Highly symmetric travelling waves in pipe flow

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The recent theoretical discovery of finite-amplitude travelling waves (TWs) in pipe flow has reignited interest in the transitional phenomena that Osborne Reynolds studied 125 years ago. Despite all being unstable, these waves are providing fresh insight into the flow dynamics. We describe two new classes of TWs, which, while possessing more restrictive symmetries than previously found TWs of Faisst & Eckhardt (2003 Phys. Rev. Lett. 91, 224502) and Wedin & Kerswell (2004 J. Fluid Mech. 508, 333–371), seem to be more fundamental to the hierarchy of exact solutions. They exhibit much higher wall shear stresses and appear at notably lower Reynolds numbers. The first M-class comprises the various discrete rotationally symmetric analogues of the mirror-symmetric wave found in Pringle & Kerswell (2007 Phys. Rev. Lett. 99, 074502), and have a distinctive double-layered structure of fast and slow streaks across the pipe radius. The second N-class has the more familiar separation of fast streaks to the exterior and slow streaks to the interior and looks like the precursor to the class of non-mirror-symmetric waves already known.

Keywords: pipe flow; travelling waves; turbulence transition

1. Introduction

The stability of the flow of a Newtonian fluid such as water through a straight pipe of constant circular cross section has fascinated scientists ever since Reynolds’ (1883) original experiments. Reynolds realized that there was a single non-dimensional control parameter for the flow, \( Re = \frac{UD}{\nu} \) (where \( U \) is the mean velocity; \( D \) is the pipe diameter; and \( \nu \) is the kinematic viscosity of the fluid), but that there was no unique value of this parameter for transition to occur. Instead, he noticed that the value of \( Re \) required for transition depended critically on how cleanly he performed the experiment, i.e. on the ambient level of noise. Later theoretical and computational work confirmed that the steady, unidirectional, parabolic, laminar ‘Hagen–Poiseuille’ flow (Hagen 1839; Poiseuille 1840) realized at low \( Re \) is linearly stable at least to \( Re = 10^7 \) (Meseguer & Trefethen 2003), and that pipe flow transition must be a finite-amplitude process. The variability

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in the critical $Re$ across experimental realizations has indicated that this process is also sensitive to the exact form of disturbances present in the flow (e.g. Peixinho & Mullin 2007).

Once triggered, transition is generally abrupt and can lead to both localized and global forms of turbulence depending on $Re$. For $1760 \leq Re \leq 2300$, turbulence is localized into ‘puffs’ (Wygnanski & Champagne 1973), which have a weak front, a sharp trailing edge and maintain their length at about $20D$ as they propagate along the pipe. For $2300 \leq Re$, the turbulence delocalizes into ‘slug’ turbulence (Wygnanski et al. 1975), which has a sharp front travelling faster than the mean flow and a trailing edge travelling slower so that it expands while propagating downstream. As a result of this varied behaviour and the sensitivity of the transition process, pipe flow has become the canonical example of shear flow transition.

From a dynamical systems viewpoint, the existence of turbulence requires the presence of simpler solutions of the governing equations to provide a framework in state space to sustain turbulent trajectories from simply relaminarizing. Natural questions are then whether the emergence of such alternative solutions can be used to predict the threshold (lowest) $Re$ for transition and what role they play in the process, as well as the final turbulent state. Steady states and travelling waves (TWs)—fixed points in an appropriately chosen Galilean frame—are the simplest solutions.

The first step towards identifying such states was taken by Smith & Bodonyi (1982), who predicted the existence of helical modes (with swirl) through a critical layer analysis. Landman (1990), however, failed to find any numerical evidence for their existence. In the same year, Nagata (1990) used a homotopy approach to find finite-amplitude solutions in plane Couette flow starting from the linearly unstable rotating situation. A similar attempt using rotating pipe flow, however, failed to reach the vanishing rotation limit (Barnes & Kerswell 2000). By this time, a more systematic approach was emerging based upon the idea of a self-sustaining process. Hamilton et al. (1995) suggested that streamwise rolls could create streamwise streaks, which in turn could reinforce the original rolls by some unspecified nonlinear process. Waleffe (1997, 1998) closed the loop by showing that the streaks had inflectional instabilities, which led to what he described as ‘wriggling’ in such a manner that energy is fed back into the initial rolls. He then identified such solutions in plane Couette flow using a constructive homotopic approach. Faisst & Eckhardt (2003) and Wedin & Kerswell (2004) employed this technique within the geometry of pipe flow to find the expected TWs, although all were without swirl and therefore unrelated to the predictions of Smith & Bodonyi (1982). Experimental observations (Hof et al. 2004, 2005) and numerical computations (Kerswell & Tutty 2007; Schneider et al. 2007a; Willis & Kerswell 2008) then followed, which indicated that these waves are realized, albeit transiently, as coherent structures in turbulent pipe flow. These waves are all born in saddle-node bifurcations with the lowest at $Re=1251$ for TWs with a discrete rotational symmetry. Pringle & Kerswell (2007) substantially lowered this to 773 by finding TWs with no rotational symmetry. The $Re$ gap between when alternative solutions exist and when sustained turbulence occurs is then certainly positive but also too large to be a useful predictor.

Recently, Duguet et al. (2008) have adopted an entirely different approach to finding TWs based upon identifying episodes in time-dependent pipe flow, which exhibit temporal periodicity. To generate a long-time signal for this search,
sustained flow dynamics away from the laminar state is required. The obvious source for this is studying flow on the turbulent attractor, but they realized that there was a more promising alternative where the flow is less energetic and periodic episodes are more clearly evident: the laminar–turbulent boundary. This is the set of flows which separates initial conditions that immediately relaminarize from those that undergo a turbulent episode. An initial condition chosen on this boundary or ‘edge’ (Skufca et al. 2006) will evolve in time forever treading a tightrope between relaminarizing and becoming turbulent (Toh & Itano 1999; Itano & Toh 2001; Schneider et al. 2007b). Duguet et al. (2008) identified three new TWs using velocity fields from near-time periodic episodes in this energetically intermediate flow dynamics as starting guesses for an iterative Newton–Krylov solver.

The purpose of this paper is to show how the TWs found by Duguet et al. (2008) led to the discovery of two new classes of mirror-symmetric TWs, hereafter referred to as the M- and N-classes, and a distinctly new family of mirror-symmetric TWs. Each class is partitioned by the discrete rotational symmetry (see (2.1) below) enjoyed by the waves, which defines a family within the class. This suggests a natural labelling system where, for example, the family of M-class TWs with \( m \) (integer) degree of discrete rotational symmetry will be referred to as the \( Mm \) family. We argue that the mirror-symmetric wave of Pringle & Kerswell (2007) is actually part of the first family (M1) of the M-class, and ‘A3’ in the nomenclature of Duguet et al. (2008) is a member of the second family (M2). We further report the families M3, M4, M5 and M6 which, unlike anything seen so far, all have a double-layered structure of fast and slow streaks across the pipe radius. The second N-class of TWs has ‘C3’ in Duguet et al. (2008) as part of its second family (N2), with the families N1, N3, N4 and N5 described here for the first time. In these waves, the fast streaks are positioned near the pipe wall and slow streaks in the interior, a separation familiar from the original non-mirror-symmetric TWs found by Faisst & Eckhardt (2003) and Wedin & Kerswell (2004). Since these waves only have shift-and-reflect symmetry (see §3 below), we refer to this original set as the S-class hereafter for convenience.

The plan of the paper is to start by discussing the various symmetries of the TWs which serve to categorize them in §2. Beyond these, all the TWs are parametrized by their axial wavenumber \( \alpha \), which is a continuous variable with a finite range at a given \( Re \). These waves are exact solutions of the Navier–Stokes equations when periodic flow conditions are imposed across the pipe with a spatial period equal to an integer multiple of \( 2\pi/\alpha \). The three TWs found recently by Duguet et al. (2008) are then introduced in §3. Numerical codes developed by Wedin & Kerswell (2004) are used to trace out these solutions of the Navier–Stokes equations, which are steady in a Galilean frame moving at a phase velocity of \( C\hat{z} \),

\[
-C \frac{\partial u}{\partial z} + u \cdot \nabla u + \nabla p = \frac{1}{2} \lambda \hat{z} + \frac{1}{Re} \nabla^2 u, \tag{1.1}
\]

through parameter space. Here, the mean velocity \( U \) and diameter \( D \) have been used to non-dimensionalize the system, \( \lambda \) is the friction factor measuring the imposed pressure gradient, \( p \) is strictly periodic and the usual cylindrical coordinates \((s,\phi,z)\) and the corresponding velocity components \((u,v,w)\) are adopted throughout. The numerical representation of the TWs is triply spectral with Fourier modes in the two periodic directions (axial and azimuthal) and
Chebyshev polynomials in the radial direction (see Wedin & Kerswell 2004 for details). A new TW classification is then introduced in §4 stimulated by the discoveries of many related families. A final discussion follows in §5.

2. Travelling wave symmetries

The original work by Faisst & Eckhardt (2003) and Wedin & Kerswell (2004) used the homotopy approach developed by Waleffe (1997) to look for exact steady solutions in an appropriately chosen Galilean frame. In this, an artificial body forcing is added to the Navier–Stokes equations in order to drive streamwise rolls. These advect the underlying shear towards and away from the pipe wall, thereby generating positive and negative streamwise anomalies in the mean flow called ‘streaks’. If these streaks are of sufficient amplitude, they become inflexionally unstable to streamwise-dependent flows, which if allowed to grow to sufficient amplitude, can replace the artificial forcing as the energy source for the rolls. By choosing the forcing functions to have certain discrete rotational symmetries defined by

\[ \mathbf{R}_m : (u, v, w, p)(s, \phi, z) \rightarrow (u, v, w, p)(s, \phi + 2\pi/m, z), \]

and selecting streamwise-dependent instabilities having the shift-and-reflect symmetry

\[ \mathbf{S} : (u, v, w, p)(s, \phi, z) \rightarrow (u, -v, w, p)(s, -\phi, z + \pi/\alpha), \]

where \( 2\pi/\alpha \) is the wavelength, exact TWs were found by continuing back to zero forcing in \( m = 2, 3, 4, 5 \) and 6 discrete rotational symmetry subspaces. All appear through saddle-node bifurcations with the lowest at \( Re = 1251 \) corresponding to a \( \mathbf{R}_3 \)-symmetric TW. Each TW family is parametrized by its axial wavenumber \( \alpha \), which, at a given \( Re \), has a finite range (see figs 1 and 2 of Kerswell & Tutty 2007).

Later, motivated by a chaotic flow state found on the laminar–turbulent boundary by Schneider et al. (2007b), Pringle & Kerswell (2007) generalized the forcing function to have no rotational symmetry (see also Mellibovsky & Meseguer 2007). As a result, they found an asymmetric TW (formally the missing \( \mathbf{R}_1 \) state possessing one pair of streamwise rolls) with \( \mathbf{S} \) symmetry. This TW originates through a symmetry-breaking bifurcation from a state that satisfies the additional shift-and-rotate symmetry

\[ \mathbf{\Omega}_m : (u, v, w, p)(s, \phi, z) \rightarrow (u, v, w, p)(s, \phi + \pi/m, z + \pi/\alpha), \]

with \( m = 1 \), and is therefore mirror symmetric about \( \phi = \pi/2 \) (the plane \( \phi = 0 \) being set by the shift-and-reflect symmetry). For general \( m \),

\[ S\Omega_m^j = Z_{j \pi/2m} \quad j = 1, 3, 5, ..., 2m - 1, \]

where \( \Omega_m^j \) implies \( \Omega_m \) is applied \( j \) times and

\[ Z_\psi : (u, v, w, p)(s, \phi, z) \rightarrow (u, -v, w, p)(s, 2\psi - \phi, z) \]

represents reflection in the plane \( \psi = \Psi \). Helical generalizations that satisfy a modified shift-and-rotate symmetry

\[ \Omega_1^\delta : (u, v, w, p)(s, \phi, z) \rightarrow (u, v, w, p)\left(s, \phi + \left(1 + \frac{\delta}{\alpha}\right)\pi, z + \pi/\alpha\right) \]

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3. TWs found on the laminar–turbulent boundary

The new TWs discovered by Duguet et al. (2008) were found by studying the flow dynamics on the laminar–turbulent boundary within the $R_2$-symmetric subspace. Surprisingly, at $Re=2875$ and periodic pipe length $5D$, the rotationally unrestricted situation only ever revealed the already-known asymmetric TW (Pringle & Kerswell 2007) and a weakly rotating version of it ($A'$ in Duguet et al. 2008). The fact that new $R_2$-symmetric TWs were observed as transient coherent structures only in the $R_2$-symmetric subspace calculations is due to the reduced number of unstable directions they have, as non-$R_2$-symmetric flows are removed from the dynamics. The chance of the flow ‘visiting’ a TW in phase space is

(where $\beta$ is the helicity; Pringle & Kerswell 2007), but no $S$ symmetry, were also found but always at higher $Re$, indicating that the flow prefers the streaks to be aligned with the flow rather than twisted around it.

Figure 1. The TWs (a) A3, (b) C3 and (c) Z2. (d) S2 is also included for comparison purposes (2$b$-1.25 in Kerswell & Tutty 2007). All are shown at their respective values of $a'$ (table 1) and at $Re=2400$ except for Z2, which is shown at $Re=3860$. The contours indicate the magnitude of the streamwise velocity difference from the underlying laminar flow. The contours and shading are standardized across the plots (and, more generally, figures 3, 4 and 10) with contour increments of 0.02 $U$ running from $-0.19U$ (grey) to $0.19U$ (white) (the shading adjacent to the pipe wall indicates zero). Arrows indicate the cross-stream velocities.
presumably related to how unstable it is (the spatial periodicity of the flow was also shortened to improve the stability of the TWs). In the notation of Duguet et al. (2008), the three new TWs were labelled as ‘A3’, ‘C3’ and ‘D2’ and all are mirror symmetric, even though this symmetry was not imposed on the flow. A3 and C3 were readily converged in the continuation codes of Wedin & Kerswell (2004), but D2 could not be due, we suspect, to insufficient axial resolution being achievable in the continuation code. The roll structure of D2 was, however, used to design a forcing in the homotopy approach, which successfully yielded a similar looking TW—called Z2—with the same symmetries as D2. That they were, in fact, different waves became apparent when Z2 could not be continued below \( Re = 3250 \), whereas D2 was discovered at \( Re = 2875 \). Despite the remarkably similar form of Z2 and D2, there was no readily apparent evidence that the calculated trajectories visited Z2 as well. The structure of A3, C3 and Z2 is shown in figure 1, compared with a known TW in S2, which, at \( \alpha = 1.25 \) and \( Re = 2400 \) is already known to be on the laminar–turbulent boundary (this is 2\( b \), 1.25 in the nomenclature of Kerswell & Tutty 2007).

While all the new families possess the apparently universal features of exact coherent structures known in wall-bounded shear flows—wavy streaks with staggered quasi-streamwise vortices (Nagata 1990; Waleffe 1998, 2001, 2003;)

<table>
<thead>
<tr>
<th>source</th>
<th>symmetries</th>
<th>( Re_{\text{lowest}} )</th>
<th>( \alpha^* )</th>
</tr>
</thead>
<tbody>
<tr>
<td>FE03 &amp; WK04</td>
<td>S &amp; R₂</td>
<td>1358</td>
<td>1.55</td>
</tr>
<tr>
<td>S3</td>
<td>S &amp; R₃</td>
<td>1251</td>
<td>2.44</td>
</tr>
<tr>
<td>S4</td>
<td>S &amp; R₁</td>
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<tr>
<td>S5</td>
<td>S &amp; R₅</td>
<td>2485</td>
<td>4.11</td>
</tr>
<tr>
<td>WK04</td>
<td>S &amp; R₀</td>
<td>2869</td>
<td>4.73</td>
</tr>
<tr>
<td>U₁</td>
<td>S &amp; Ω₁</td>
<td>3046</td>
<td>2.17</td>
</tr>
<tr>
<td>PK07</td>
<td>S &amp; R₄</td>
<td>773</td>
<td>1.44</td>
</tr>
<tr>
<td>here</td>
<td>S &amp; Ω₅</td>
<td>773</td>
<td>1.44</td>
</tr>
<tr>
<td>M2 (A3)</td>
<td>S, R₂ &amp; Ω₂</td>
<td>1125</td>
<td>2.0</td>
</tr>
<tr>
<td>M3</td>
<td>S, R₃ &amp; Ω₃</td>
<td>1552</td>
<td>2.2</td>
</tr>
<tr>
<td>M4</td>
<td>S, R₁ &amp; Ω₄</td>
<td>1824</td>
<td>2.6</td>
</tr>
<tr>
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<tr>
<td>M6</td>
<td>S, R₆ &amp; Ω₆</td>
<td>2531</td>
<td>3.5</td>
</tr>
<tr>
<td>N1</td>
<td>S &amp; Ω₈</td>
<td>1554</td>
<td>0.9</td>
</tr>
<tr>
<td>N2 (C3)</td>
<td>S, R₂ &amp; Ω₂</td>
<td>1141</td>
<td>1.2</td>
</tr>
<tr>
<td>N3</td>
<td>S, R₃ &amp; Ω₃</td>
<td>1037</td>
<td>2.0</td>
</tr>
<tr>
<td>N4</td>
<td>S, R₁ &amp; Ω₄</td>
<td>1290</td>
<td>2.5</td>
</tr>
<tr>
<td>N5</td>
<td>S, R₅ &amp; Ω₅</td>
<td>1622</td>
<td>2.9</td>
</tr>
<tr>
<td>Z2</td>
<td>Z &amp; R₂</td>
<td>( \approx 3250 )</td>
<td>0.8</td>
</tr>
<tr>
<td>D2</td>
<td>Z &amp; R₂</td>
<td>&lt;2875</td>
<td>?</td>
</tr>
</tbody>
</table>

Table 1. The various symmetries of all the TWs currently known. (FE04, Faisst & Eckhardt (2003); WK04, Wedin & Kerswell (2004); PK07, Pringle & Kerswell (2007). The ‘U1’ entry for WK04 is to indicate that they found a mirror-symmetric TW with no rotational symmetry but at high \( Re \) and its relationship with M1 and N1 is unclear.)
Figure 2. Stability of C3 within the $R_2$-symmetric subspace against the phase speed $C$ as $Re$ increases away from the saddle-node bifurcation at 1141 ($\alpha = \alpha^*$ and the corresponding $Re$ is across the upper $x$-axis). Each curve either indicates the locus of a real eigenvalue (solid curves) or a complex conjugate pair (dashed curves) as $Re$ changes. The black curves correspond to those that are symmetric under $S$, while the grey curves are antisymmetric under $S$. The vertical dotted line indicates the saddle node, of which the inset is a close-up. There are bifurcations at $C = 1.33, 1.38, 1.385, 1.388, 1.389$ (the saddle node), 1.44, 1.47 and 1.50. (The grey dotted line indicates a stable complex conjugate pair that was difficult to resolve.)

Faisst & Eckhardt 2003; Wedin & Kerswell 2004; Pringle & Kerswell 2007)—there are new structural features. A3 has a strikingly different cross-sectional profile from previously known TWs (e.g. S2), in that both fast and slow streaks are concentrated at the pipe wall leaving the interior relatively quiescent. Figure 1 shows that C3 and Z2 are also notably different too. However, it is the axial structure of Z2 that really makes it stand out. Z2 is neither $S$ symmetric nor $\Omega$ symmetric, but does have mirror symmetry about two perpendicular planes. There is no a priori reason to expect the flow to prefer one of these symmetries over any of the others, but to date all TWs in any of the canonical shear flows have always been $S$ symmetric (except, trivially, the helical modes of Pringle & Kerswell (2007); Waleffe, private communication). Therefore, Z2 and D2 are the first TWs found to possess only $Z$ symmetry ($Z_\psi$ with $\psi$ suppressed as there is no longer an origin for $\phi$). Both Z2 and D2 are, however, close to being $S$ symmetric in the sense that the simple indicator

$$\frac{\int_0^{2\pi} \int_0^1 (u - Su)^2 |_{z=0} s \, ds \, d\phi}{\int_0^{2\pi} \int_0^1 (u + Su)^2 |_{z=0} s \, ds \, d\phi} = O(10^{-3}).$$

(3.1)
By contrast, A3 and C3 are both $S$ and $\Omega_2$ symmetric, and hence also reflectionally symmetric about a plane at $\pm \pi/4$ to the plane of shift-and-reflect symmetry (horizontal in figure 1), i.e. they have $Z_{\pm \pi/4}$ symmetry. The modes C3 and A3 appear through saddle-node bifurcations at $Re = 1141$ and 1125 with optimal wavenumbers $\alpha^* = 1.2$ and 2.0, respectively, giving these lowest saddle-node bifurcations (table 1).

(a) Stability

An interesting feature of all TWs found so far is that they are unstable but only with an unstable manifold of very small dimension (invariably less than 10 for those checked so far; Faisst & Eckhardt 2003; Kerswell & Tutty 2007). The typical situation is illustrated in figure 2, which displays the spectrum of the C3 wave at the wavenumber $\alpha^*$ which gives its lowest saddle-node bifurcation. While $Re$ is double valued near the bifurcation, the phase speed $C$ is not, monotonically decreasing in value as the bifurcation point is crossed from lower to upper branch. This then provides a convenient abscissa to show how the stability changes on both upper and lower branches as $Re$ increases away from the saddle-node bifurcation. The stability analysis is restricted to disturbances periodic over the same wavelength as C3 and to the $R_2$-symmetric subspace. C3 is particularly interesting for two reasons: (i) the number of unstable
directions decreases down to one for $Re \geq 1826$ along the lower branch, and (ii) there are several bifurcations involving real eigenvalues. The implication of the first observation is that C3, which is on the laminar–turbulent boundary (in a pipe 2.5$D$ long at $Re = 2400$; Duguet et al. 2008), will become a local attractor there beyond $Re = 1826$ for $R_2$-symmetric flow (confirmed by Duguet et al. 2008). The second observation means that new solutions bifurcate from C3, which are also steady in an appropriately translating Galilean frame—that is to say that they are TWs (this is how the asymmetric TW of Pringle & Kerswell 2007 appears). There are four possibilities, a transcritical bifurcation where no symmetry is broken, and three types of symmetry-breaking pitchfork in which only $S$, $Z$ or $\Omega_2$ symmetry is retained (recall two symmetries imply the third). For example, the bifurcation at $(C,Re) = (1.47, 1449)$ is a mirror-symmetry-breaking pitchfork, which can be followed to produce TWs that resemble those of S2. D2 is plausibly the product of a $S$-symmetry-breaking pitchfork from C3 (D2 is very similar to Z2 and hence C3; figure 1).

Hopf bifurcations are the generic scenario, however, leading to more complicated relative periodic orbits such as those traced by Duguet et al. (2008). Within the $Re$ range of figure 2, C3 experiences three Hopf bifurcations (corresponding to the phase speeds $C = 1.33, 1.385, 1.50$), which all lead to relative periodic orbits with the $Z$ symmetry broken.

Figure 4. Slices of the lower branch N-class TW solutions at $Re = 2400$ and at each TW’s respective value of $a^*$. (a) N2 (C3), (b) N3, (c) N4 and (d) N5. Contour levels of the streamwise velocity perturbation are in increments of 0.02$U$ running from $-0.19U$ (grey) to $0.19U$ (white).
4. New mirror-symmetric classes

The fact that A3 and C3 appear at such low $Re$ strongly suggests that there are analogous waves in different rotational symmetry classes also existent at pre- and transitional $Re \lesssim 2400$. This indeed turns out to be the case with all except two (N1 and M3) of the new families of mirror-symmetric modes being easily found using the homotopy approach, once the appropriate $\Omega$ symmetry is incorporated. Selecting a roll forcing of $R_{2m}$ symmetry defined, in the notation of
Wedin & Kerswell (2004), by \( \lambda_{2m} \) with \( i=2, 3 \) or 4 invariably led to a subharmonic \( \mathbf{R}_m \)-symmetric streak instability, which could be easily tracked back to zero forcing (this was exactly the strategy used by Wedin & Kerswell to find their \( \mathbf{R}_1 \) TW (‘U1’ in table 1), which has two roll pairs and appears beyond \( Re=3046 \)). The remaining waves were found by homotopy in \( m \) starting from an \( M5 \) wave (\( M5 \) being used rather than \( M4 \) because of the similar parity of velocity components with respect to the radius) for \( M3 \) and from \( N3 \) for \( N1 \).

Inspection of the new lower branch TWs (figures 3 and 4\(^1\)) indicates that the mirror-symmetric family reported in Pringle & Kerswell (2007) is the first family (\( M1 \)) and \( A3 \) is a member of the second family (\( M2 \)) of a class of TWs (\( Mm \) with \( m=2, 3, 4, 5, 6, \ldots \)) with two layers of fast and slow streaks. Furthermore, \( C3 \) appears to be a member of a \( \mathbf{R}_2 \)-symmetric family (\( N2 \)) of another class of TWs (\( Nm \) with \( m=1, 2, 3, 4, 5, \ldots \)). In figure 5, \( M1 \) and \( N1 \) are shown alongside \( U1 \) found in Wedin & Kerswell (2004). That all three TWs can possess exactly the same symmetries yet have such different velocity fields and locations in phase space is in itself noteworthy. The ordering of the families within the respective classes is clear from the phase speed and wavenumber data shown in figure 6.

Extending the \( Mm \) and \( Nm \) families to higher rotational symmetry \( m \) is straightforward and could have been continued if the general trend had not already emerged. The bifurcation points for \( Mm \) TWs monotonically increase in

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\(^1\) Videos included in the electronic supplementary material are particularly illuminating.
Re and decrease in C with m (figure 7). This behaviour is not strictly true for the N-class where N3 has the lowest bifurcation point but is for m ≥ 3. N3 appears before N1 and N2 mimics the situation for the original S-class (non-mirror symmetric) TWs. The fact that the lowest Re for each Nm family as well as M1 and M2 is consistently smaller than the lowest Re for the equivalent Sm family is suggestive that the latter all bifurcate off the former in mirror-symmetry-breaking bifurcations. So far, there are two known examples of this: S1 (the asymmetric waves of Pringle & Kerswell 2007) bifurcating off M1, and a recently discovered connection between S3 and N3 at Re=2400. This latter bifurcation is shown clearly in a $L$ versus $\alpha$ plot (figure 8), which also highlights the smoothness of the N3 curve compared with the convoluted nature of the S3 curve (first plotted in fig. 1 of Kerswell & Tutty 2007). Another important feature is the very large friction factors (compared with the typical turbulent values) achieved by the upper branch N3 TWs. This is a general property of the N-class, as shown in figure 9. The new upper branch N- and M-class TWs (figure 10) all show an expected intensification of the slow and fast streaks, and the positioning of the fast streaks closer to the pipe wall than the corresponding lower branch solutions. This localization is particularly noticeable for the N-class TWs and leads to their large friction factors. Figure 9 also shows that the friction factors associated with some of the lower branch N-class TWs as well as M1 and M2 are significantly higher than time-averaged experimental values (as shown by the upper dotted line) and those of the S-class TWs.

Figure 8. The S3–N3 family connection at $Re=2400$ on a friction factor $A$ versus $\alpha$ plot. This plot for S3 appeared in fig. 1 of Kerswell & Tutty (2007). Now the bifurcation from the more symmetric N3 family is clear. The resolution is (8, 30, 6) in the nomenclature of Wedin & Kerswell (2004) throughout (the N3 curve is essentially unchanged for the higher resolution of (10, 30, 8)). The dotted line represents the lower bound given by the Hagen–Poiseuille solution ($A_{lam}=64/Re$), and the dashed line corresponds to the $Re=2400$ value of the log-law parametrization of the experimental data, $1/\sqrt{A}=2.0\log(Re\sqrt{A})-0.8$ (see Schlichting 1968, eqn (20.30)).
Finally, it is worth remarking that the M-class TWs start to appear at fascinatingly low $Re$, given their intricate double radial layer structure: for example, M4 appears at $Re = 1824$ (figure 7; table 1). That they exist at all is not a surprise (presumably TWs with three radial layers are possible too), but that they appear so early in $Re$ surely is and contrasts starkly with Z2, which does not emerge until $Re \approx 3250$.

5. Discussion

In this paper, we have described two new classes of mirror-symmetric TWs—the M- and N-classes which are also both shift-and-reflect symmetric—and a new mirror-symmetric family Z2 which is not. The M- and N-class waves appear earlier in $Re$ than the original non-mirror-symmetric waves of Faisst & Eckhart (2003) and Wedin & Kerswell (2004), suggesting that the latter are borne through generally supercritical symmetry-breaking bifurcations from them. This was already found to be the case for the M1 waves (Pringle & Kerswell 2007) but seems more generally true now that the various TW families making up each class have been found. The stability analysis of N2 (C3) presented here provides a timely reminder of the bifurcation possibilities: S or Z symmetries can be broken at pitchfork bifurcations leading to either Z-class TWs or S-class TWs.
found originally (Faisst & Eckhardt 2003; Wedin & Kerswell 2004). Hopf bifurcations, of course, give rise to periodic orbits in the Galilean frame of the TW or quasi-periodic orbits in the pipe (laboratory) frame. Tracing these requires a more sophisticated approach based upon time stepping the equations and searching for an exact return of an appropriately chosen Poincaré map (Viswanath 2007, 2008; Duguet et al. 2008).

The original motivation for searching for all the new TWs discussed here was the discovery of previously unknown TWs in the laminar–turbulent boundary by Duguet et al. (2008). It is likely that some of the lower branch solutions of these new waves similarly populate this boundary so that their stable manifolds also play a role in determining whether a given initial condition will lead to relaminarization or a turbulent evolution. There are, however, many interesting issues surrounding this assumption. For example, if a lower branch TW is on the boundary at one Re, will it still be at 10Re, and if not, how did it leave? In a long pipe where the TWs are continuously parametrized by their wavenumber over a finite range, at what critical wavenumber does the lower branch TW leave the boundary on its way to becoming an upper branch TW? The asymmetric wave (Pringle & Kerswell 2007) provides an obvious example being on the boundary for $\alpha=0.625$ (at $Re=2875$; Schneider et al. 2007b; Duguet et al. 2008), but presumably not for $\alpha=1.44$ where its wall shear stress is high (see fig. 6 of Pringle & Kerswell 2007). Hopefully, these issues will be discussed elsewhere in this celebratory issue.
The discovery that the upper branch solutions of the new N-class waves have such high wall shear stresses is, however, potentially more important. A recent attempt to extract coherent fast streak structures from pipe turbulence (Willis & Kerswell 2008) concluded that the TWs currently known were not energetic enough to be part of the turbulent attractor, as previously supposed and speculated about others as yet undiscovered. The new N-class of mirror-symmetric TWs may well be these missing waves as illustrated so starkly by figure 8, and the issue clearly needs to be revisited.

In this report, we hope a step forward has been made in appreciating the hierarchy of TWs that exist in pipe flow. Generally, the picture is that TWs with shift-and-reflect symmetry, mirror symmetry and a low degree of rotational symmetry seem to appear first and spawn further, less symmetric, TWs through bifurcations as Re increases. However, given the notorious complexity of the Navier–Stokes equations, it would be foolhardy to be too rigid about this. What should be absolutely clear, though, is that the pipe flow problem continues to fascinate and intrigue us fully 125 years after Reynolds’ original experiments.

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References


