43, 685 (1994).
[3] R. Höhler and S. Cohen-Addad, J. Phys.: Condens. Matter 17, R1041 (2005).
[4] D. J. Durian, Phys. Rev. Lett. 75, 4780 (1995).
[5] D. J. Durian, Phys. Rev. E 55, 1739 (1997).
[6] S. Tewari, D. Schiemann, D. J. Durian, C. M. Knobler, S. A. Langer, and A. J. Liu, Phys. Rev. E 60, 4385 (1999).
[7] I. K. Ono, S. Tewari, S. A. Langer, and A. J. Liu, Phys. Rev. E 67, 061503 (2003).
[8] S. A. Khan, C. A. Schnepper, and R. C. Armstrong, J. Rheol. 32, 69 (1988).
[9] P. Marmottant and F. Graner, Eur. Phys. J. E 23, 337 (2007).
[10] L. W. Schwartz and H. M. Princen, J. Colloid Interface Sci. 118, 201 (1987).
[11] H. M. Princen and A. D. Kiss, J. Colloid Interface Sci. 128, 176 (1989).
[12] A. D. Gopal and D. J. Durian, J. Colloid Interface Sci. 213, 169 (1999).
[13] J. Lauridsen, G. Chanon, and M. Dennin, Phys. Rev. Lett. 93, 018303 (2004).
[14] N. D. Denkov, V. Subramanian, D. Gurovich, and A. Lips, Colloids Surf., A 263, 129 (2005).
[15] G. Katgert, M. E. Möbius, and M. van Hecke, Phys. Rev. Lett. 101, 058301 (2008).
[16] N. N. Thondavadi and R. Lemlich, Ind. Eng. Chem. Process Des. Dev. 24, 748 (1985).
[17] D. Weaire and S. Hutzler, Philos. Mag. 83, 2747 (2003).
[18] F. Rioual, S. Hutzler, and D. Weaire, Colloids Surf., A 263, 117 (2005).
[19] R. A. Bagnold, The Physics of Blown Sand and Desert Dunes (Methuen, London, 1941).
[20] D. E. Wolf, in Computational Physics, edited by K. H. Hoffmann and M. Schreiber (Springer, Berlin, 1996), pp. 64–95.
[21] G. Debrégeas, H. Tabuteau, and J.-M. di Meglio, Phys. Rev. Lett. 87, 178305 (2001).
[22] Y. Wang, K. Krishan, and M. Dennin, Phys. Rev. E 73, 031401 (2006).
[23] A. Kabla and G. Debrégeas, Phys. Rev. Lett. 90, 258303 (2003).
[24] A. Wyn, I. T. Davies, and S. J. Cox, Eur. Phys. J. E 26, 81 (2008).
[25] J. Janiaud, D. Weaire, and S. Hutzler, Phys. Rev. Lett. 97, 038302 (2006).
[26] R. J. Clancy, E. Janiaud, D. Weaire, and S. Hutzler, Eur. Phys. J. E. 21, 123 (2006).
[27] E. Janiaud, D. Weaire, and S. Hutzler, Colloids Surf., A 309, 125 (2007).
[28] D. Weaire, S. Hutzler, V. J. Langlois, and R. J. Clancy, Philos. Mag. Lett. (to be published).
[29] J. F. Hutton, H. A. Barnes, and K. Walters, An Introduction to Rheology (Elsevier, Amsterdam, 1989).
[30] N. D. Denkov, S. Tcholakova, K. Golemanov, K. P. Ananthapadmanabhan, and A. Lips, Phys. Rev. Lett. 100, 138301 (2008).
[31] P. L. Marze, A. Saint-Jalmes, and D. Langevin, Colloids Surf., A 263, 121 (2005).
[32] F. P. Bretherton, J. Fluid Mech. 10, 166 (1961).