

Viscoplastic rimming flow inside a rotating cylinder

Thomasina V. Ball¹ and Neil J. Balmforth²

¹*Mathematics Institute, University of Warwick, Coventry CV4 7AL, United Kingdom*

²*Department of Mathematics, University of British Columbia, Vancouver, British Columbia V6T 1Z4, Canada*

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A theoretical analysis is presented for the flow of a Herschel-Bulkley fluid around the inside surface of a rotating cylinder, at rotation speeds for which the fluid largely collects in a prominent pool in the lower part of the cylinder. The analysis, based on lubrication theory, predicts the steady states typically reached after a small number of rotations. The analysis is also modified to consider the drainage of the film around a stationary cylinder, which allows an exploration of the dynamics when a rotating drum is suddenly stopped. The predictions of the theory are compared with experiments in which a Carbopol suspension is rotated inside an acrylic drum.

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I. INTRODUCTION

The flow of a viscous fluid around the surface of a rotating cylinder has proved a popular configuration in which to study the fluid mechanics of coating flows [1,10,28,33,37,39]. Beginning with Moffatt [37], an initial point of interest was the development of a corner and shocks in leading-order lubrication theory, as a more uniform coating on the inside of a horizontal cylinder became increasingly influenced by gravity [28,39]. Such irregularities in the shape of the free surface become smoothed by surface tension, higher-order gravity effects, or over lengthscales violating the scale separation implicit in lubrication theory [1,10,33,44]. Instabilities leading to three-dimensional flow structures have also been of interest as illustrations of pattern formation [11,25,42,43].

A key drawback of the leading-order lubrication model used to study flows with weak gravitational effects is that it fails to deal with situations in which gravity plays a dominant role, dragging most of the coating fluid around the rotating cylinder into a pool sitting near the lowest point. This situation led Tirumkudulu and Acrivos [44] and Ashmore, Hosoi, and Stone [1] to restructure the lubrication model, adding some important higher-order terms (incorporating either surface tension, gravity corrections, or both). The restructured model accounts for pooling at the bottom of the cylinder, together with the thin film that becomes drawn out of the pool, then coats the remainder of the cylinder surface before reentering the pool. The draw-out and reentry of the film at the edges of the pool correspond to the classical Landau-Levich problem and its variants [1,10,14,49].

Despite the prevalence of analyses employing Newtonian fluid, the coating flows that feature in a large number of industrial processes [48] often involve complex fluids. Partly with this in mind, a number of studies have explored rimming flows of viscoelastic fluids and suspensions (see Ref. [42]). With more relevance to the present work, Johnson [28] considered coating by power-law fluids and Ross, Duffy, and Wilson [41] generalized Moffatt's analysis to allow the coating fluid to have a yield stress. Neither of these studies, however, deal with situations with gravitational pooling at the bottom of the cylinder. When pooling does arise, that reservoir of fluid is not motionless, but continues to overturn driven by the rotation. In fact, for viscoplastic fluid, the resulting states have much in common with the steady surges encountered in laboratory experiments using inclined conveyor belts, designed to simulate geophysical flows [15,16,21,31]. The rotating drum arguably

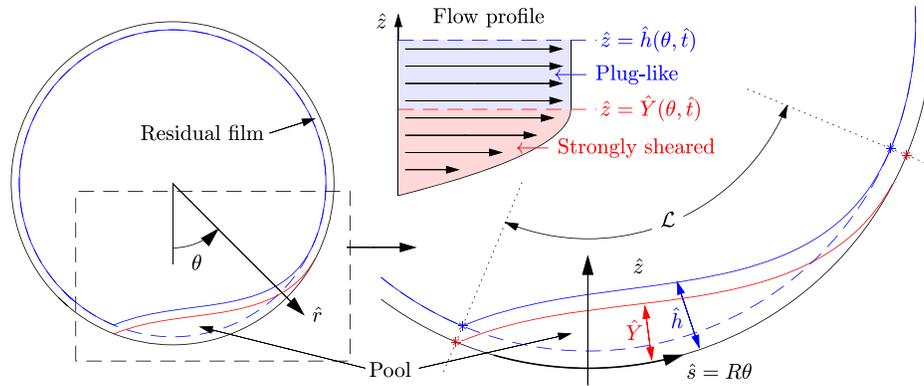


FIG. 1. A sketch showing the geometry of the rimming flow. The red and blue stars mark the ends of the “pool” over which most of the fluid collects and the depth \hat{h} varies significantly; beyond these points there is a residual, potentially plugged up film. The pool has a typical length $\mathcal{L} \ll R$. Lubrication theory predicts that the flow profile takes a distinctive form in which a strongly sheared layer intervenes between the drum and a plug-like region, the “pseudoplug,” bordering the free surface; the strongly sheared layer has thickness \hat{Y} . Beyond the pool, one anticipates a residual, plugged-up film to coat the remainder of the cylinder (with $\hat{Y} = 0$).

provides a simpler configuration to explore such surges, providing a key motivation for the present paper.

We organize the paper as follows: in Sec. II we mathematically formulate the problem and use a shallow-layer approximation (lubrication theory) to derive a reduced model for the flow of a thin film of a viscoplastic fluid described by a model constitutive law (the Herschel-Bulkley law [6]) along a circular surface (cf. Refs. [4,29]). That said, our formulation is more general, applying to any nearly horizontal, slowly varying surface and may therefore prove useful in other settings. However, by avoiding any geometrical approximations in the gravitational terms, the model also recovers the viscoplastic generalization of the leading-order lubrication analysis for a more uniform coating [37,41] in a different limit of the model’s dimensionless parameters. Thus, by bridging between the two asymptotic limits, the model captures the dynamics of both the pool at the bottom of the cylinder and the residual film coating it elsewhere (cf. Ref. [44]). In Sec. III, we use the model to explore the steady rimming solutions that are predicted to develop within a rotating drum. We follow this up in Sec. IV, with a discussion of the dynamics when the cylinder does not rotate, which allows us to explore what happens when a rotating cylinder is suddenly stopped. In particular, we consider how the coating drains to rest, which inevitably becomes controlled by the yield stress (if one is present). The theoretical results are then compared with some laboratory experiments in Sec. V. We conclude in Sec. VI. The Appendix contains some details of the theoretical construction of steady rimming flows without surface tension (which is otherwise included in our model formulation).

II. THE MODEL

A. Formulation

Consider the flow of a two-dimensional, incompressible film of Bingham fluid on the inside of a cylinder of radius R , rotating anticlockwise with rate Ω . We neglect inertia, and represent the geometry of the cylinder using polar coordinates system (\hat{r}, θ) , with $\theta = 0$ corresponding to the vertically down direction, as illustrated in Fig. 1(a). The shallow fluid layer coating the inner surface is described using a local curvilinear coordinate system (\hat{s}, \hat{z}) in which $\hat{s} = R\theta$ is arc length along

71 the cylinder and $\hat{z} = R - \hat{r}$ [see Fig. 1(b)]. The fluid velocity and pressure are $(\hat{u}(\hat{s}, \hat{z}), \hat{w}(\hat{s}, \hat{z}))$ and
 72 $\hat{p}(\hat{s}, \hat{z})$, satisfying

$$\frac{\partial \hat{u}}{\partial \hat{s}} + (1 - \kappa \hat{z}) \frac{\partial \hat{w}}{\partial \hat{z}} - \kappa \hat{w} = 0, \quad (1)$$

$$\frac{\partial \hat{\tau}_{ss}}{\partial \hat{s}} - \frac{\partial \hat{p}}{\partial \hat{s}} + (1 - \kappa \hat{z}) \frac{\partial \hat{\tau}_{sz}}{\partial \hat{z}} - 2\kappa \hat{\tau}_{sz} = \rho g(1 - \kappa \hat{z}) \sin \theta, \quad (2)$$

$$\frac{\partial \hat{\tau}_{sz}}{\partial \hat{s}} + (1 - \kappa \hat{z}) \left(\frac{\partial \hat{\tau}_{zz}}{\partial \hat{z}} - \frac{\partial \hat{p}}{\partial \hat{z}} \right) + \kappa (\hat{\tau}_{ss} - \hat{\tau}_{zz}) = \rho g(1 - \kappa \hat{z}) \cos \theta, \quad (3)$$

73 where the curvature is

$$\kappa = \frac{\partial \theta}{\partial \hat{s}} \equiv R^{-1}, \quad (4)$$

74 $\hat{\tau}_{ij}$ is the deviatoric stress tensor, and subscripts (in Roman font) indicate tensor components. The
 75 Bingham law relates the stress to the strain rates by

$$\hat{\tau}_{ij} = \left(\mu + \frac{\tau_Y}{\hat{\gamma}} \right) \hat{\gamma}_{ij} \quad \text{for} \quad \hat{\tau} > \tau_Y, \quad (5)$$

76 and $\hat{\gamma}_{ij} = 0$ otherwise, where μ is the (plastic) viscosity, τ_Y is the yield stress, and

$$\{\hat{\gamma}_{ij}\} = \begin{pmatrix} 2(\hat{u}_s - \kappa \hat{w})/(1 - \kappa \hat{z}) & (\hat{w}_s + \kappa \hat{u})/(1 - \kappa \hat{z}) + \hat{u}_z \\ (\hat{w}_s + \kappa \hat{u})/(1 - \kappa \hat{z}) + \hat{u}_z & 2\hat{w}_z \end{pmatrix}. \quad (6)$$

77 Here subscripts of \hat{s} and \hat{z} on the velocity components denote partial derivatives, and $\hat{\gamma} =$
 78 $\sqrt{\frac{1}{2} \sum_{ij} \hat{\gamma}_{ij} \hat{\gamma}_{ij}}$ and $\hat{\tau} = \sqrt{\frac{1}{2} \sum_{ij} \hat{\tau}_{ij} \hat{\tau}_{ij}}$ denote the tensor second invariants.

79 We have further employed a hat notation to denote dimensional variables; this decoration is
 80 removed using suitable scalings in our nondimensionalization of the problem below.

81 At the free surface of the fluid film, $\hat{z} = \hat{h}(\hat{s}, \hat{t})$, ignoring surface tension,

$$(1 - \kappa \hat{h}) \hat{\tau}_{sz} - (\hat{\tau}_{ss} - \hat{p}) \frac{\partial \hat{h}}{\partial \hat{s}} = (1 - \kappa \hat{h})(\hat{\tau}_{zz} - \hat{p}) - \hat{\tau}_{sz} \frac{\partial \hat{h}}{\partial \hat{s}} = 0 \quad (7)$$

82 and

$$\hat{w} = \frac{\partial \hat{h}}{\partial \hat{t}} + \frac{\hat{u}}{1 - \kappa \hat{h}} \frac{\partial \hat{h}}{\partial \hat{s}}. \quad (8)$$

83 On the cylinder, no slip demands that

$$\hat{u} = R\Omega \quad \text{and} \quad \hat{w} = 0. \quad (9)$$

84 B. Scaling

85 The pool at the bottom of the cylinder has a characteristic length \mathcal{L} and a typical depth \mathcal{H} . The
 86 fluid is shallow, so that $\epsilon = \mathcal{H}/\mathcal{L} \ll 1$, and the pool occupies a small fraction of the circumference,
 87 $\mathcal{L} \ll R$; see Fig. 1. More specifically, to account properly for the effect of gravity on the pool,
 88 the gravitational force along the surface of the cylinder, $\rho g \sin \theta \sim \rho g \hat{s}/R$, must be comparable to
 89 hydrostatic pressure gradients in the fluid, $\hat{p}_\theta \sim \rho g \frac{\partial \hat{h}}{\partial \hat{s}}$.

90 Balancing the magnitudes of these two terms indicates that $\mathcal{H} = O(\epsilon^2 R)$ or $\mathcal{L} = O(\epsilon R)$. Further-
 91 more, when the pool contains a significant fraction of the fluid, the pool's area, of $O(\mathcal{H}\mathcal{L})$, must be
 92 comparable to the total fluid area, $\pi R^2 f$, where $f \ll 1$ is the ‘‘fill fraction’’ of the drum. In other
 93 words, $\epsilon = O(f^{\frac{1}{3}})$.

94 We now remove the dimensions from the equations by defining the new variables,

$$s = \frac{\hat{s}}{\mathcal{L}}, \quad (z, h) = \frac{(\hat{z}, \hat{h})}{\mathcal{H}}, \quad u = \frac{\hat{u}}{R\Omega}, \quad w = \frac{\hat{w}}{\epsilon R\Omega}, \quad t_* = \frac{R\Omega \hat{t}}{\mathcal{L}}, \quad (10)$$

$$p = \frac{\hat{p}}{\rho g \mathcal{H}}, \quad \tau_{ij} = \frac{\hat{\tau}_{ij}}{\epsilon \rho g \mathcal{H}}, \quad \dot{\gamma}_{ij} = \frac{\hat{\dot{\gamma}}_{ij} \mathcal{H}}{R \Omega}. \quad (11)$$

To leading order, the dimensionless equations become

$$u_s + w_z = 0, \quad (12)$$

$$\frac{\partial \tau_{sz}}{\partial z} - \frac{\partial p}{\partial s} = g_s \equiv \frac{\sin \theta}{\epsilon}, \quad (13)$$

$$-\frac{\partial p}{\partial z} = g_z \equiv \cos \theta. \quad (14)$$

Because the pool has a limited angular extent, $\epsilon^{-1} \sin \theta \sim s = O(1)$, and the gravity components (g_s, g_z) are both order one. Rather than introducing their leading order replacements, $(s, 1)$, however, we retain the trigonometric functions in (g_s, g_z) to make contact with existing literature and allow the model to capture the dynamics of the residual film coating the cylinder outside the pool (see Sec. II E).

The dimensionless, yielded version of the constitutive law is

$$\tau_{ij} = \left(\Lambda_*^{-1} + \frac{\text{Bi}_*}{\dot{\gamma}} \right) \dot{\gamma}_{ij}, \quad \{\dot{\gamma}_{ij}\} = \begin{pmatrix} 2\epsilon u_s + O(\epsilon^3) & u_z + O(\epsilon^2) \\ u_z + O(\epsilon^2) & 2\epsilon w_z \end{pmatrix}, \quad (15)$$

where the dimensionless parameters,

$$\Lambda_* = \frac{\rho g \mathcal{H}^3}{\mu \mathcal{L} R \Omega} \quad \text{and} \quad \text{Bi}_* = \frac{\tau_y \mathcal{L}}{\rho g \mathcal{H}^2}. \quad (16)$$

Provided $u_z = O(1)$, Eq. (15) indicates that the shear stress dominates the other deviatoric stress components and the yield condition reduces to $|\tau_{sz}| > \text{Bi}_*$. In this situation, the fluid is relatively strongly sheared. There is, however, a second way to deal with the constitutive relation: if we take $u_z = O(\epsilon)$, we must retain all the components of the deformation rate tensor at the same $O(\epsilon)$ order. This leads to the perfectly plastic relation $\tau_{ij} = \text{Bi}_* \dot{\gamma}_{ij} / \dot{\gamma}$ to leading order. Because this applies where $u_z = O(\epsilon)$, $|\tau_{sz}| < \text{Bi}_*$ and $\sqrt{\tau_{ss}^2 + \tau_{sz}^2} \sim \text{Bi}_*$, the flow here is plug-like, but still yielded; this is Walton and Bittleston's pseudoplug [2,47]. It turns out that shallow viscoplastic flows typically contain both a strongly sheared region and a pseudoplug [2,4].

Finally, the surface boundary conditions reduce to

$$p = \tau_{sz} = 0 \quad \text{and} \quad w = \frac{\partial h}{\partial t_*} + u \frac{\partial h}{\partial s}, \quad \text{on} \quad z = h. \quad (17)$$

The dimensionless no-slip conditions are

$$u = 1 \quad \text{and} \quad w = 0 \quad \text{on} \quad z = 0. \quad (18)$$

C. Reduction

Equations (13) and (14) imply that

$$p = g_z(h - z) \quad \text{and} \quad \tau_{sz} = -(g_s + g_z h_s)(h - z) \quad (19)$$

[bearing in mind that $p_s = g_z h_s + O(\epsilon^2)$]. Unless the entire film becomes plugged up, the flow therefore splits up into a sheared region near the cylinder, where

$$|\tau_{sz}| > \text{Bi}_*, \quad u_z = -\Lambda_*(g_s + g_z h_s)(Y - z), \quad Y = h - \frac{\text{Bi}_*}{|g_s + g_z h_s|}, \quad 0 < z < Y, \quad (20)$$

and an overlying pseudoplug with

$$|\tau_{sz}| < \text{Bi}_*, \quad u_z = O(\epsilon), \quad Y < z < h \quad (21)$$

118 (cf. Fig. 1). This furnishes the velocity profile,

$$u = \begin{cases} 1 - \frac{1}{2}\Lambda_*(g_s + g_z h_s)(2Y - z)z, & 0 < z < Y, \\ u_p = 1 - \frac{1}{2}\Lambda_*(g_s + g_z h_s)Y^2, & Y < z < h, \end{cases} \quad (22)$$

119 and the angular flux

$$q_*(s, t_*) = \int_0^h u \, dz = h - \frac{1}{6}\Lambda_*(g_s + g_z h_s)Y^2(3h - Y). \quad (23)$$

120 Mass conservation, integrated over the thickness of the film, then demands that

$$\frac{\partial h}{\partial t_*} + \frac{\partial q_*}{\partial s} = 0. \quad (24)$$

121 When Y , as defined in Eq. (20), approaches zero the pseudoplug is predicted to fill the entire layer
 122 and the velocity vanishes. The condition, $h \leq \text{Bi}_*/|g_s + g_z h_s|$, then corresponds to the creation of
 123 a genuine plug attached to the cylinder. In this situation, the flux q_* must equal h . To recover the
 124 correct flux from Eq. (23) we therefore replace the definition of Y in Eq. (20) by

$$Y = \text{Max}\left(0, h - \frac{\text{Bi}_*}{|g_s + g_z h_s|}\right). \quad (25)$$

125 Equations (23)–(25) adapt the shallow viscoplastic fluid model of Liu and Mei [4,29] to the
 126 rotating drum configuration. As written, however, the model in Eqs. (23) and (24) is more general,
 127 applying to shallow viscoplastic flow over an arbitrarily shaped, nearly horizontal surface, as long
 128 as one revises the connection between arc length and angle [from $\hat{s} = R\theta$ to $\theta_s = \kappa(\hat{s})$].

129 D. Inclusion of surface tension and shear thinning

130 When surface tension is included, we must modify the normal force condition at $z = h$. In
 131 particular, the surface pressure, $p = P$ at $z = h$, no longer vanishes and we replace Eq. (19) with

$$p = P + g_z(h - z) \quad \text{and} \quad \tau_{sz} = -(g_s + g_z h_s + P_s)(h - z), \quad (26)$$

132 where the surface pressure is

$$P = \text{const.} - \frac{\sigma h_{ss}}{\rho g \mathcal{L}^2}, \quad (27)$$

133 the surface tension is σ and we take the dimensionless group $\sigma/(\rho g \mathcal{L}^2)$ (related to a Bond number)
 134 to be order one. The reduction of the equations now follows as before, but for the addition of P_s to
 135 the hydrostatic pressure gradient $g_s + g_z h_s$.

136 If we use the Herschel-Bulkley law instead of Bingham, the key change is the replacement of the
 137 effective viscosity $(\mu + \tau_y \hat{\gamma}^{-1})$ in Eq. (5) with $(K \hat{\gamma}^{n-1} + \tau_y \hat{\gamma}^{-1})$. The consistency K and power-
 138 law index n can now be employed to redefine the viscosity scale as $\mu = K(R\Omega/\mathcal{H})^{n-1}$, and the
 139 nondimensionalization of the model then proceeds as above. The dimensionless viscosity function
 140 is $[(\dot{\gamma}/\Lambda_*)^n + \text{Bi}_*] \dot{\gamma}^{-1}$, with

$$\Lambda_* = \Omega^{-1} \left(\frac{\epsilon \rho g R}{K} \right)^{\frac{1}{n}} \left(\frac{\mathcal{H}}{R} \right)^{1 + \frac{1}{n}}, \quad (28)$$

141 and the flux in Eq. (23) becomes

$$q_* = h - \frac{n\Lambda_* Y^{1+\frac{1}{n}} [(1+2n)h - nY]}{(n+1)(2n+1)} |g_s + g_z h_s + P_s|^{\frac{1}{n}} \text{sgn}(g_s + g_z h_s + P_s). \quad (29)$$

More details of the reduction for the Herschel-Bulkley model can be found in Refs. [4,31] There is also a relatively thin buffer between the strongly sheared region and the pseudoplug over which the solutions for the two become matched, but which needs no explicit consideration [2,20,47].

E. Summary

We now return to the original angular variable θ in place of s , which has the awkward feature that the length of the pool \mathcal{L} is selected as the scale during nondimensionalization, but this is not known ahead of time. We further follow Ashmore *et al.* [1] in defining some new variables based on the fill fraction of the cylinder, f , which for a shallow film is

$$f = \int_0^{2\pi} \hat{h}(\theta, \hat{t}) \frac{d\theta}{\pi R}. \quad (30)$$

In particular, we set

$$t = \Omega \hat{t} \equiv \frac{\mathcal{L} t_*}{R}, \quad \eta(\theta, t) = \frac{\hat{h}}{Rf} \equiv \frac{\mathcal{H}h}{Rf}, \quad \mathcal{Y}(\theta, t) = \frac{\mathcal{H}Y}{Rf}, \quad q(\theta, t) = \frac{\mathcal{H}q_*}{Rf}, \quad (31)$$

which indicate that

$$\eta_t + q_\theta = 0, \quad q = \mathcal{Q}(\eta, \Gamma) = \eta - \frac{\Lambda \mathcal{Y}^{1+\frac{1}{n}} |\Gamma|^{\frac{1}{n}}}{n+1} [(1+2n)\eta - n\mathcal{Y}] \operatorname{sgn}(\Gamma), \quad (32)$$

along with the constraint,

$$1 = \int_0^{2\pi} \eta(\theta, t) \frac{d\theta}{\pi}, \quad (33)$$

where

$$\Gamma = \sin \theta + f \eta_\theta \cos \theta - S \eta_{\theta\theta\theta}, \quad \mathcal{Y} = \operatorname{Max}\left(0, \eta - \frac{\operatorname{Bi}}{|\Gamma|}\right), \quad (34)$$

$$\Lambda = \frac{n \Lambda_* \mathcal{L}^{\frac{1}{n}} (fR)^{1+\frac{1}{n}}}{(2n+1) \mathcal{H}^{1+\frac{2}{n}}} = \frac{n f^{1+\frac{1}{n}}}{(2n+1) \Omega} \left(\frac{\rho g R}{K}\right)^{\frac{1}{n}}, \quad \operatorname{Bi} = \frac{\tau_v}{\rho g R f} \equiv \frac{\mathcal{H}^2 \operatorname{Bi}_*}{R \mathcal{L} f}, \quad S = \frac{\sigma f}{\rho g R^2}, \quad (35)$$

and we have extended our subscript notation to partial time and angle derivatives.

A final modification that is possible is to make the nonasymptotic replacement,

$$S \eta_{\theta\theta\theta} \rightarrow \left[\frac{S(\eta + \eta_{\theta\theta})}{(1 + f^2 \eta_\theta^2)^{3/2}} \right]_\theta, \quad (36)$$

for the surface tension term in Eq. (34), a commonly employed device that seeks to better incorporate surface curvature [1,49]. We follow this suit in the computations reported below, although it has little impact on the solutions presented. Indeed, surface tension altogether has a relatively minor effect. That said, the evolution equation (32) is problematic when $S = 0$ because the resulting nonlinear diffusion equation has a negative diffusivity (being proportional to $\cos \theta$) in the upper half of the cylinder. This awkward behavior is conveniently removed when surface tension is included, as in the Newtonian version of the problem [1,10].

With the definition of the new variables and parameters, the scales \mathcal{H} and \mathcal{L} now disappear from the formulation of the model (indicating that it is no longer necessary to select either). Nevertheless, our asymptotic reduction applies in a certain limit of the problem corresponding to $\epsilon = \frac{\mathcal{H}}{\mathcal{L}} \ll 1$. In particular, in our shallow-layer limit, the dimensionless parameters and variables in Eqs. (32) and (33) are not all order one, with their size depending on how ϵ ultimately becomes related to the fill fraction f [formally, $\epsilon = O(f^{\frac{1}{3}})$].

169 Note that the model in Eqs. (32) and (33) also applies in a different limit of the coating problem,
 170 corresponding to $f \ll 1$ and $\theta = O(1)$. In this case, the film coats the cylinder more uniformly, and
 171 a different asymptotic reduction leads to Eqs. (32) and (33), but with the sole modification that one
 172 must drop the free-surface gravity term, $f\eta_\theta \cos \theta$, in comparison to the other terms, $\sin \theta - S\eta_{\theta\theta}$,
 173 in Γ . This alternative reduction corresponds to the Herschel-Bulkley generalization of the leading-
 174 order lubrication model considered by Moffat and others (as accomplished for a Bingham fluid by
 175 Ross *et al.* [41]).

176 Such considerations lead us to now relax any constraints on the parameters in Eqs. (32) and (33),
 177 allowing us to consider both the pool and residual film, or coatings in which the lengthscale \mathcal{L} is
 178 larger, and indeed comparable to R . In other words, we use our model to bridge between the limit
 179 considered by Moffat and others, and the one underlying our asymptotic reduction (cf. Ref. [44]).

180 III. MODEL SOLUTIONS

181 We solve the evolution equation (32) as an initial-value problem, choosing a fixed grid in angle
 182 and employing spectral differentiation matrices to evaluate spatial derivatives [45]. This reduces
 183 Eq. (32) to a set of coupled ordinary differential equations (ODEs), which we integrate using
 184 Matlab's stiff integrator ODE15s. To ease the computation and take into account any unyielded
 185 plugs, we regularize the constitutive law by replacing the second relation in Eq. (34) by

$$\mathcal{Y}_\varepsilon = \eta - \frac{\text{Bi}}{|\Gamma|}, \quad \mathcal{Y} = \frac{1}{2}(\mathcal{Y}_\varepsilon + \sqrt{\varepsilon^2 + \mathcal{Y}_\varepsilon^2}), \quad (37)$$

186 where ε is a regularization parameter. Practically we take ε to be 10^{-4} or smaller, and verify that
 187 its precise value has little effect on the solutions other than smoothing out some of the finer scale
 188 features near the yield points.

189 As found by Ashmore *et al.* in the Newtonian version of the problem, solutions rapidly converge
 190 to steady state whenever an appreciable pool forms near the bottom of the cylinder. These steady
 191 states can be found more directly by abandoning (32) and setting instead $q = Q(\eta, \Gamma) = \text{const}$.
 192 With surface tension, this corresponds to a differential eigenvalue problem with periodic boundary
 193 conditions and the constraint (33), which can be solved using Matlab's BVP4c solver [again
 194 regularizing via Eq. (37)]. The flux condition must be inverted to write Γ in terms of η and q as
 195 part of this solution strategy. The BVP4c steady solution solver has the advantage of exploiting an
 196 adaptive grid to resolve any finer features in the solution.

197 A. Sample solutions

198 Figure 2 displays sample steady solutions for a Bingham fluid ($n = 1$) at fixed $\Lambda = 10$, $f = 0.01$
 199 and $S = 10^{-4}$ with varying Bi. Plotted are the angular profiles of the free surface, $h(\theta)$, and fake
 200 yield surface, $\mathcal{Y}(\theta)$. At these parameter settings, surface tension has only a minor effect and a narrow
 201 pool forms near the bottom of the cylinder. The solutions feature prominent pseudoplugs over the
 202 pool and film, which thicken to become true plugs (where $\mathcal{Y} = 0$) over sections to the left of the pool
 203 and at the top of the cylinder. Over these true plugs, $\eta = q$. The pseudoplugs are shaded according
 204 to the local plug speed, and the yield points where $\mathcal{Y} \rightarrow 0$ are indicated. Insets show illustrations of
 205 how the pool would appear within the rotating drum. Increasing the yield stress from the Newtonian
 206 limit reduces the size of the pool and drags it up the side, whilst thickening the residual film coating
 207 the remainder of the cylinder; the pool disappears for $\text{Bi} \rightarrow \frac{1}{2}$. This trend, which is shown in more
 208 detail in Fig. 3 for a wider set of solutions, is similar to an increase in rotation rate in the Newtonian
 209 problem [1].

210 A sequence more like an experimental series is shown in Fig. 4, which shows solutions with fixed
 211 rheological parameters ($\text{Bi} = 0.158$ and $n = 0.38$), fill fraction ($f = 0.0838$) and surface tension
 212 ($S = 10^{-4}$), but varying rotation rate (Λ). The trends with increasing rotation rate (decreasing Λ)
 213 reproduce those for increasing yield stress in Fig. 2. The choices for f , n , and Bi in Fig. 4 actually

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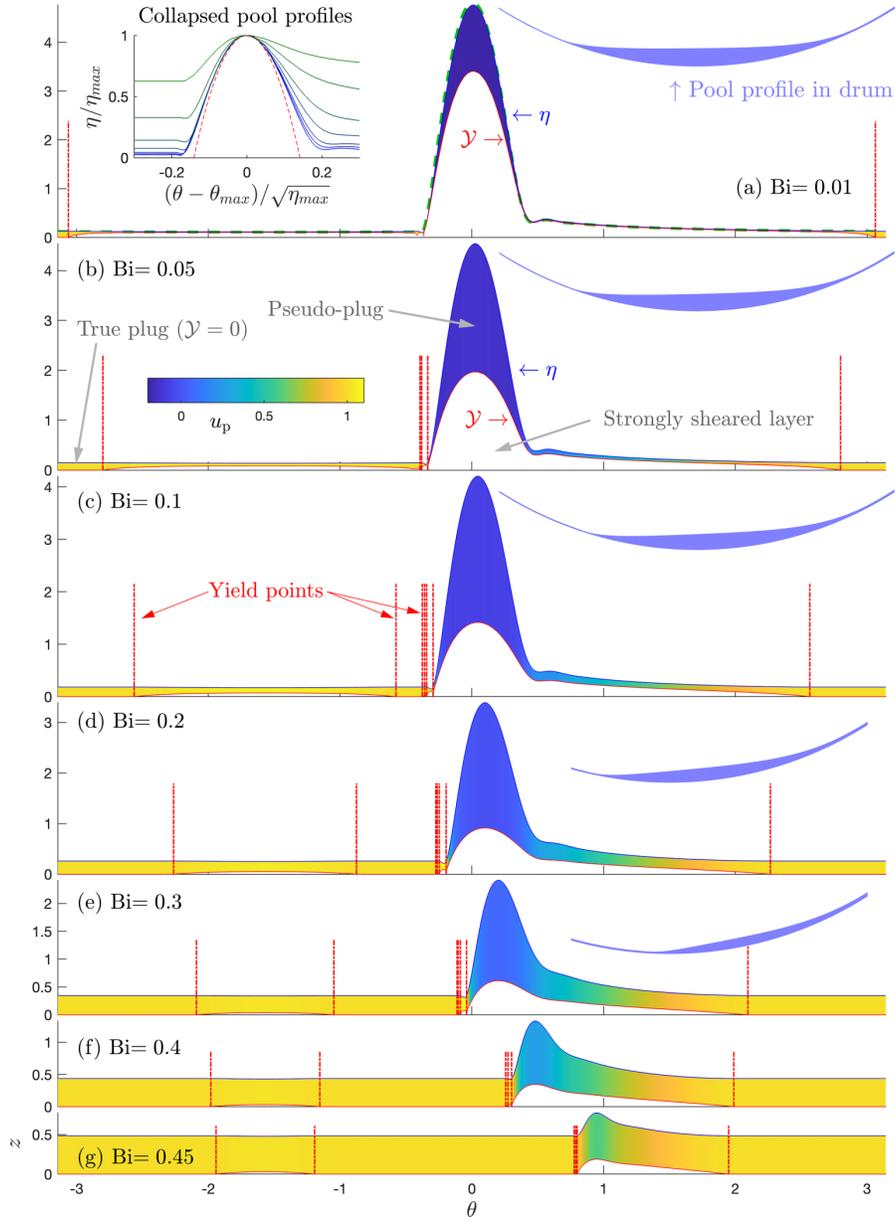


FIG. 2. Steady solutions with varying Bi (as indicated) for $(\Lambda, f, n, S) = (10, 0.01, 1, 10^{-4})$. Shown are the positions of the free surface η and yield surface \mathcal{Y} , with the pseudoplug and plug shaded according to the local value of plug speed, $u_p \equiv 1 - \frac{3}{2}\Lambda\Gamma\mathcal{Y}^2$. The vertical dot-dashed lines show the yield points. In panel (a), the dashed line shows the corresponding Newtonian solution, and the left-hand inset replots all the surface profiles (from blue to green) using the scaled variables η/η_{max} and $(\theta - \theta_{max})/\sqrt{\eta_{max}}$, and comparing with the approximation in Eq. (38) (red dashed line). The right-hand insets show the profiles of some of the solutions as they would appear inside the drum.

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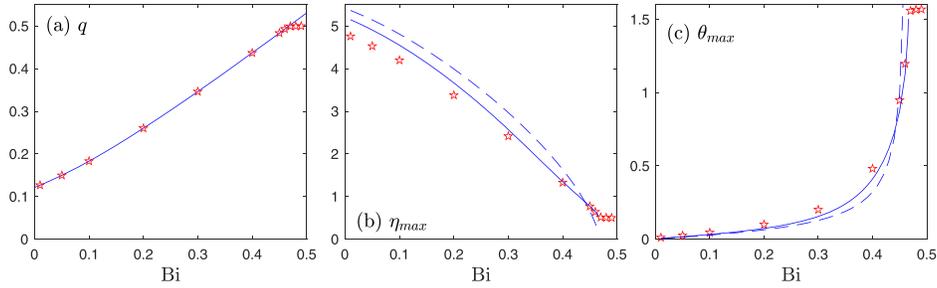


FIG. 3. (a) Flux, q , and (b,c) peak height and position, η_{max} and θ_{max} , plotted against Bi for $(\Lambda, f, S) = (10, 0.01, 10^{-4})$. The solid line in panel (a) shows the prediction from Eq. (40). The solid lines in panels (b) and (c) show the results from the analysis of the Appendix in which surface tension is neglected, and the dashed lines show the predictions from Eqs. (44) and (46).

214 correspond to the experiments discussed below in Sec. V. The choice for S , though, is too large by
 215 a factor of about four, assuming that the experimental fluid, Carbopol, has a surface tension slightly
 216 lower than that of water [13,23]. This is not significant, given the minor role played by surface
 217 tension in the solutions.

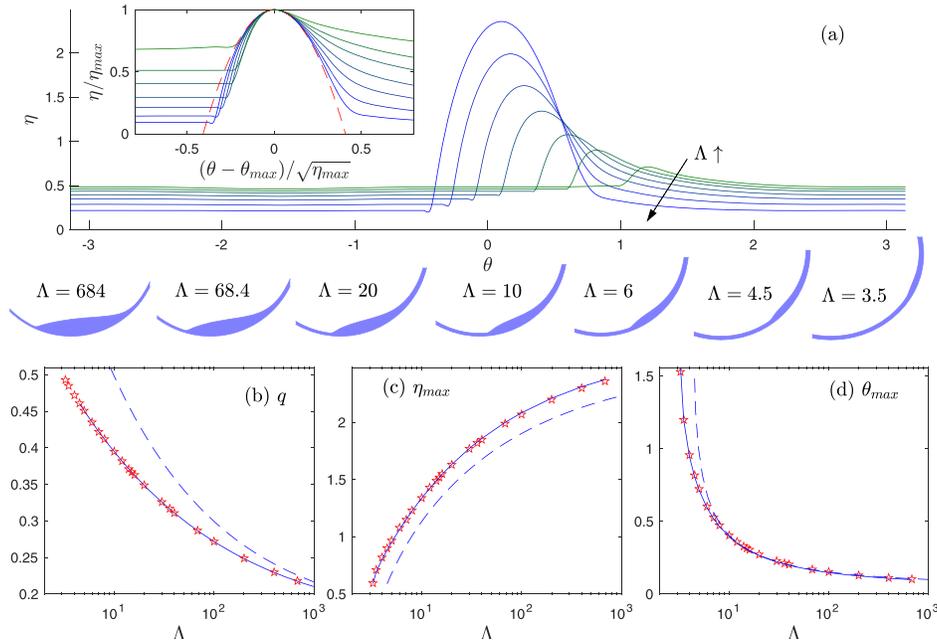


FIG. 4. Steady solutions with increasing Λ (from green to blue) for $(f, Bi, n, S) = (0.0838, 0.158, 0.38, 10^{-4})$. In panel (a) the free surface profile is shown for the values of Λ indicated. The upper left inset shows a collapse of these profiles over the pool, along with the prediction in Eq. (38). In panels (b–d), the flux q , and the peak height and position, η_{max} and θ_{max} , are plotted against Λ for a wider set of solutions (stars). The solid lines in panel (b) shows the predictions from Eq. (41), and the solid lines in panels (c, d) indicate the results from the $S = 0$ analysis of the Appendix. The dashed lines in panels (b, c, d) show the predictions from Eqs. (42), (44), and (46).

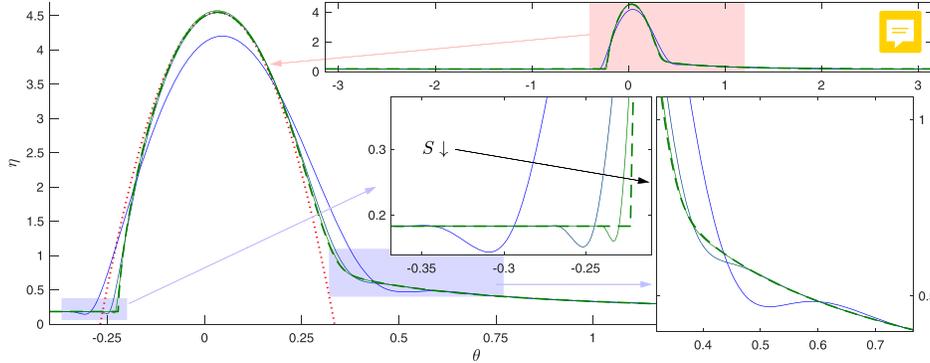


FIG. 5. Steady solutions with $S = 10^{-4}$, 10^{-5} , and 10^{-6} (solid, from blue to green), for $(\Lambda, \text{Bi}, n, f) = (10, 0.1, 1, 0.01)$. The dashed line shows the construction of the Appendix for $S = 0$. The dotted line shows the approximation (38). The main panel shows the pool solutions; the insets show the entire solutions and magnifications near the borders of the pool (as indicated by the shading and arrows).

To emphasize this point, we show solutions with varying S in Fig. 5, for a Bingham fluid with $(\Lambda, f, \text{Bi}) = (10, 0.01, 0.1)$. The solution with $S = 10^{-4}$ is still slightly sensitive to this surface tension parameter; once S is decreased to 10^{-5} or smaller, however, only some of the finer details of the solutions depend on S . This figure also includes a construction of the steady solution without surface tension altogether, as described more fully in the Appendix. The suite of solutions shown in Fig. 4, with the Herschel-Bulkley parameters chosen to match the experiments, are even less sensitive to S , as highlighted by the match of the peak depth and position between solutions computed for $S = 10^{-4}$ and $S = 0$ (unlike in Fig. 3, where there is a minor difference).

B. Anatomy of the steady, pooled solution

More details of a steady, pooled solution are displayed in Fig. 6. A similar picture, but for a solution in which the pool has been dragged well up the side of the cylinder, is shown in Fig. 7. The latter example corresponds to the viscoplastic generalization of one of the Newtonian solutions with shocks considered in previous literature (e.g., Refs. [1,10]). Note that the significantly pooled example in Fig. 6 possesses a substantial recirculation cell (as revealed by the streamline pattern also drawn), whereas the more shock-like case in Fig. 7 displays no recirculation.

As illustrated in Fig. 6, the pool is mostly relatively narrow and of high amplitude for $f \ll 1$ and $\Lambda \gg 1$, remaining localized to the bottom of the cylinder ($\theta \ll 1$). The flux relation $q = Q(\eta, \Gamma)$ then indicates that Γ must be $O(\Lambda^{-n})$ over this region. A local approximation of the first relation in Eq. (34) then implies that $f\eta_\theta + \theta \approx 0$, provided surface tension remains unimportant. Hence, $\eta \approx \eta_{\max} - \theta^2/(2f)$, which reflects how the surface of the pool becomes relatively flat, and can be cast as a more formal asymptotic solution relevant to the limit $\Lambda \gg 1$. The pool solutions in Figs. 2, 4, and 5 match fairly well with the approximation, even when the pool becomes shifted away from $\theta_{\max} = 0$, as long as we make the replacement $\theta^2 \rightarrow (\theta - \theta_{\max})^2$, i.e., the solution for the pool is approximately

$$\eta \approx \eta_{\max} - \frac{(\theta - \theta_{\max})^2}{2f}. \quad (38)$$

The pool therefore has a width of approximately $2\sqrt{2f\eta_{\max}}$, and contains about $\frac{4}{3}\eta_{\max}\sqrt{2f\eta_{\max}}$ of the total fluid area.

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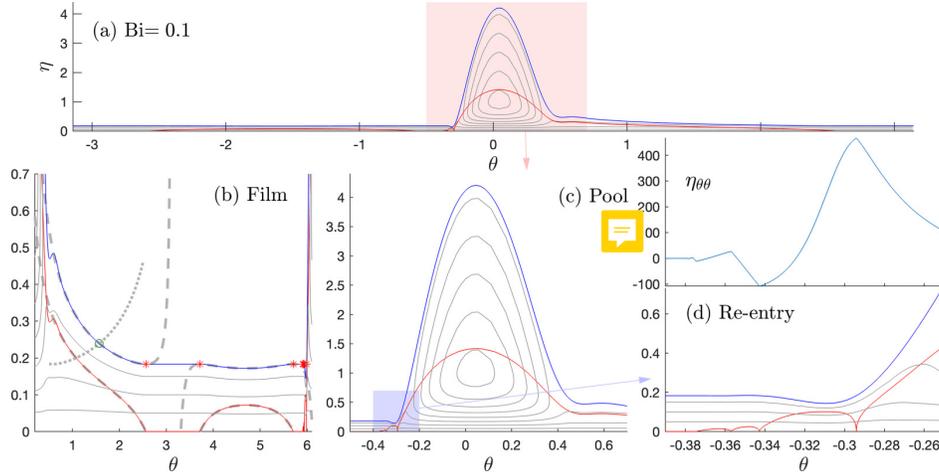


FIG. 6. Anatomy of the solution with $(\Lambda, f, Bi, n, S) = (10, 0.01, 0.1, 1, 10^{-4})$. Plotted are η (blue) and \mathcal{Y} (red), along with a selection of streamlines (light gray solid lines). In panel (b), the residual film is shown in more detail. The dashed gray lines show the root of the cubic $q = \text{constant}$ with $\Gamma = \sin \theta$, to which the solution converges over the film; the gray dotted line shows a second root, which intersects the first at the green circle at $\theta = \frac{1}{2}\pi$. A magnification of the pool is shown in panel (c). Details of where the film reenters the pool are plotted in panel (d), along with $\eta_{\theta\theta}$.

244 Outside the pool, over the residual film coating the remainder of the cylinder, we have $q =$
 245 $Q(\eta, \Gamma) = \text{const.}$ and $\Gamma \approx \sin \theta$, which for $n = 1$ boils down to a cubic equation for η :

$$\Lambda \sin \theta \left(\eta^3 - \frac{3}{2} b \eta^2 + \frac{1}{2} b^3 \right) - \eta + q = 0, \quad b = \frac{Bi}{|\sin \theta|}. \quad (39)$$

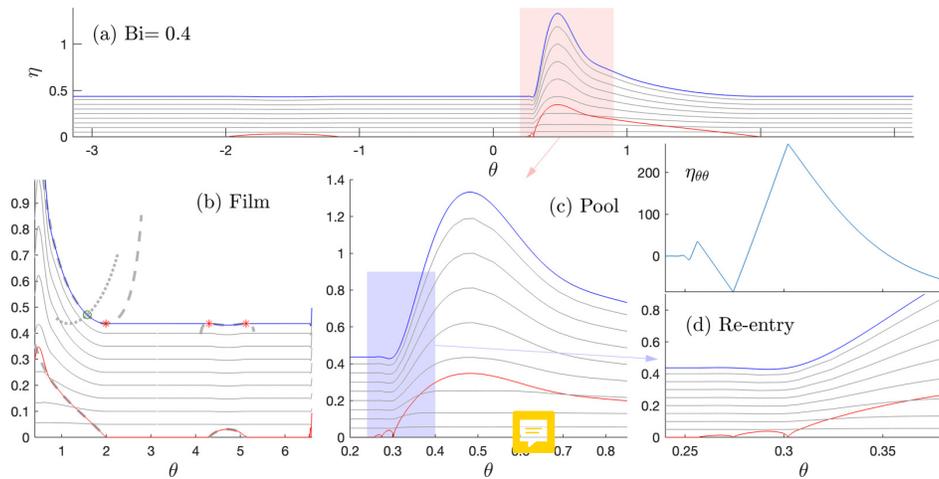


FIG. 7. A similar set of plots to Fig. 6, but for $Bi = 0.4$.

As shown in Figs. 6 and 7, the steady solutions to the full problem do indeed converge to one of the roots of the cubic over the film. Importantly, at $\theta = \frac{1}{2}\pi$ ($\Gamma \approx 1$), two of the roots intersect, to allow the film thickness to continue to decrease to the right of the pool (cf. Refs. [10,39]). The intersection (i.e., the presence of a double root) implies that

$$q = \frac{1}{4}\Lambda^{-\frac{1}{2}}\left[\left(\beta^2 + \frac{4}{3}\right)^{\frac{3}{2}} + 2\beta - \beta^3\right], \quad \beta = \Lambda^{\frac{1}{2}}\text{Bi}, \quad (40)$$

as illustrated in Fig. 3(a). For $\beta \ll 1$, we recover the Newtonian result $q \sim \frac{1}{4}\Lambda^{-\frac{1}{2}}\left(\frac{4}{3}\right)^{\frac{3}{2}}$. When $\beta \gg 1$, we enter instead a plastic limit (cf. Sec. III C) with $q \sim \text{Bi}$. When $n < 1$, similar predictions follow from demanding

$$Q(\eta, \sin \theta) - q = 0 \quad \text{and} \quad \frac{\partial}{\partial \eta}Q(\eta, \sin \theta) = 0, \quad (41)$$

although the results are not explicit and must be found by solving an algebraic problem numerically. As long as $\text{Bi}\Lambda^{\frac{n}{n+1}} \gg 1$, the second relation in Eq. (41) can be solved in the limit $\mathcal{Y} \ll 1$, to find $\mathcal{Y} \sim [(2+n^{-1})\Lambda\eta]^n \sim [(2+n^{-1})\Lambda\text{Bi}]^n$, and the first relation then implies that

$$q \approx \text{Bi} + \frac{1}{n+1}\left[\frac{n}{(2n+1)\text{Bi}\Lambda}\right]^n. \quad (42)$$

These results indicate that $\eta = O(q)$ over the film, which is small when there is a significant pool near the bottom of the cylinder. Consequently, the fluid is mostly contained in the pool, and so

$$\frac{4}{3}\eta_{\max}\sqrt{2f\eta_{\max}} \sim \pi \quad \text{or} \quad \eta_{\max} \sim \left(\frac{9\pi^2}{32f}\right)^{\frac{1}{3}}, \quad (43)$$

which is independent of all parameters but f . The pool solutions plotted in Figs. 2–7 have not yet reached the limit in Eq. (43), and still depend on Bi and Λ . To advance past the limitations of Eq. (43), we keep track of the contribution of the residual film to the fluid area. In particular, by taking $\eta \approx q$ over the film, and retaining its contribution to the constraint (33) [i.e., adding $2\pi q$ to the left of the first relation in Eq. (43)], we arrive at

$$\eta_{\max} \sim \left[\frac{9\pi^2(1-2q)^2}{32f}\right]^{\frac{1}{3}}. \quad (44)$$

Note also that the peak position, where $\Gamma = \sin \theta_{\max}$ if $|S\eta_{\theta\theta\theta}| \ll 1$, follows from solving

$$q = Q(\eta_{\max}, \sin \theta_{\max}), \quad (45)$$

or Eq. (39) if $n = 1$. Approximately (when \mathcal{Y} is small and $\eta_{\max} \gg q$), we have

$$\sin \theta_{\max} \sim \frac{\text{Bi}}{\eta_{\max}} + \left[\frac{(n+1)\text{Bi}}{(2n+1)\Lambda\eta_{\max}^{\frac{2+\frac{1}{n}}}}\right]^{\frac{n}{n+1}}. \quad (46)$$

The predictions from Eqs. (40), (44), and (46) are included in Fig. 3, and those from Eqs. (42), (44), and (46) in Fig. 4. The predictions for the flux in Eq. (40) or Eq. (41) follow from discarding the $O(f)$ and surface tension terms from Eq. (34), a relatively accurate approximation as seen in Fig. 3(c). By contrast, the cruder approximations leading to Eqs. (42), (44), and (46) introduce noticeable discrepancies with the numerical solutions in Figs. 3(b) and 3(c) and 4(b)–4(d).

Finally we comment on the structure of the transition layers between the pool and residual film. In Figs. 2 and 4–7, the layer to the right contains mild undulations owing to the importance of surface tension; by analogy with the Newtonian problem, the asymptotic structure of these standing capillary waves is likely complicated [10], although the impact of the yield stress is not known. Nevertheless, these undulations disappear once the surface tension parameter S becomes sufficiently small (cf. the Appendix), leaving a relatively smooth transition. The layer to the left of the pool is

276 more complicated: here the solutions with surface tension ($S > 0$) again pass through a wavetrain. In
 277 the viscoplastic problem, this wavetrain takes a distinctive form, with the curvature $\eta_{\theta\theta}$ converging
 278 to decaying sawtooth wave (see Figs. 6 and 7), as explored further in a related problem by Jalaal
 279 *et al.* [26,27]. The convergence to the limit $S \rightarrow 0$ is also further obscured by the presence of
 280 discontinuities in surface slope for $S = 0$ (see Fig. 5 and the Appendix), which must be smoothed
 281 on a sufficiently small scale by surface tension whenever S is finite.

282 **C. Plastic limit**

283 The solutions presented above indicate that the most significant effect of the yield stress arises
 284 at the slowest rotation rates ($\Lambda \gg 1$): as evident from Fig. 4 and Eq. (46), the viscoplastic fluid
 285 coating converges to a shape with finite slope in this limit. This feature contrasts sharply with the
 286 Newtonian problem, for which the free surface must become horizontal when surface tension is not
 287 significant.

288 For fluid with a yield stress, $\Lambda \rightarrow \infty$ actually corresponds to the “perfectly plastic” limit,
 289 wherein the yield stress dominates the viscous stress except over a thin boundary layer adjacent
 290 to the cylinder. That is, $\mathcal{Y} \rightarrow 0$, with the pseudoplug spanning most of the fluid depth and a good
 291 fraction of the residual film fully plugged up. Because $\mathcal{Y} \rightarrow 0$, the yielded regions of the limiting
 292 state satisfy $\eta \approx \text{Bi}/|\Gamma|$, or the ODE,

$$\eta_{\theta} \approx \frac{\text{Bi} - \eta \sin \theta}{f \eta \cos \theta}, \quad (47)$$

293 when surface tension is neglected. This ODE can be integrated from a yield point $\theta = \theta_L$ at the left
 294 border of the pool to the right-hand yield point $\theta = \theta_R$, which approaches $\theta = \frac{1}{2}\pi$ in the plastic
 295 limit. The slope is discontinuous at θ_L , but smoother at θ_R (cf. Figs. 2, 6, and 7). The position of
 296 the left-hand yield point θ_L must be selected to satisfy the constraint (33) (practically, this can be
 297 accomplished by adopting a trial value for θ_L , computing the solution and the implied fluid area,
 298 then performing Newton iteration until the constraint is satisfied). Sample solutions constructed in
 299 this fashion are illustrated in Fig. 8. With surface tension, the construction is more complicated in
 300 view of the behavior near the left-hand yield point [cf. Fig. 6(d)].

301 For $f \ll 1$, the pool is approximately given by Eq. (38) [as follows from neglecting Bi on the
 302 right-hand side of Eq. (47) and taking $\theta \ll 1$], whereas the film is given by $\eta = \text{Bi}/\sin \theta$ [dropping
 303 the left-hand side of Eq. (47)]. The pool then has the maximum depth in Eq. (43), or more accurately
 304 Eq. (44) with $q = \text{Bi}$, and an almost flat surface with a finite inclination given by $\sin \theta_{\max} = \text{Bi}/\eta_{\max}$
 305 [see Figs. 8(b) and 8(c)]. These results can be seen much more readily by first observing that,
 306 given the fill fraction f , the maximum depth of the segment of the drum occupied by the pool is
 307 approximately Eq. (43), or Eq. (44) if the residual film is taken into account. Then, per unit width,
 308 the net gravitational force on the central cross-section of the segment (through $\theta = \theta_{\max}$) has an
 309 angular component $\rho g \hat{h}_{\max} \sin \theta_{\max}$ that must be largely balanced by the yield stress, τ_y . That is,

$$\sin \theta_{\max} = \frac{\tau_y}{\rho g \hat{h}_{\max}} \equiv \frac{\text{Bi}}{\eta_{\max}}. \quad (48)$$

310 This result mirrors that for the onset of flow in a uniform sheet down an incline, which has been
 311 suggested as a practical rheometer [19]. Similarly, one could envision using a slowly rotating drum
 312 to infer yield stresses.

313 The second term in Eq. (42) illustrates how corrections to q arise at $O(\Lambda^{-n})$. As can be seen
 314 from combining Eqs. (42) and (44), this suggests that η_{\max} and $\theta_{\max} \sim \sin^{-1}(\text{Bi}/\eta_{\max})$ also contain
 315 corrections terms of that order. Thus, η_{\max} and θ_{\max} should converge to the plastic limit at the
 316 rate Λ^{-n} , as appears to be borne out in Figs. 8(c) and 8(d), in which these quantities are plotted
 317 against Λ^{-n} . However, as evident from the second term on the right of Eq. (46), there are also
 318 potentially larger corrections of $O(\Lambda^{-\frac{n}{n+1}})$. In fact, these corrections, which do ultimately control
 319 the convergence to the plastic limit, are hidden at relatively small values of Λ^{-n} in Figs. 8(c) and

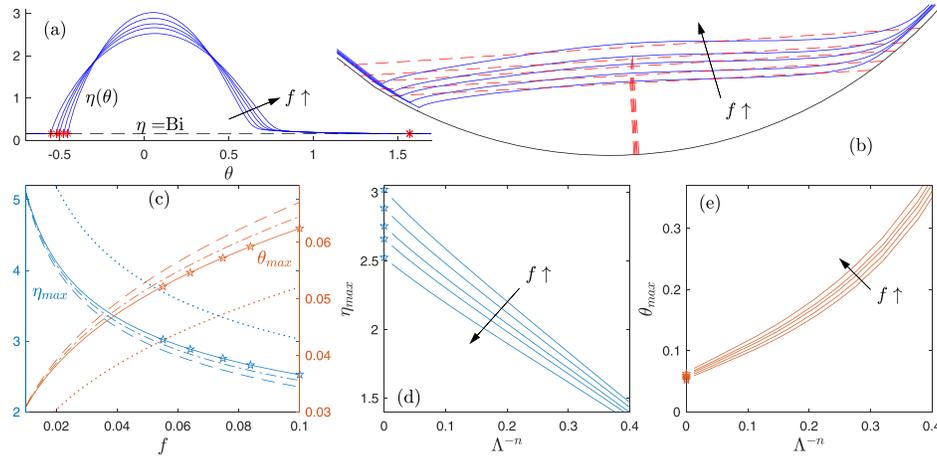


FIG. 8. (a) Solutions in the plastic limit for $f = 0.055, 0.064, 0.0747, 0.0838,$ and 0.1 , with $S = 0$ and $(n, \text{Bi}) = (0.38, 0.158)$; the stars indicate the yield points. The solutions are replotted as they would appear in the drum in panel (b), with the dashed lines indicating the approximation of the pool as a segment given by Eqs. (A2) and (48). (c) Peak depth η_{\max} and its position θ_{\max} plotted against f . The dotted lines show Eqs. (43) and (48); the improvements in Eqs. (44) and (A2) are shown by dashed and dot-dashed lines, respectively. The stars indicate the values of f shown in panel (a). (d, e) Peak depth η_{\max} and its position θ_{\max} plotted against Λ^{-n} for the same values of f as in panel (a). The limits for $\Lambda \rightarrow \infty$, from panel (b), are indicated by stars.

8(d). Instead, a conspiracy between the parameter settings ensures that the intermediate rate of convergence Λ^{-n} , or Ω^n in terms of the dimensional rotation rate, is more obvious.

IV. DRAINAGE SOLUTIONS

If the flux function $Q(\eta, \Gamma)$ in Eq. (32) is modified to delete the first term, the background rotation of the cylinder is removed (the scalings of the problem then employ an arbitrary rate Ω , which could be removed by another scaling of time; we avoid this extra rescaling here). The evolution equation then models the gravitationally driven drainage of the fluid coating. For a viscous fluid, the outcome of such drainage is a flat pool at the bottom of the cylinder with an ever-thinning residual film elsewhere. When there is a yield stress, however, drainage ends with pools with finite free surface slopes, analogous to the slumped states resulting from dambreaks on flat surfaces [4,5,30,32,35], together with a residual film of finite thickness.

A. Viscous (and power-law) drainage

Solutions for viscous drainage [$(n, \text{Bi}) = (1, 0)$] are shown in Fig. 9. These initial-value problems commence with either a uniform film [$\eta(\theta, 0) = \frac{1}{2}$] or a state corresponding to a steadily rotating solution (with $\Lambda = 1$), simulating the sudden stop of the cylinder. The latter is left-right asymmetrical at the outset, but the solution quickly symmetrizes and the long-time behavior of both solutions is similar. In particular, a slowly growing pool forms the bottom of the cylinder, as a residual film slowly drains downwards.

The long-time drainage solutions can be described analytically, as follows. First, over the pool, $\Gamma \sim \sin \theta + f \eta_\theta \cos \theta \sim 0$, with the approximate solution in Eq. (38). Second, over the film in the right-half of the cylinder ($0 < \theta < \pi$ and $\sin \theta > 0$), we have $\Gamma \sim \sin \theta$ and the evolution equation

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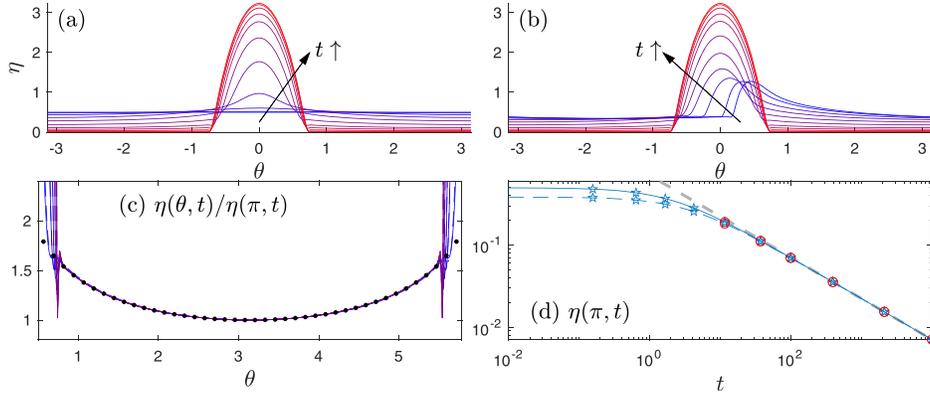


FIG. 9. Two viscous drainage solutions for $f = 0.0838$ and $S = 10^{-4}$ ($\Lambda = 1$). The initial conditions are (a) $\eta(\theta, 0) = \frac{1}{2}$ and (b) a steady rotating solution with $\Lambda = 1$. The residual film is scaled by $\eta(\pi, t)$ in panel (c), and the time series of $\eta(\pi, t)$ is plotted in panel (d). The times of the snapshots in panels (a, b), which increase from blue to red, are plotted as dots in panel (d); the circles show the slightly fewer snapshots plotted in panel (c). The dots in panel (c) and gray dashed line in panel (d) show the prediction from Eq. (50) with $n = 1$.

341 for $\text{Bi} = 0$ reduces to

$$\eta_t - \Lambda[\eta^{2+\frac{1}{n}}(\sin \theta)^{\frac{1}{n}}]_{\theta} \sim 0 \quad (49)$$

342 [the long-time solution is symmetric, with $\eta(\theta, t) = \eta(-\theta, t)$]. We note the separable solution,

$$\eta \sim \left[\frac{n}{(2n+1)\Lambda t} (\sin \theta)^{-\frac{n+1}{2n+1}} \int_{\theta}^{\pi} (\sin \theta)^{-\frac{1}{2n+1}} d\theta \right]^{\frac{n}{n+1}}, \quad (50)$$

343 which is compared to late-time numerical solutions in Fig. 9(c) for $n = 1$. The viscous film therefore
 344 drains as $t^{-\frac{1}{2}}$, as seen in Fig. 9(d). For $n < 1$, and the fluid becomes more shear thinning, the film
 345 drains more slowly as $t^{-n/(n+1)}$.

346 B. Drainage with a yield stress

347 Sample drainage solutions with a yield stress are displayed in Fig. 10. In the first of these, a flat
 348 pool at the bottom of the cylinder is suddenly rotated clockwise by 90° , and then allowed to slump
 349 to rest. The slump leaves behind a thin residual layer held at the yield stress (with $\eta \approx \text{Bi}/|\sin \theta|$),
 350 spanning the location of the original pool. To generate a second solution, the final state from this
 351 slump is then rotated again by 90° , and left to slump again. Repeating this exercise two more times
 352 then gives four rotate-and-slump solutions, the final states of which are shown in Fig. 10(b). As the
 353 slumps are sequentially rotated, more fluid is left coating the cylinder and the area of the pool slowly
 354 decreases. For the final rotate-and-slump, the slumping fluid encounters the residual film left by the
 355 first slump on its way to the bottom of the cylinder.

356 The rotate-and-slump scenario leads to fluid slumping in one direction and an asymmetrical
 357 final deposit. Other examples, with less asymmetry are shown in Figs. 10(c) and 10(d). In these
 358 two cases, the initial conditions correspond to steady solutions for a rotating drum at two different
 359 rotation rates, each simulating a sudden stop. Because these solutions feature both an offset pool and
 360 a residual film thickened by rotation, when the fluid begins to drain, flow occurs down both sides
 361 of the cylinder. For the slower example in Fig. 10(c), the drainage down the left side of the cylinder

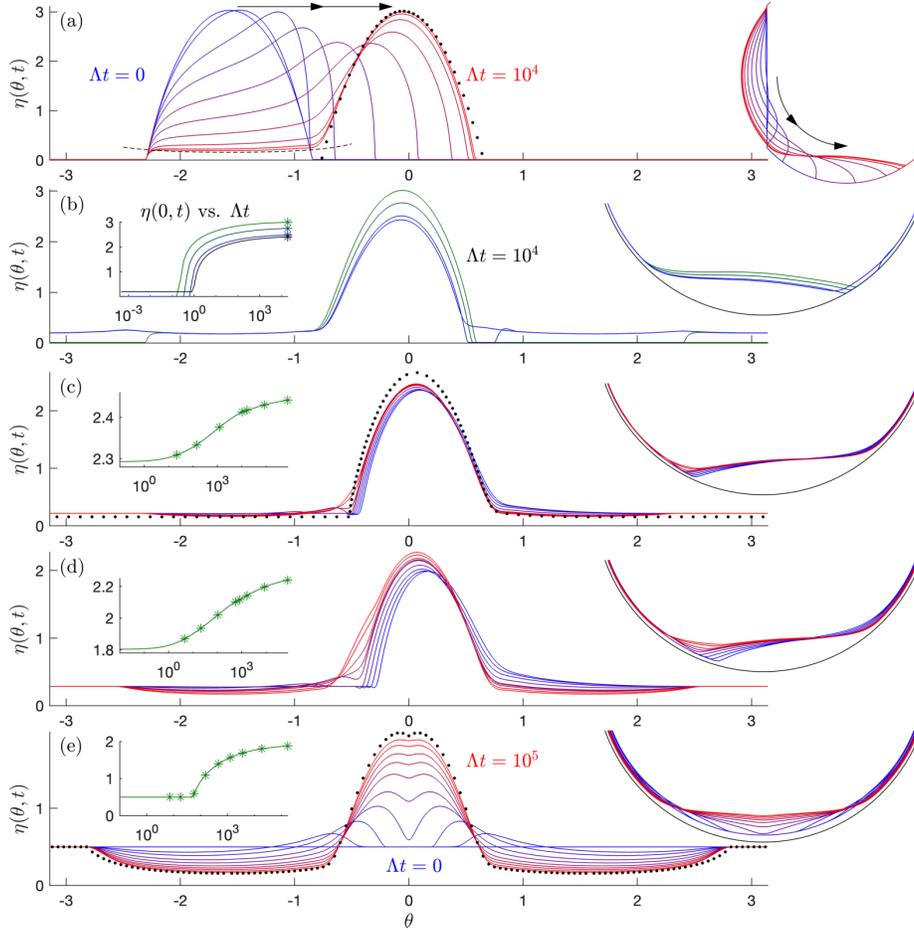


FIG. 10. Drainage solutions for $(f, \text{Bi}, n, S) = (0.0838, 0.158, 0.38, 10^{-4})$. In panel (a), a flat pool at the bottom of the cylinder is suddenly rotated clockwise by 90 degrees; shown are snapshots of $\eta(\theta, t)$ as the fluid slumps back to rest (from blue to red). The black dashed line shows $\eta \sim \text{Bi}/|\sin \theta|$, and the dots plot (38). In panel (b), the final state is plotted, along with three others generated by successive rotate-and-slump events (see main text, with the progression plotted from green to blue). In panels (c, d), slumps from steady states in rotating drums are used (respectively, $\Lambda = 684$ and 68.4), simulating sudden stops. For panel (e), the slump occurs from a uniform coating [$\eta(\theta, 0) = \frac{1}{2}$]. The right-hand insets show the solutions as they would appear in the cylinder, and the insets on the left show the time series of $\eta(0, t)$, with the stars showing the times of the snapshots in the main panels. For the solutions in panels (a) and (b), to avoid any issues with contact lines, a thin prewetted film of thickness 10^{-2} is included to ensure that $\eta(\theta, 0) > 0$ everywhere. The black dots in panels (c) and (e) plot the results predicted in the plastic limit for $S = 0$ [i.e., from integrating Eq. (47)] subject to either $\eta = \text{Bi}$ at the left of the pool and $\theta = \frac{1}{2}\pi$ (c), or symmetry about $\theta = 0$ (e).

is mild and the final deposit is almost as asymmetric as those from the rotate-and-slump examples. More drainage arises on the left for the faster example in Fig. 10(d).

When all the drainage takes place down one side, the final state is the same as the plastic limit ($\Lambda \rightarrow \infty$) of the steadily rotating solutions of Sec. III C, except for any differences arising over the

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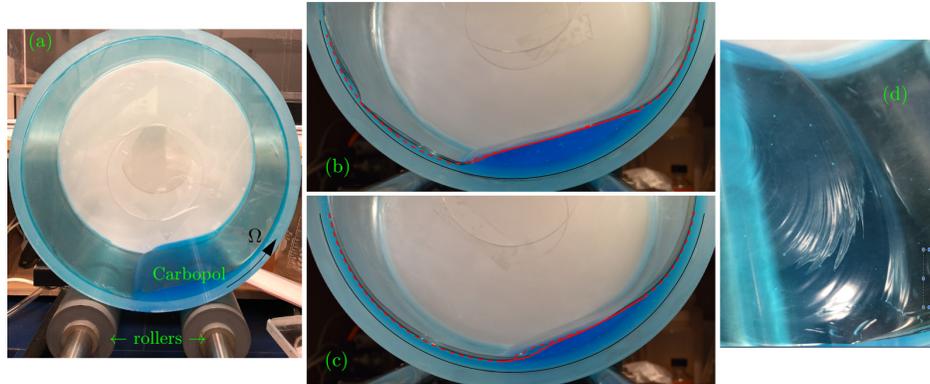


FIG. 11. Images of the rotating drum, with fits to the surface profile of an enclosed Carbopol suspension, dyed blue with food colouring. In panels (b, c), with rotation rates are $\Omega = 0.024$ rad/s and 0.14 rad/s, respectively, two fits are shown. The solid line shows a fit to the profile along the front face; the dashed line is a fit with a weaker threshold that captures more of the profile further inside the drum as well as the residual film outside the pool. The two fits are mostly the same in panel (b), but at the faster rotation speed in panel (c), the three-dimensionality of the front edge of the pool becomes more pronounced; the second fit takes some account of this to furnish a more averaged profile. An oblique view of a three-dimensional surface profile is shown in panel (d), along with some surface wrinkles generated by the flow dynamics. $f = 0.0838$.

366 plugged sections of the residual film. This correspondence arises because, for final solution from the
 367 initial-value problem to approach rest, the yield condition $\mathcal{Y} \rightarrow 0$ must be reached over the flowing
 368 regions, implying Eq. (47). The limiting, steadily rotating solution for $\Lambda \rightarrow \infty$ is also plotted in
 369 Fig. 10(c). Despite the common form expected for the solutions over the yielded sections of the
 370 film, differences can remain over the plugs, since the rotating solution has $\eta = \text{Bi}$ here, whereas the
 371 drainage solution may have $\eta < \text{Bi}$.

372 Finally, the last example of Fig. 10(e) shows a slump from a uniform initial coating [$\eta(\theta, 0) = \frac{1}{2}$].
 373 This example is necessarily left-right symmetrical. The maximum depth also now no longer arises at
 374 the center of the pool, where, instead, a mild minimum arises at the meeting of the counter-flowing
 375 drainage currents. The final state in this case is more directly computed by neglecting surface tension
 376 and integrating the ODE (47) from a trial value of η at $\theta = 0$, to the right and up to to $\theta = \frac{1}{2}\pi$. As
 377 the integration of Eq. (47) proceeds past the pool, the solution converges to the residual film, with
 378 $\eta \approx \text{Bi}/\sin \theta$, which can then be used to continue the solution past the singular point $\theta = \frac{1}{2}\pi$ (as in
 379 the construction of the Appendix) and up to the plug at $\theta \approx \sin^{-1}(2\text{Bi})$, where $\eta = \frac{1}{2}$. The solution
 380 for $\theta < 0$ is then given by reflection. This computation can be iterated again, adjusting the value of
 381 η at $\theta = 0$, such that the constraint (33) is eventually satisfied. The dots in Fig. 10(e) show the result
 382 of this construction. There is a minor difference between the construction and the final snapshot of
 383 the initial-value computation because surface tension plays a slight role near the center of the pool
 384 for the latter (cf. Fig. 5), smoothing out the corner at the local minimum.

385 **V. EXPERIMENTS**

386 To complement the theoretical analysis, we conducted some experiments using an aqueous
 387 suspension of Carbopol in an acrylic cylinder; see Fig. 11. The Carbopol (Ultrez 21) suspension was
 388 neutralized with sodium hydroxide and had a concentration of about 0.2% and density $\rho \approx 1$ g/cm³.
 389 A Herschel-Bulkley fit of the flow curves measured in a rheometer (Kinexus, Malvern, fitted with

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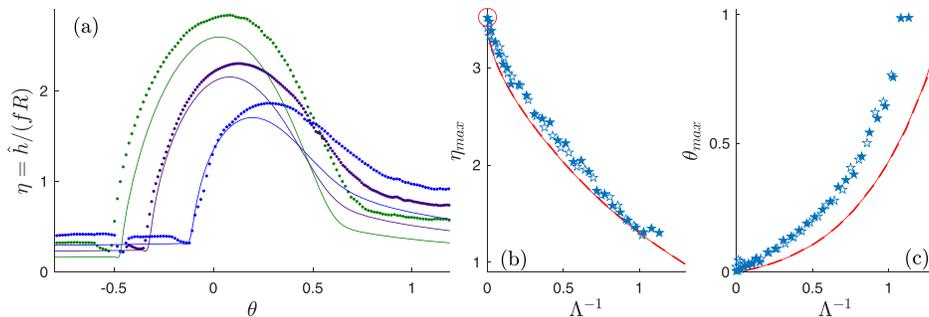


FIG. 12. (a) Experimental syrup profiles at the dimensionless rotation rates, $\Lambda^{-1} = 3\mu\Omega/(\rho g R f^2) = (0.21, 0.43, 0.67)$ (from green to blue), with $f = 0.0687$. (b) Maximum depths \hat{h}_{max} and (c) angular positions θ_{max} plotted against Λ^{-1} . The open stars correspond to the drum with inner radius and width, $(R, W) = (27.8, 13.1)$ cm; the filled stars correspond to a wider drum with $(R, W) = (28.8, 20)$ cm. The solid and dashed lines in panels (b, c) show equivalent theoretical results with $S = 10^{-5}$ and $S = 0$, respectively. The red circle in panel (b) shows the depth expected for $\Omega = 0$. The fits of the surface in the experiments are all to the syrup profile as observed at the front face of the drum.

roughened parallel plates) gave the parameters,

$$\tau_v = 18 \text{ Pa}, \quad K = 13 \text{ Pa s}^n, \quad n = 0.38. \quad (51)$$

The inside surface of the cylinder was roughened with 60 grit sandpaper in an effort to eliminate the effective slip that Carbopol suffers against acrylic surfaces [8,18,20]. The front and back surfaces were left smooth, however, to allow the Carbopol to slide more freely over them, and try to maintain a more two-dimensional flow pattern. The cylinder, with inside diameter of 27.8cm and a length of 13.1cm, was placed on friction rollers attached to a stepper motor to drive a prescribed rotation.

To conduct a set of experiments, fluid was first loaded into the drum up to a given fill fraction. The drum then sealed, to minimize evaporation, and its rotation rate was first increased and then decreased in a step-like manner to sweep over a range of Ω . At each rotation rate, the free surface was imaged through the clear, front, acrylic face. Figure 11 shows two examples illustrating our fit of the surface profile using image processing (based on a threshold intensity for the difference between the blue and red channels of the RGB image). Whenever the fluid drained into a pool near the bottom of the cylinder, a steady state was reached. Several images were taken of such steady states and the extracted profiles combined to furnish an average. Except at the lowest rotation rates, these averages were taken over at least two rotations of the drum. For the lowest rotation rates, each average spanned more than thirty minutes, and the sweep in rotation rate performed over the course of hours.

As a comparison against a suite of control experiments with a Newtonian fluid, we also conducted tests in which the cylinder was filled with golden syrup (density $\rho = 1.54 \text{ g/cm}^3$, viscosity 11 Pa s) to a fill fraction of $f = 0.0737$. Results for these controls are shown in Fig. 12. Displayed are a selection of measured depth profiles, and the peak depth and its position, plotted against rotation rate. The figure includes corresponding results from the theoretical model (for $n = 1$ and $\text{Bi} = 0$, adopting either $S = 0$ or $S = 10^{-5}$, the latter corresponding to a dimensional surface tension parameter of 0.5 Nm^{-1} ; the overlap of the two theoretical lines highlights how the results are insensitive to surface tension effects). There is some agreement between theory and experiment, although the location of the pool is somewhat shifted between the two, particularly at higher rotation rate. Better agreement is reported by Tirumkudulu and Acrivos [44] who used a different drum geometry and fluid, but similar fill fractions. In particular their drum was rather longer than ours (in comparison to drum radius). It is therefore possible that the three-dimensionality of the flow

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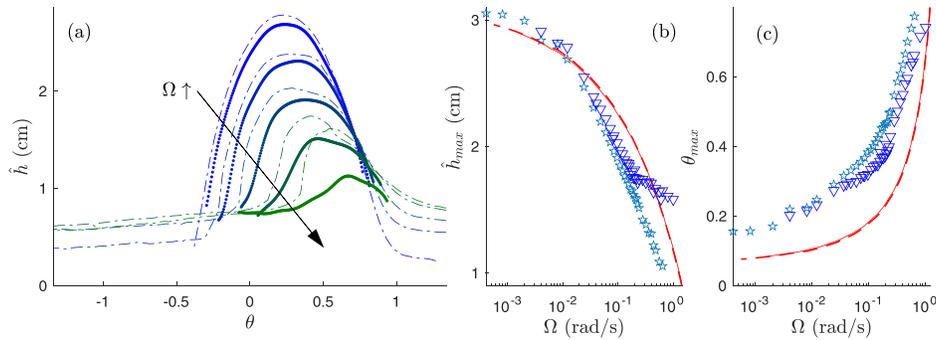


FIG. 13. (a) Experimental Carbopol profiles for $f = 0.0838$ at the rotation rates, $\Omega = (0.012, 0.036, 0.097, 0.24, 0.44)$ rad/s, corresponding to $\Lambda = (457, 152, 57, 23, 12)$ (from blue to green) and $(n, Bi) = (0.38, 0.158)$. The dots show the profiles fitted for the front face; the dash-dotted lines show the fits including more of the interior (see Fig. 11). (b) Maximum depths \hat{h}_{\max} and (c) angular positions θ_{\max} plotted against rotation rate Ω . The solid and dashed lines in panels (b, c) translate the theoretical results (with $S = 10^{-4}$ and $S = 0$, respectively) from Fig. 4; the stars and triangles show the two fits (front face and including interior, respectively).

419 field is responsible for the discrepancy in our experiments. In fact, there were noticeable residual
 420 films of syrup coating annuli on the front and back faces of the cylinder, the length of the pool was
 421 comparable to the drum width, and the shape of the pool became visibly more three dimensional in
 422 the experiments as the rotation rate increased (cf. Fig. 11). Nevertheless, repeating the experiments
 423 in a longer drum did not substantially change the results for the peak depth and its position; see
 424 Fig. 12.

425 Figure 13(a) shows a sequence of surface profiles for increasing rotation rate for experiments with
 426 Carbopol at fill fraction $f = 0.0838$. These tests correspond to a set of values for the dimensionless
 427 parameter Λ similar to those in Fig. 4. Also plotted are the (dimensional) maximum depth, \hat{h}_{\max} ,
 428 and its angular location, θ_{\max} , against Ω for a much wider set of experiments. Again, there is some
 429 qualitative agreement between the experimental results in Fig. 13 and those predicted theoretically,
 430 but the two disagree quantitatively. As in the Newtonian case, this disappointment may arise partly
 431 because of the three-dimensionality of the experimental flows. The effect of the third dimension is
 432 illustrated by the comparison between the profiles extracted at the front face of the drum, and those
 433 found when the thresholding in the image analysis is weakened to capture more of the profile further
 434 inside the drum as well as the residual film outside the pool (Fig. 11). At lower rotation rates, the
 435 two profiles largely match, but they disagree significantly for $\Omega > 0.1$ rad/s.

436 The discrepancy between theory and experiment may also arise because the fill fraction in the
 437 experiments, $f = 0.0838$, corresponds to $\epsilon \approx 0.3$ in the asymptotic theory. The relatively large
 438 size of ϵ suggests that the theory may not be very accurate. Indeed, the surface slopes observed
 439 experimentally were fairly steep, potentially invalidating the shallow-flow theory. Imaging problems
 440 (the need to view the flow through the sidewall and the local meniscus there) and the incomplete
 441 elimination of wall slip may also both contribute.

442 Experimental results closer to the plastic limit ($\Omega \rightarrow 0$) and for varying fill fraction are shown
 443 in Fig. 14. Here, following the theoretical predictions in Fig. 8, we plot the maximum depth and its
 444 angular position, $h_{\max} = \hat{h}_{\max}/R \equiv f\eta_{\max}$ and θ_{\max} against Ω^n . According to the theory, this power
 445 of rotation rate should straighten the data into straight lines. Figure 14 does indeed suggest this to be
 446 the case: the data is certainly straightened in comparison to a plot using Ω , although any exponent
 447 between 0.3 and 0.45 is as successful. Linear fits to the data plotted in this way can be used to
 448 estimate the limiting maximum depth and its position for $\Omega \rightarrow 0$, as also shown in the figure. The

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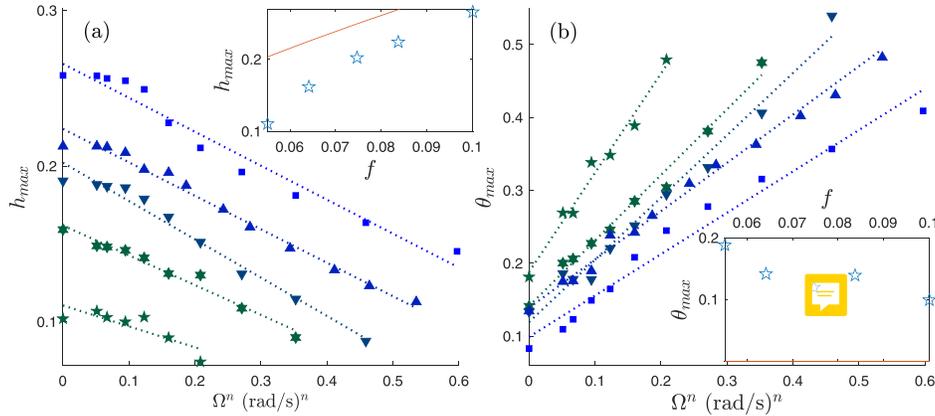


FIG. 14. Experimental results for Carbopol with varying fill fraction ($f = 0.055, 0.064, 0.0747, 0.0838$ and 0.1 , from green to blue), showing (a) maximum depth $h_{max} = \hat{h}_{max}/R \equiv f\eta_{max}$ and (b) position θ_{max} plotted against Ω^n (with Ω in rad/s). The dotted lines show linear fits, and the insets plot the implied intercepts (for $\Omega \rightarrow 0$) against f . The solid lines in the insets show Eqs. (44) and (48) (with $q = Bi$).

limit for the maximum depth compares somewhat favourably with the prediction in Eq. (44) (with $q = Bi$), but the experimental results for the position are again too high in comparison to Eq. (48).

Finally, Fig. 15 displays some experimental results for the drainage problem. As predicted theoretically, the final state reached after drainage supports finite surface slopes and has a varying degree of left-right symmetry depending on how the draining flow was initiated, and, in particular, how much of the fluid drained down each side. The experiments shown in Fig. 15 roughly correspond to the theoretical solutions plotted in Fig. 10. In particular, shown are a succession of rotate-and-slump events, or adjustments taking place when a rotating cylinder, with its fluid coating in steady state, is abruptly stopped. For the final example of the latter three tests, the cylinder was rotated at sufficiently high speed that the fluid coating became relatively uniform; the drainage that resulted after stopping the cylinder was then almost left-right symmetric.

VI. CONCLUSIONS

In this paper, we have considered the viscoplastic version of the problem of viscous rimming flow around a rotating drum. Our analysis extends some earlier work that considered relatively rapidly rotating coatings of either power-law or Bingham fluids [28,41]. In particular, we allowed the fluid to be described by the Herschel-Bulkley constitutive law and to rotate more slowly, so that the fluid mostly gathers into a pool near the bottom of the drum. The main effect of introducing a yield stress is to arrest and thicken the residual film coating the cylinder outside the pool, thereby drawing more fluid out of the pool and forcing that feature to ascend up the side of the drum. We complemented the theory with some experiments with a Carbopol suspension, finding qualitative, but not quantitative agreement. The discrepancies between theory and experiment may result from, in the experiments, three-dimensional flow fields (the theory is two-dimensional), the presence of wall slip [8,18,20], or overly steep surface slopes (we employ a shallow flow theory).

A distinctive feature of shallow viscoplastic fluids is the development of a superficial, plug-like flow. In most situations, this plug-like region cannot be rigid because there must also be weak straining motion along the film, leading to previous work referring to this region as a pseudoplug [2,6,47]. Such pseudoplugs also feature in the viscoplastic rimming flows considered here. Nevertheless, in some settings, the pseudoplug can itself fall below the yield stress, to lead to genuine floating plugs [3,22,31]. In the curved geometry of the shallow coating flow, it is more difficult to understand how

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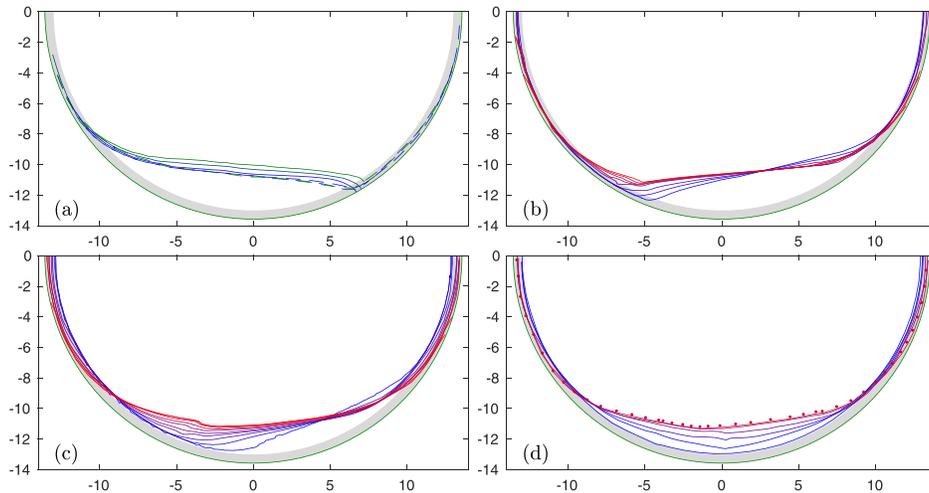


FIG. 15. Experimental drainage tests with Carbopol; $f = 0.0838$ and $(n, Bi) = (0.38, 0.158)$. In panel (a), a mound of material emplaced at the bottom of the cylinder was rotated clockwise by about 90° to force an adjustment to rest (green curve). The new mound was then rotated again twice (by a similar angle) to prompt two more adjustments (each leaving a longer residual coating behind; solid lines, from green to blue). The cylinder was then rotated three more times, with the fluid now encountering the coating from the previous rotate-and-slump events (dashed lines, from green to blue). In panels (b–d), steady-state cylinders rotating with rates of $\Omega = 0.008$ rad/s, 0.08 rad/s and 2 rad/s were abruptly stopped, leaving the fluid to drain (outlines from blue to red). In panel (d), the dots show the outcome of a repetition of the experiment. In all panels, the shaded region indicates the uniform film expected with a fill fraction of $f = 0.0838$, and the units are in cm.

478 the pseudoplug could become truly rigid in the present context, except near the angular positions
 479 where the speed of the pseudoplug becomes relatively flat near an extremum. Such positions do arise
 480 near the center of the pool, suggesting that this feature could partly plug up. In fact, it is possible
 481 to extend the analysis presented in Refs. [3,22,31,47] to explicitly build in the plugged parts of
 482 the pool. It is also conceivable that when the film is no longer shallow (at higher fill fractions and
 483 slow rotation rates), the entire pool plugs up. Indeed, in some experiments that we conducted with a
 484 thicker Carbopol suspension (not reported in the main text), pools with complicated surface features
 485 could be emplaced in the drum; these features survived the rotation of the drum, persisting for the
 486 length of the experiment.

487 Rotating drums are popular devices to study the flow of various materials, besides viscous
 488 fluids [42]. In particular, rotating drums have been used to help formulate constitutive theories
 489 for granular media [36]. At low rotation rates, granular flow in rotating drums becomes unsteady,
 490 continuous avalanching (or “rolling”) becoming replaced by episodic avalanching (“slumping”;
 491 e.g., Refs. [7,42]). Although this transition may be a unique feature of a granular medium, one of
 492 our interests in conducting the experiments of the present work was to explore whether nonideal
 493 viscoplastic effects near the yield stress [6] might prompt analogous or novel, unsteady flow
 494 dynamics at very low rotation rates. None were evident for the Carbopol suspension, suggesting
 495 that a fluid with more hysteretic material behavior might be needed in the drum [17,24,38]. At
 496 higher rotation rates, however, surface features with a relatively small spatial scale did become
 497 evident for the Carbopol [Fig. 11(d)]. These features resembled the wrinkling of elastic or viscous
 498 films [12,40,46] and may have some correspondence with surface patterns seen in other viscoplastic
 499 flows [34].

Finally, for Newtonian rimming flow, it has been of interest to move beyond the Stokes limit and add inertia to the thin film problem to explore the potential for hydrodynamic instability (e.g., Ref. [9]). One could also continue along this vein for viscoplastic rimming flow. Indeed, one expects that inertial viscoplastic surges should become unstable and time-dependent at sufficiently high speeds, a possibility that does not appear to have been recorded in the conveyor belt experiments of Chambon and co-workers, but was noted in related numerical computations [31]. Time-dependent motions were observed in our rotating drum at high rotation rates, but at these speeds the flow was also fully three-dimensional and well beyond the scope of our two-dimensional theory (fluid flowed sideways, collecting into localized patches that could be carried all the way around the drum or dripped off the upper surface).

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APPENDIX: STEADY SOLUTIONS WITHOUT SURFACE TENSION

Without surface tension, the steady solutions, with $q = Q(\eta, \Gamma) = \text{const.}$, satisfy a first-order differential-algebraic problem. For the Newtonian problem, Tirumkudulu and Acrivos [44] pointed out some of the numerical difficulties inherent in the Newtonian version of this problem. For small f , however, the solutions can be found using a relatively simple shooting technique: we first parametrize the solution by $\eta_{\max} = \eta(\theta_{\max})$, the height of the peak in the pool. Because the flux is given by Eq. (40) for $n = 1$, or more generally by Eq. (41), the peak location θ_{\max} (where $\eta_\theta = 0$) follows from $q = Q(\eta_{\max}, \sin \theta_{\max})$. We may then integrate the ODE for $\eta(\theta)$ to the left and right of the peak. Progressing to the right, the pool ends when the solution reaches the residual film, dictated approximately by $q = Q(\eta, \sin \theta)$, or Eq. (39) for $n = 1$. The integration can then be terminated before reaching the singular point at $\theta = \frac{1}{2}\pi$, and the solution continued using the approximation for the film. Further right, the film encounters a yield point where $\eta = q$, and the film plugs up thereafter. To the left, the integration either again reaches the film solution, or hits another yield point (again with $\eta = q$). Over the rigid sections, the film thickness maintains the value $\eta = q$, until one reaches the other yielded section surrounding $\theta = \pi$, where the solution is again given approximately by $q = Q(\eta, \sin \theta)$. In the Newtonian case ($\text{Bi} = 0$ and $n = 1$), the strategy is the same, but for the omission of any plugs. The last part of the exercise is to iterate the procedure, adjusting the value of η_{\max} until the mass constraint (33) is satisfied. Sample constructions are shown in Fig. 16.

Evidently, when the pool ends to the left at a yield point, the free surface slope becomes discontinuous (see Fig. 5). This feature becomes smoothed by surface tension to generate the viscoplastic wavetrain seen in Figs. 6 and 7, and discussed more thoroughly by Jalaal *et al.* [26,27]. The transition from the pool to the film is smoother to the right, and only features undulations with sufficient surface tension (cf. Fig. 5).

For $\Lambda \gg 1$, we enter the plastic limit of the problem in which the bulk of the fluid flow is controlled by the yield stress. In this situation, the pseudoplug fills most of the layer, $\mathcal{Y} \rightarrow 0$, and $q \rightarrow \text{Bi}$ (see Sec. III B). The solution can then be written down more explicitly: the residual film is mostly plugged up, with $\eta = q = \text{Bi}$, except over the region spanning the edge of the pool to $\theta = \pi$. Here, and over the pool, the problem reduces to the ODE (47), with the approximate solution discussed in Sec. III C. The contributions of each region to the constraint (33) is as follows: the area of the pool is approximately $\frac{4}{3}\eta_{\max}\sqrt{2f\eta_{\max}}$ (see Fig. 6), whereas the rigid section of the residual film has area $\frac{3}{2}\pi\text{Bi}$. The yield section of the film extends from the right-hand border of the pool, at position $\theta_{\max} + \sqrt{2f\eta_{\max}} \sim O(f^{\frac{1}{3}})$ to $\frac{1}{2}\pi$; the corresponding area is then, to leading-order, $\frac{1}{3}\text{Bi} \log f^{-1}$. The constraint then implies that

$$\frac{4}{3}\eta_{\max}\sqrt{2f\eta_{\max}} + \frac{3}{2}\pi\text{Bi} + \frac{1}{3}\text{Bi} \log f^{-1} = \pi, \tag{A1}$$

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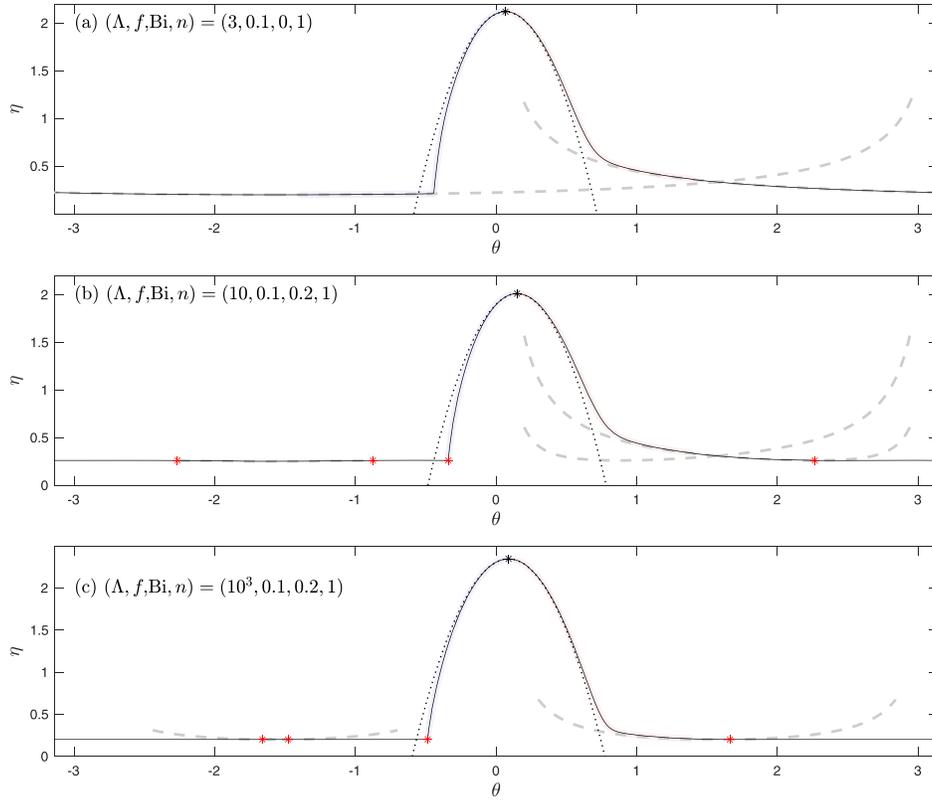


FIG. 16. Construction of steady solutions without surface tension for (a) Newtonian and (b, c) Bingham fluid (parameter settings indicated). Beginning from the peak, we integrate left (blue solid) and right (red solid) until the solution converges to one of the residual film solutions (gray dashed) or a yield point (red stars). Over any plugs, $\eta = q$. The thin black line shows the constructed solutions. The dotted line shows the approximation for the pool in panel (38). For panel (c), the steady solution is given approximately by Eq. (47) and the film solution is just $\eta \approx \text{Bi}/|\sin \theta|$ or $\eta \approx \text{Bi}$.

546 OR

$$\eta_{\max} = \left(\pi - \frac{3}{2}\pi \text{Bi} - \frac{1}{3}\text{Bi} \log f^{-1} \right)^{\frac{2}{3}} \left(\frac{9}{32f} \right)^{\frac{1}{3}}. \quad (\text{A2})$$

547 This approximation fails for $f \rightarrow 0$, corresponding to a fill-fraction regime where the left-hand edge
548 of the yielded residual film has greatest area.

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