Experimental Evidence of a Phase Transition in a Closed Turbulent Flow

P.-P. Cortet, Arnaud Chiffaudel, François Daviaud, Bérengère Dubrulle

To cite this version:

HAL Id: cea-01373378
https://cea.hal.science/cea-01373378
Submitted on 28 Sep 2016

HAL is a multi-disciplinary open access archive for the deposit and dissemination of scientific research documents, whether they are published or not. The documents may come from teaching and research institutions in France or abroad, or from public or private research centers.

L’archive ouverte pluridisciplinaire HAL, est destinée au dépôt et à la diffusion de documents scientifiques de niveau recherche, publiés ou non, émanant des établissements d’enseignement et de recherche français ou étrangers, des laboratoires publics ou privés.
Experimental Evidence of a Phase Transition in a Closed Turbulent Flow

P.-P. Cortet, A. Chiffaudel, F. Daviaud, and B. Dubrulle

CEA, IRAMIS, SPEC, CNRS URA 2464, Groupe Instabilités et Turbulence, 91191 Gif-sur-Yvette, France

(Received 3 May 2010; published 19 November 2010)

We experimentally study the susceptibility to symmetry breaking of a closed turbulent von Kármán swirling flow from \( \text{Re} = 150 \) to \( \text{Re} \approx 10^6 \). We report a divergence of this susceptibility at an intermediate Reynolds number \( \text{Re} = \text{Re}_\chi \approx 90,000 \), which gives experimental evidence that such a highly space and time fluctuating system can undergo a “phase transition.” This transition is further associated with a peak in the amplitude of fluctuations of the instantaneous flow symmetry corresponding to intermittencies between spontaneously symmetry breaking metastable states.

DOI: 10.1103/PhysRevLett.105.214501 PACS numbers: 47.20.Ky, 47.27.Cn

Phase transitions are ubiquitous in physical systems and are generally associated with symmetry breaking. For example, ferromagnetic systems are well known to undergo a phase transition from paramagnetism to ferromagnetism at the Curie temperature \( T_c \). This transition is associated with a symmetry breaking from the disordered paramagnetic—associated with a zero magnetization—toward the ordered ferromagnetic phase—associated with a finite magnetization [1]. In the vicinity of \( T_c \), a singular behavior characterized by critical exponents is observed, e.g., for the magnetic susceptibility to an external field.

Our experimental setup consists of a Plexiglas cylinder of radius \( R = 100 \) mm filled up with either water or water-glycerol mixtures. The fluid is mechanically stirred by a pair of coaxial impellers rotating in opposite directions (Fig. 1). The impellers are flat disks of radius 0.925R, fitted with 16 radial blades of height 0.2R and curvature radius 0.4625R. The disks’ inner surfaces are 1.8R apart setting the axial distance between impellers from blades to blades to 1.4R. The impellers rotate, with the convex face of the blades pushing the fluid forward, driven by two independent brushless 1.8 kW motors. The rotation frequencies \( f_1 \) and \( f_2 \) can be varied independently from 1 to 12 Hz.

Velocity measurements are performed with a stereoscopic PIV system provided by DANTEC Dynamics. The data provide the radial \( u_r \), axial \( u_z \), and azimuthal \( u_\phi \) velocity components in a meridian plane on a 95 × 66 points grid with 2.08 mm spatial resolution through time series of 400 to 27,000 fields regularly sampled, at frequencies from 1 to 15 Hz, depending on the turbulence intensity and the related need for statistics. The control parameters of the
studied von Kármán flow are the Reynolds number $\text{Re} = \pi (f_1 + f_2) R^2 / \nu$, where $\nu$ is the fluid viscosity, which controls the intensity of turbulence and the rotation number $	heta = (f_1 - f_2)/(f_1 + f_2)$, which controls the asymmetry of the forcing conditions. The rotation frequencies $f_{1,2}$ are regulated by servo loop control and we obtain for $\theta$ a typical absolute precision of $1 \times 10^{-3}$ and time fluctuations of the order of $\pm 2 \times 10^{-4}$. The correlation between these fluctuations and the flow dynamics is negligible.

When $\theta = 0$, the experimental system is symmetric with respect to any $\mathcal{R}_\pi$ rotation exchanging the two impellers: the problem conditions are invariant under $\pi$ rotation around any radial axis passing through the center of the cylinder (Fig. 1). The symmetry group for such experimental systems is $O(2)$ [12]. When $\theta \neq 0$, the experimental system is no longer $\mathcal{R}_\pi$ symmetric, and the symmetry switches to the $SO(2)$ group of rotations. However, the parameter $\theta$, when small but nonzero, can be considered as a measure of the distance to the exact $O(2)$ symmetry: the $SO(2)$ system at small $\theta$ can be considered as a slightly broken $O(2)$ system [13,14]. Depending on the value of $\theta$, the flow can respond by displaying different symmetries: (i) the exact $\mathcal{R}_\pi$-symmetric flow composed of two toric recirculation cells separated by an azimuthal shear layer located at $z = 0$ when $\theta = 0$ [Figs. 2(a) and 2(c)]; (ii) an asymmetric two-cells flow, the shear layer being closer to the slowest impeller ($z \neq 0$), when $\theta \neq 0$ [Figs. 2(b) and 2(d)]; (iii) and finally, a fully nonsymmetric one-cell flow, the whole shear layer being concentrated in between the blades of the slowest impeller, when $\theta$ becomes large enough [7,10,15].

In order to quantify the distance of the flow to the $\mathcal{R}_\pi$ symmetry, we use the normalized and space-averaged angular momentum $I(\text{Re}, \theta, t)$ as order parameter:

$$I(t) = \frac{1}{\mathcal{V}} \int_V r \, dr \, dz \, \frac{ru_\varphi(t)}{\pi R^2 (f_1 + f_2)},$$

where $\mathcal{V}$ is the volume of the flow [16]. An example of time variation of $I(t)$ at $\theta = 0$ in the turbulent regime is provided in Fig. 3. We assume that ergodicity holds, meaning that the instantaneous turbulent flow is exploring in time its energy landscape according to its statistical probability. In this framework, the time-average value $\bar{I}$ of $I(t)$ is equivalent to a statistical mechanics ensemble average providing the average is performed over a long enough duration in order to correctly sample the slowest time scales. Then, using this ensemble averaged order parameter, we define a susceptibility of the flow to symmetry breaking as $\chi_I = \partial \bar{I} / \partial \theta |_{\theta=0}$. Note that $\bar{I}$ is proportional to the mean altitude $z_\varphi$ of the shear layer which is the natural measure of the flow symmetry. Contrary to $z_\varphi$, $I(t)$ is defined for any instantaneous velocity field, including turbulent ones.

In the nonfluctuating laminar case, when $\theta = 0$, $\bar{I} = 0$ due to the symmetry of the flow. In contrast, as $\theta$ drifts away from 0, the value of the angular momentum $\bar{I}$ becomes more and more remote from zero as the asymmetry of the flow grows. In such a framework, there is a formal analogy between ferromagnetic and turbulent systems. For ferromagnetism [respectively, turbulence], the order parameter is the magnetization $M(T, h)$ [the angular momentum $I(\text{Re}, \theta)$], the symmetry-breaking parameter is the external applied field $h$ [the relative driving asymmetry $\theta$], the control parameter is the temperature $T$ [the Reynolds number Re, or a function of it].

In the following, we first investigate the influence of turbulence on $\bar{I}$ and $\chi_I$ as $\text{Re}$ increases from 150 to $10^6$. In the laminar flow at $\text{Re} = 150$, the symmetry parameter $\bar{I} = I(t)$ evolves linearly with $\theta$ [Fig. 4(c)] and the susceptibility is $\chi_I = 0.240 \pm 0.005$. Increasing the Reynolds number, one expects to reach fully developed turbulence around $\text{Re} = 10000$ [10]. In such turbulent regimes, velocity fields

![FIG. 2 (color online).](image)

(a),(b) Schematic drawings of the flow topology and (c),(d) corresponding experimental maps of mean velocity field of the turbulent von Kármán flow at $\text{Re} = 800000$ for (c) $\theta = 0$ ($\bar{I} = 0$) and (d) $\theta = -0.0147$ ($\bar{I} < 0$). The color maps the azimuthal velocity $u_\varphi$, from blue to red ("jet" color map), whereas the arrows map the $(u_r, u_z)$ field. The resolution has been reduced by a factor of 2 for better visibility. The $r = -r$ symmetry in (c) and (d) reveals that the time-averaged mean fields are axisymmetric.

![FIG. 3 (color online).](image)

Global angular momentum $I(t)$ as a function of time for an experiment performed at $\text{Re} = 43000$ for $\theta = 0$. Green (light gray) dots are PIV data sampled at 15 Hz, and blue (dark gray) dots correspond to 1 Hz low-pass filtered data $I_f(t)$. Eye-guiding lines have been drawn at $I(t) = \pm I_\varphi = \pm 0.04$. 

214501-2
of von Kármán flows are characterized by a high level of intrinsic fluctuations, i.e., fluctuations of the same order of magnitude as the mean values [15]. Therefore, even when $\theta = 0$, the $R_\pi$ symmetry is of course broken for the instantaneous flow. However, as is usually observed for classical turbulence, this symmetry is restored for the time-averaged flow [Fig. 2(c)], which proves that time averages are long enough to correctly sample the slowest flow time scales. Then, as in the case of the laminar flow, when $\theta$ is varied, we observe the breaking of the $R_\pi$ symmetry of the mean flow [Fig. 2(d)].

In Fig. 4(c), we see that, at $Re = 800000$, in the close vicinity of $\theta = 0$, $\bar{I}(\theta)$ evolves actually much more rapidly with $\theta$ than in the laminar case with the susceptibility being larger by more than 1 order of magnitude: $\chi_I = 9 \pm 1$. Therefore, turbulence seems to enhance dramatically the sensitivity of the flow to symmetry breaking. Furthermore, for intermediate $Re = 65000$, the slope of $\bar{I}(\theta)$ around $\theta = 0$ is even much steeper—$\chi_I = 43 \pm 1$—than for $Re = 800000$. In Fig. 4(a), we plot the susceptibility with respect to $Re$. We see that the susceptibility actually grows by more than 2 orders of magnitude—from 0.24 to 46—between $Re = 150$ and $Re \approx 90000$, before decreasing by a factor of 4 between $Re \approx 90000$ and $Re = 800000$. These results suggest a critical behavior for $\chi_I(Re)$ near $Re = Re_c = 90000 \pm 10000$: a divergence—revealing a continuous second order phase transition—or only a maximum—revealing either a subcritical bifurcation or a continuous transition with finite-size effect. This cannot be experimentally tested further since the highest measured $\chi_I$ are already of the order of the highest measurable value considering the $\theta$ precision of our setup. For higher $|\theta|$, we observe a crossover—for the slope—in the curve $\bar{I}(\theta)$ at $|\theta| = (6 \pm 1) \times 10^{-3}$ for $Re = 800000$ and very close to $\theta = 0$, at $|\theta| = (0.9 \pm 0.15) \times 10^{-3}$, for $Re = 65000$ [Fig. 4(c)]. For $|\theta| > |\theta|$ at large $Re$, we recover the laminar slow evolution of $I$ with $\theta$ up to $\theta = \pm 0.1$ where the flow bifurcates to the one-cell topology (not shown).

Since $\bar{I}(\theta)$ is quite independent of $Re$ for $|\theta| > |\theta|$ at large $Re$, we can extrapolate this linear behavior to $\theta = 0$. This extrapolation describes the ideal behavior at critical Reynolds number $Re_c$, if $\chi_I$ diverges: a jump of $\bar{I}$ between $-I_0$ and $+I_0$, where $I_0 \approx 0.05$. This can be interpreted as a spontaneous “turbulent magnetization” $I_0$ at $\theta = 0$—possibly affected by finite-size effects—by analogy with the spontaneous magnetization $M_0$ at zero external field for ferromagnetism. It is also similar to the experimental results of [17].

A signature of this magnetization can be seen on the instantaneous global angular momentum $I(t)$ for $Re$ near the peak of susceptibility and $\theta = 0$ (e.g., in Fig. 3). Indeed, one observes that $I(t)$ does not remain near zero (its mean value) but shows a tendency to lock preferentially on the plateaus $\pm I_0$ with $I_0 = 0.04 \pm 0.15$. Therefore, the turbulent flow explores a continuum of metastable symmetry-breaking patterns evidenced by $-0.04 \leq I_I(t) \leq 0.04$, $I_I(t)$ being the 1 Hz low-pass filtered value of $I(t)$. The global angular momentum actually fluctuates very much along time with two separate time scales: fast fluctuations related to “traditional” small scale turbulence and time intermittencies corresponding to residence time of few tens of seconds. If one performs a time average over one of these intermittent periods only, one obtains a time localized “mean” flow, which breaks spontaneously the symmetry, analogous to what is obtained for true mean flows when $\theta \neq 0$, as presented in Figs. 2(b) and 2(d).

The presence of strong fluctuations is not surprising here: close to a phase transition we expect critical fluctuations. To check this, we compute the standard deviation $\sigma_I(Re, \theta)$. Figure 4(b) shows, for $\theta = 0$, how $\sigma_I$ varies from zero in the laminar case to finite values for highly turbulent flows going through a maximum at Reynolds number $Re = Re_c = 45,000 \pm 10,000$ located below the peak of susceptibility at $Re_c$. Additionally, in Fig. 4(d), the dependence of $\sigma_I$ as a function of the symmetry control parameter $\theta$ reveals a strong difference between the two Reynolds numbers shown: $\sigma_I(\theta)$ presents a sharp and narrow peak at $\theta = 0$ for $Re = 65,000$, which does not exist for $Re = 800000$. The amplitude of this peak from its bottom to its top actually measures the additional amount of low frequency symmetry fluctuations due to the multiscale stability. This amount of fluctuations appears to be connected to the susceptibility increase below $Re_c$.

The previous experimental results set a strong connection between the spontaneous symmetry fluctuations of the flow
and the mean flow response to the system symmetry breaking: the interpretation of the large fluctuations of \( I(t) \) in terms of multistability suggests that the strong observed linear response of the mean flow [Fig. 4(c)] with respect to \( \theta \) in the close vicinity of \( \theta = 0 \) is the result of a temporal mixing between the metastable states in different proportions.

Nevertheless, as the Reynolds number is varied, two distinct maxima have been evidenced in our turbulent flow: one for the susceptibility \( \chi_t \) to the \( O(2) \) symmetry breaking near \( Re = Re_s \approx 90 000 \), and the other for the standard deviation of the global angular momentum \( I(t) \) near \( Re = Re_p \approx 45 000 \). In this Reynolds number range, the turbulence is generally expected to be fully developed; i.e., any nondimensional characteristic quantity of the flow should be \( Re \) independent. This is definitively not the case in this von Kármán experiment. Actually, visual observations of the flow reveal an increase of the average azimuthal number \( m \) of large scale vortices in the shear layer from \( m = 3 \) to \( m = 4 \) through an Eckhaus-type transition, between \( Re_p \) and \( Re_s \). In the following, we make the hypothesis that at least one critical phenomenon exists in the range \( 50 000 \leq Re \leq 100 000 \). Using the turbulent von Kármán-ferromagnetism analogy, we can check how classical mean field predictions for second order phase transition apply to our system. In terms of susceptibilities, it predicts a critical divergence at a temperature \( T_c \): \( \chi \propto |T - T_c|^{-1} \). Since the logarithm of the Reynolds number has already been proposed as the control parameter governing the statistical temperature of turbulent flows—\( T \sim 1/\log Re \) [18]—this prediction translates in our case into \( \chi_t \propto 1/\log Re - 1/\log Re_s |^{-1} \). This formula reasonably describes our data with \( Re_s \) between typically 70 000 and 90 000 [Fig. 4(a)] supporting the asymmetry between the two branches even with a unique exponent \(-1\). As far as fluctuations are concerned, it is difficult to find a reasonable critical exponent for \( \sigma_f \). However, the results show that the maxima for \( \chi_t \) and \( \sigma_f \) are clearly separated. This is at variance with classical phase transition theory, stating, e.g., in the fluctuation dissipation theorem, that \( \sigma_f \) should be proportional to \( \chi_t \). A reason for this discrepancy lies in the high level of intrinsic fluctuations in our system that makes this transition nonclassical. Instabilities or bifurcations occurring on highly fluctuating systems are commonly found in natural systems, and the literature reports transitions and symmetry breaking at high \( Re \) (see, e.g., Refs. [5–11,17,19–21]), but the corresponding theoretical tools are still not well settled. Existing studies of phase transitions in the presence of fluctuations generally consider systems in which an external noise—additive or multiplicative—is introduced [22]. In particular, it has been shown in models that multiplicative noise can produce an ordered symmetry-breaking state through a nonequilibrium phase transition [23]. This behavior could be at the origin of our observed transition. We can notice that, as \( Re \) increases, the turbulent momentization \( I_0 \) first increases and then decreases, contrary to the magnetization in the usual paraferromagnetic transition. This result is reminiscent of a reentrant noise-induced phase transition similar to that observed in the annealed Ising model [24,25]. The study of the evolution of \( I_0 \) and/or \( I_s \) with \( Re \) requires more statistics and is left for future work. Finally, our turbulent system, in which we can have access to both the spatiotemporal evolution of the states and the mean thermodynamic variables, appears as a unique tool to study out-of-equilibrium phase transitions in strongly fluctuating systems.

We thank K. Mallick for fruitful discussions. P.-P. C. was supported by Triangle de la Physique.

[16] Practically, \( I(t) \) is computed from PIV data restricted to a meridian plane only, but, since azimuthal flow fluctuations are strong, time average over several rotation periods—statistically equivalent to spatial azimuthal averaging—estimates correctly the 3D value of \( I(t) \) (see details in Ref. [15]).