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Statistical features of rapidly rotating decaying turbulence: Enstrophy and energy spectra and coherent structures

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In this paper, we investigate the properties of rapidly rotating decaying turbulence using numerical simulations and phenomenological modeling. We find that as the turbulent flow evolves in time, the Rossby number decreases to $\sim 10^{-3}$, and the flow becomes quasi-two-dimensional with strong coherent columnar structures arising due to the inverse cascade of energy. We establish that a major fraction of energy is confined in Fourier modes $(\pm 1, 0, 0)$ and $(0, \pm 1, 0)$ that correspond to the largest columnar structure in the flow. For wavenumbers (k) greater than the enstrophy dissipation wavenumber (k_d), our phenomenological arguments and numerical study show that the enstrophy flux and spectrum of a horizontal cross section perpendicular to the axis of rotation are given by $\epsilon_\omega \exp(-C(k/k_d)^2)$ and $C\epsilon_\omega^{2/3}k^{-1} \exp(-C(k/k_d)^2)$, respectively; for this 2D flow, ϵ_ω is the enstrophy dissipation rate, and C is a constant. Using these results, we propose a new form for the energy spectrum of rapidly rotating decaying turbulence: $E(k) = C\epsilon_\omega^{2/3}k^{-3} \exp(-C(k/k_d)^2)$. This model of the energy spectrum is based on wavenumber-dependent enstrophy flux, and it deviates significantly from power law energy spectrum reported earlier. *Published by AIP Publishing.* <https://doi.org/10.1063/1.5018346>

I. INTRODUCTION

One of the least understood non-equilibrium statistical mechanical systems is a fully developed turbulent flow.^{1–3} Kolmogorov^{4,5} proposed a theory for homogeneous and isotropic three-dimensional (3D) hydrodynamic turbulence, according to which the inertial-range energy spectrum $E(k) \sim \epsilon^{2/3}k^{-5/3}$. Here ϵ is the energy dissipation rate that equals the energy flux, and k is the wavenumber. This theory successfully explains many experimental and numerical findings.^{6–12} Kraichnan¹³ however showed that the two-dimensional (2D) hydrodynamic turbulence has further complexities—here the small wavenumber Fourier modes exhibit inverse cascade of energy with $E(k) \sim k^{-5/3}$, while the large wavenumber Fourier modes exhibit forward enstrophy cascade with $E(k) \sim k^{-3}$.

The fluid flows in nature and in laboratory are generally quite complicated. For example, they may involve external magnetic field, rotation, or buoyancy.^{6–12} The aforementioned hydrodynamic turbulence phenomenologies play a major role in modeling these flows. Researchers have shown that the external magnetic field^{14–16} and rotation^{7,17,18} typically affect the energy spectrum in the inertial range. In this paper, we address the turbulence phenomenology of rapidly rotating decaying turbulence.

Rotating flows are ubiquitous in nature, e.g., in ocean, atmosphere, celestial bodies, and in engineering applications. In the rotating frame of reference, the flow is affected by

the Coriolis and the centrifugal forces. While the centrifugal force may be absorbed into the pressure gradient term of the Navier–Stokes equation, the Coriolis force, which is perpendicular to the direction of rotation, tends to make the flow quasi-two-dimensional (2D). The Taylor–Proudman theorem⁶ predicts formation of the Taylor columns and emergence of quasi-2D behaviour. Note however that the Taylor–Proudman theorem is applicable in the linear limit of rapidly rotating steady flow. The quasi-2D behaviour of the turbulent rotating fluid is however much more subtle and is full of defining signatures that are still not fully understood.^{19–22} As mentioned earlier, 2D turbulence has its own caveats. The velocity component along the rotation axis, although relatively weaker than the perpendicular component, plays a significant role in rotating turbulence. Additionally, one cannot ignore the nonlinear advection term even when the fluid is rotating extremely fast. Thus such flows are quite delicate to model, and they are being studied vigorously.

Several models have been proposed to study the kinetic energy spectrum of rotating turbulence. Zeman²³ proposed a dual spectrum: the large- k modes exhibit Kolmogorov's spectrum ($k^{-5/3}$), while small- k modes show $k^{-11/5}$ spectrum. Zhou²⁴ proposed that $E(k) \sim k^{-2}$ for the entire inertial range when the rotation rate is very high, whereas Smith and Waleffe²⁵ argued that $E(k_\perp) \sim k_\perp^{-3}$, where k_\perp, k_\parallel are, respectively, the components of wavevector \mathbf{k} perpendicular and parallel to the rotation axis. Chakraborty²⁶ argued that $E(k) \sim k^m$, where $m \in [-2, -3]$; the spectral range can be further confined to $m \in [-2, -7/3]$ by kinetic helicity. Additionally, in the limit of very strong rotation, Kraichnan²⁷ proposed that $E(k) \sim \exp(-2\nu k^2 t)$, where ν is the kinematic viscosity and t is the time elapsed. This model assumes

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the absence of nonlinearity. In this paper, we show that the nonlinearity, though weak, is present, and it produces nonzero energy and enstrophy fluxes.

Many researchers^{28–32,41} have attempted to verify the aforementioned models of rotating turbulence using experiments. Morize *et al.*³³ studied decaying turbulence in a rotating tank and showed that the energy spectrum steepens from $k^{-5/3}$ to k^{-2} or even further as the rotation speed is increased. Morize *et al.*³³ and Staplehurst *et al.*³⁴ demonstrated asymmetry between cyclones and anticyclones. Moisy *et al.*³⁵ studied decay laws, anisotropy, and cyclone-anticyclone asymmetry in decaying rotating turbulence. Campagne *et al.*³⁶ carried out experiments on rotating turbulence and demonstrated a dual cascade of kinetic energy.

A large number of numerical experiments have been performed on rotating turbulence.^{37–40} Yang and Domaradzki,⁴² Müller and Thiele,⁴³ Mininni *et al.*,⁴⁴ and Biferale *et al.*⁴⁵ showed that the energy spectrum of rotating turbulence is approximately k^{-2} or k_{\perp}^{-2} . However, Smith and Lee⁴⁶ and Sen *et al.*⁴⁷ argued that the energy spectrum is proportional to k_{\perp}^{-3} . Deusebio *et al.*⁴⁸ studied how rotating flow transitions from three-dimensional to quasi-two-dimensional. Mininni⁴⁹ studied the rotating helical turbulence numerically and found that for $k > k_{\Omega}$, the energy spectrum exponent is -2.2 , and the helicity spectrum exponent is -1.8 ; here k_{Ω} is the Zeman wavenumber. For $k > k_{\Omega}$, they observed that the system becomes isotropic, and both energy and helicity spectra exhibit $k^{-5/3}$ scaling.

Baqui and Davidson⁵⁰ and Baqui *et al.*⁵¹ constructed a phenomenological theory of rotating turbulence and argued that the flow is anisotropic with $E(k_{\perp}) \sim \epsilon^{2/3} k_{\perp}^{-5/3}$ and $E(k_{\parallel}) \sim \epsilon^{2/3} (L_{\parallel}/L_{\perp})^{2/3} k_{\parallel}^{-5/3}$, where L_{\parallel} and L_{\perp} are, respectively, the integral length scales parallel and perpendicular to the direction of rotation. However, an intriguing aspect of the results of Baqui and co-workers is the absence of power law scaling in the inertial range of $E(k)$. This phenomenon, to the best of our knowledge, stands unexplained, and the quantitative nature of the spectrum is unreported. In this paper, we present a model to quantify this spectrum; this is the main result of this paper.

Another feature of rotating turbulence is that it has strong columnar structures. Similar structures are predicted by the Taylor–Proudman theorem in the linear limit [i.e., when the convective term in Eq. (1) is negligible] of rapidly rotating laminar flow.⁶ Nonlinear interactions among the inertial waves too yield quasi-2D behaviour^{45,52–54}—a result more relevant to the present study of turbulent rotating flow. Note that two-dimensionalization of the flow leads to an inverse cascade of energy (cf. Refs. 55–57) that strengthens the columnar structures. These features go beyond the linear limit. Using large-scale simulations on 4096³ grid, Biferale *et al.*⁴⁵ studied in detail the complex structures of rotating turbulence, in particular, the vortical structures and 3D anisotropic fluctuations.

For rapidly rotating flows, researchers have observed that $u_{\perp} \gg u_{\parallel}$, yet $u_{\parallel} \neq 0$. Here, u_{\perp} and u_{\parallel} are, respectively, the magnitudes of the velocity components perpendicular and parallel to the rotation axis. Thus rapidly rotating flow is quasi-2D, not 2D. As a result, the properties of rotating turbulence differ

from 2D hydrodynamic turbulence. For example, the energy spectrum of rotating turbulence at small wavenumbers differs from $k^{-5/3}$, in contrast to $k^{-5/3}$ spectrum observed for 2D hydrodynamic turbulence at low wavenumbers. As we show later in the paper, the columnar structures are observed in strongly rotating flows. Note that similar features have been observed in magnetohydrodynamic (MHD) turbulence⁵⁸ and quasi-static (QS) MHD turbulence.^{59–61} Interestingly, there is non-trivial energy exchange between the perpendicular and parallel components of the velocity.^{45,52–54,58–61} There have been various attempts to quantify anisotropy in rotating turbulence. For example, Delache *et al.*⁶² studied scale-by-scale anisotropy using the ring spectrum and showed that the flow was nearly isotropic for $k > k_{\Omega}$, but strongly anisotropic for $k < k_{\Omega}$.

Rotation suppresses the energy cascade, hence several researchers have revisited the decay law of total energy. Thiele and Müller⁶³ observed that the total energy $E \sim t^{-1.5}$ for non-rotating flows, but the exponent decreases from 1.5 to ≈ 0.5 as the rotation speed is increased. Baqui and Davidson⁵⁰ also arrived at similar conclusions. Teitelbaum and Mininni⁶⁴ however argued that $E \sim t^{-1}$ for nonhelical rotating flows, but $E \sim t^{-1/3}$ for the helical ones. Here the inverse cascade of energy suppresses the decay of turbulence.

In this paper, we investigate the energy spectrum and structures of rapidly rotating decaying turbulence. Here we focus on the nonhelical flows (zero kinetic helicity). We perform sufficiently high resolution spectral simulations that help us to analyze the asymptotic regime. We go up to 1024³ grid resolution and up to 155 eddy turnover time. In accordance with the earlier studies on rotating turbulence, we observe a strong inverse cascade of energy that strengthens the coherent columnar structures. The kinetic energy trapped in such structures dissipates very slowly, and hence the Reynolds number remains quite large with slow variation in time.

However, for rapidly rotating decaying turbulence, the energy contents of the intermediate and the small scales are quite small. We show that the Reynolds number based on the rms speed of these modes is quite small. Following the arguments similar to Verma *et al.*,⁶⁵ we show that the energy spectrum rapidly rotating decaying turbulence is $E(k) \sim k^{-3} \exp(-C(k/k_d)^2)$, where k_d is the enstrophy dissipation wavenumber and C is a positive real constant. To the best of our knowledge, such an energy spectrum has not been reported for decaying rapidly rotating turbulence.

We remark that the above results are for decaying turbulence. It is *a priori* not obvious that the forced rotating turbulence has similar behaviour as its decaying counterpart. The energy spectrum could in principle be dependent on the type and nature of external forcing and on the injection of kinetic helicity. However these topics are beyond the scope of this paper that exclusively deals with the strongly rotating decaying turbulence.

The structure of this paper is as follows: In Sec. II, we briefly review the existing turbulence phenomenologies of rotating turbulence and then describe our model for rapidly rotating decaying turbulence. Section III contains the details of numerical simulations. In Sec. IV, we describe the properties of the coherent columnar structures in the flow. In this section, we also discuss the inverse cascade of energy that

strengthens such structures. In Sec. V, we describe the spectra and the fluxes of energy and enstrophy. We conclude in Sec. VI.

II. PHENOMENOLOGY OF ROTATING TURBULENCE

The Navier–Stokes equation of an unforced incompressible fluid in rotating reference frame is

$$\frac{\partial \mathbf{u}}{\partial t} + (\mathbf{u} \cdot \nabla) \mathbf{u} = -\nabla p - 2\boldsymbol{\Omega} \times \mathbf{u} + \nu \nabla^2 \mathbf{u}, \quad (1)$$

$$\nabla \cdot \mathbf{u} = 0, \quad (2)$$

where \mathbf{u} and p are the velocity and pressure fields, respectively, $\boldsymbol{\Omega} = \Omega \hat{z}$ is the angular velocity of the rotating reference frame, ν is the kinematic viscosity, and $-2\boldsymbol{\Omega} \times \mathbf{u}$ is the Coriolis acceleration. The centrifugal acceleration has been absorbed in the pressure gradient term. We assume the rotation to be along the z direction.

The ratio of the magnitudes of $(\mathbf{u} \cdot \nabla) \mathbf{u}$ and the Coriolis acceleration is called the Rossby number, i.e.,

$$\text{Ro} = \frac{U_0}{\Omega L_0}, \quad (3)$$

where U_0 and L_0 are the large velocity and length scales, respectively. Coriolis force drives the perpendicular component of the velocity field, $\mathbf{u}_\perp = u_x \hat{x} + u_y \hat{y}$; this is one of the reasons why the rapidly rotating flows tend to be quasi-2D with $u_\perp \gg u_\parallel$ (u_\parallel is the magnitude of the velocity component along the z direction).^{19–22,66,67} In the present paper, we focus on rapidly rotating decaying flows, i.e., for very small Ro. Note that the Reynolds number $\text{Re} = U_0 L_0 / \nu$.

Owing to the quasi-2D nature of the flow, Kolmogorov's phenomenology for the 3D hydrodynamic turbulence does not apply to rotating turbulence. Researchers have proposed different models for rotating turbulence; some of these models are described below. Zeman²³ argued that the Coriolis force dominates the nonlinear advection term for $k < k_\Omega$, and vice versa for $k > k_\Omega$, where k_Ω is called *Zeman scale*. Kolmogorov's $k^{-5/3}$ spectrum is expected to hold for $k > k_\Omega$. Zeman²³ derived an expression for k_Ω by equating the advection term and the Coriolis acceleration at $k = k_\Omega$, i.e., $k_\Omega u_{k_\Omega}^2 \sim \Omega u_{k_\Omega}$. Kolmogorov's phenomenology for hydrodynamic turbulence yields $u_k \sim \epsilon^{1/3} k^{-1/3}$, substitution of which in $k_\Omega u_{k_\Omega}^2 \sim \Omega u_{k_\Omega}$ yields the Zeman wavenumber as

$$k_\Omega = \sqrt{\frac{\Omega^3}{\epsilon}}. \quad (4)$$

For $k < k_\Omega$, Zeman²³ employed scaling arguments and argued that $E(k) \sim \epsilon^{2/5} \Omega^{4/5} k^{-11/5}$. This phenomenology resembles the energy spectrum of stably stratified turbulence with buoyancy.^{68,69}

In another phenomenology, Zhou²⁴ modeled rotating turbulence in a spirit similar to Iroshnikov's¹⁶ and Kraichnan's¹⁵ model of magnetohydrodynamic turbulence and derived that $E(k) \sim (\epsilon \Omega)^{1/2} k^{-2}$; here the relevant time scale was taken to be Ω^{-1} . Zhou²⁴ generalized the above phenomenology to include nonlinear time scale and obtained dual scaling with k^{-2} spectrum for $k \ll k_\Omega$, and $k^{-5/3}$ spectrum for $k \gg k_\Omega$. Smith and Waleffe²⁵ equated the nonlinear advection term with the Coriolis force and obtained $E(k) \sim \Omega^2 k_\perp^{-3}$, where k_\perp is

the horizontal wavenumber. Using the perturbative approach, Chakraborty²⁰ showed that $E(k) \sim k^{-2.87}$ for weakly rotating turbulent systems. In latter part of this section, we show that for rapidly rotating turbulence ($\text{Ro} \ll 1$), the energy spectrum tends to be of the form $k^{-3} \exp(-C(k/k_d)^2)$, where k_d is the enstrophy dissipation wavenumber and C is a positive real constant. It is important to note that the Coriolis force does not do any work on the fluid.

Researchers have attempted to verify the above phenomenologies using experiments and numerical simulations. Some of these studies have been described in Sec. I, and we do not repeat them here. In the following discussion, we derive a model of $E(k)$ based on the flux variation with k .

For rapid rotation, the flow tends to be strongly quasi-two-dimensional, i.e., $u_z \ll u_\perp$. The Taylor–Proudman theorem⁶ predicts such structures in the linear regime. But in the nonlinear regime, structure formation is due to the inverse cascade of energy (see Sec. I). Note that in the linear regime, $\text{Re} = 0$, and hence $\Pi(k) = 0$. However, in the nonlinear regime, as will be shown in our numerical simulations, large-scale vortices are formed due to strong nonlinear effects. Therefore we study \mathbf{u}_\perp by taking a horizontal cross section of the flow profile. It is best to relate the 2D-sectional field with 2D hydrodynamic theory of Kraichnan.¹³ It is important however to keep in mind that the rotating flow is more complex due to the u_z component that couples with \mathbf{u}_\perp . External forcing is absent in decaying turbulence, but small wavenumber modes supply energy to the large wavenumber modes. Therefore, decaying 2D turbulence and rapidly rotating turbulence exhibit forward enstrophy cascade.⁷⁰

In Fourier space, the one-dimensional energy spectrum is defined as $E(k) = \sum_{k-1 < k' \leq k} \frac{1}{2} |\mathbf{u}(\mathbf{k}')|^2$, whose evolution equation is given by

$$\frac{\partial}{\partial t} E(k, t) = -\frac{\partial}{\partial k} \Pi(k, t) - 2\nu k^2 E(k, t), \quad (5)$$

where $\Pi(k, t)$ is the energy flux emanating from a wavenumber sphere of radius k at time t . For a steady or a quasi-steady state, $\partial E(k)/\partial t \approx 0$, hence

$$\frac{d}{dk} \Pi(k) = -2\nu k^2 E(k). \quad (6)$$

The energy flux and the spectrum, $\Pi(k)$ and $E(k)$, are two unknown functions whose solution cannot be obtained from a single equation, Eq. (6). For 3D hydrodynamic turbulent flows, Pao⁷¹ assumed that $E(k)/\Pi(k)$ is independent of ν and that it depends only on ϵ and k . Under these assumptions, we obtain the following solution for the above equation:

$$E(k) = K_{\text{Ko}} \epsilon^{2/3} k^{-5/3} \exp\left(-\frac{3}{2} K_{\text{Ko}} (k/k_\eta)^{4/3}\right), \quad (7)$$

$$\Pi(k) = \epsilon \exp\left(-\frac{3}{2} K_{\text{Ko}} (k/k_\eta)^{4/3}\right), \quad (8)$$

where ϵ is the energy dissipation rate, K_{Ko} is Kolmogorov's constant, and $k_\eta = (\epsilon/\nu^3)^{1/4}$ is Kolmogorov's wavenumber. Verma *et al.*⁶⁵ generalized the aforementioned Pao's phenomenology to laminar regime ($\text{Re} \lesssim 1$).

Since the rapidly rotating flows tend to be quasi-2D, it is important to briefly describe the phenomenology of 2D hydrodynamic turbulence. Kraichnan¹³ had reported that the two-dimensional hydrodynamic turbulence has dual spectrum—the energy exhibits inverse cascade for $k < k_f$, while the enstrophy $E_\omega^{(2D)} = \int d\mathbf{r} \omega^2/2 = \int d\mathbf{r} |\nabla \times \mathbf{u}|^2/2$ exhibits forward cascade for $k > k_f$, where k_f is the forcing wavenumber. Kraichnan¹³ showed that for $k < k_f$,

$$E^{(2D)}(k) = K'_{K_0} [\Pi^{(2D)}(k)]^{2/3} k^{-5/3} \quad (9)$$

with $\Pi^{(2D)}(k) = \text{const} < 0$. However, for $k > k_f$,

$$E_\omega^{(2D)}(k) = K_\omega [\Pi_\omega^{(2D)}(k)]^{2/3} k^{-1} \quad (10)$$

with $\Pi_\omega^{(2D)}(k) = \text{const} > 0$. Here $\Pi^{(2D)}(k)$ and $\Pi_\omega^{(2D)}(k)$ are the energy and enstrophy fluxes, respectively, K'_{K_0} is Kolmogorov's constant for 2D, and K_ω is the proportionality constant for the constant enstrophy flux regime.

In the subsequent discussion, we will show that for strongly rotating turbulence the intermediate and small-scale structures contain very small amount of energy. Also, since the flow is quasi-2D, we study the enstrophy and energy of a horizontal cross section perpendicular to the rotation axis. The evolution equation for the enstrophy spectrum is

$$\frac{\partial}{\partial t} E_\omega^{(2D)}(k, t) = -\frac{\partial}{\partial k} \Pi_\omega^{(2D)}(k, t) - 2\nu k^2 E_\omega^{(2D)}(k, t), \quad (11)$$

where $\Pi_\omega^{(2D)}(k, t)$ is the enstrophy flux for the 2D flow of a horizontal cross section. For a steady or a quasi-steady state, $\partial E_\omega^{(2D)}(k, t)/\partial t \approx 0$, hence

$$\frac{d}{dk} \Pi_\omega^{(2D)}(k) = -2\nu k^2 E_\omega^{(2D)}(k). \quad (12)$$

For the intermediate and small scales, following the work of Pao,⁷¹ we assume that $E_\omega^{(2D)}(k)/\Pi_\omega^{(2D)}(k)$ is independent of ν , and it depends only on the enstrophy dissipation rate, ϵ_ω , and k . Under this ansatz, $E_\omega^{(2D)}(k)$ and $\Pi_\omega^{(2D)}(k)$ are given by

$$E_\omega^{(2D)}(k) = C \epsilon_\omega^{2/3} k^{-1} \exp(-C(k/k_d)^2), \quad (13)$$

$$\Pi_\omega^{(2D)}(k) = \epsilon_\omega \exp(-C(k/k_d)^2), \quad (14)$$

where

$$k_d = \frac{\epsilon_\omega^{1/6}}{\sqrt{\nu}} \quad (15)$$

is the enstrophy dissipation wavenumber and ϵ_ω is the enstrophy dissipation rate.

For strongly rotating turbulence discussed in this paper, $\text{Re} \gg 1$, and the flow is quasi-2D. In this paper, we show some subtle differences between the 2D hydrodynamic turbulence and the rotating turbulence. We show that Eqs. (13) and (14) match quite well with our numerical results apart from prefactors. These results will be discussed in Sec. V.

We also remark that in the linear regime where nonlinearity is absent,

$$\frac{\partial}{\partial t} E_\omega^{(2D)}(k, t) = -2\nu k^2 E_\omega^{(2D)}(k, t), \quad (16)$$

which has solution of the form¹³

$$E_\omega^{(2D)}(k, t) = E_\omega^{(2D)}(k, 0) \exp(-2\nu k^2 t). \quad (17)$$

More importantly, the enstrophy flux must be zero for this case. This prediction is not applicable to our numerical results due to the presence of the nonlinear interactions, however weak, and nonzero enstrophy flux.

In Secs. III–V, we compare the phenomenology developed in this section with numerical results.

III. NUMERICAL SIMULATION

To investigate the dynamics of decaying rotating non-helical turbulent flow, we perform direct numerical simulation (DNS) using pseudo-spectral code Tarang.^{72,73} We solve Eqs. (1) and (2) in a 3D periodic and cubic box of size $(2\pi)^3$. We use the fourth-order Runge–Kutta scheme for time stepping, Courant–Friedrichs–Lewy (CFL) condition to obtain optimal time step (Δt), and 2/3-rule for dealiasing. We take $\nu = 10^{-3}$ and $\Omega = 16$. Note that these parameters are nondimensional, in the sense that $\nu/(U'L') \rightarrow \nu$ and $\Omega L'/U' \rightarrow \Omega$ with U' , L' as the characteristic velocity and length scale of the system in proper dimension (e.g., L' in meters). We remark that $U' \neq U_0$, where U_0 is the nondimensional rms velocity of the system. For our simulations, we employ grid-resolutions of 512^3 and 1024^3 . We observe that these grids yield similar results, thus we verify grid-independence of our results. For all our runs, $k_{\max} \eta > 1$, where η is Kolmogorov's length and k_{\max} is the highest wavenumber represented by the grid points. Hence all our simulations are well resolved. Note that in the above, $k_{\max} = N/2$ with N as the grid size.

First we generate a fully developed hydrodynamic turbulence ($\Omega = 0$) with $\nu = 10^{-3}$ on 512^3 grid with random forcing in the wavenumber band of (11, 12). Our forcing is such that it supplies a constant energy and zero kinetic helicity to the flow, which is achieved by a force field

$$\mathbf{f}(\mathbf{k}) = \alpha_k \mathbf{u}(\mathbf{k}), \quad \text{where } \alpha_k = \frac{\epsilon}{N_f |\mathbf{u}(\mathbf{k})|^2}. \quad (18)$$

Here ϵ is the energy supply rate, and N_f is the total number of wavenumber modes in the forcing band where \mathbf{f} is employed.⁷⁴

Now we use the steady-state data of the three-dimensional homogeneous isotropic fully developed hydrodynamic turbulence as an initial condition for the simulation of our rapidly rotating turbulence. Note that such initial conditions have been widely used in earlier simulations of decaying and rotating turbulence.^{50,62,75} We employ $\Omega = 16$. Note that the (nondimensional) frequency associated with the nonlinearity is $O(1)$; hence, our simulation with $\Omega = 16$ is reasonably fast rotating. We carry out simulations of the rotating flow with the same forcing as hydrodynamic simulations for 6 nondimensional time units after which the forcing is turned off. We set $t = 0$ here. The strongly rotating and decaying simulation starts at this stage, and it is carried out until $t = t_{\text{final}} = 155$. We also carry out simulations of the rotating flow on 1024^3 grid with the same set of parameters and initial conditions. Since 1024^3 simulation is much more expensive, we end this simulation at $t = 49$.

In Fig. 1, we show temporal evolution of various quantities for 512^3 (blue) and 1024^3 (red) grids. Figure 1(a) exhibits the evolution of total energy, $E = \int d\mathbf{r} u^2/2$, that decays from 0.8 to approximately 0.08 following a power law $E \sim t^{-0.33}$. This result is in general agreement with those of Thiele and

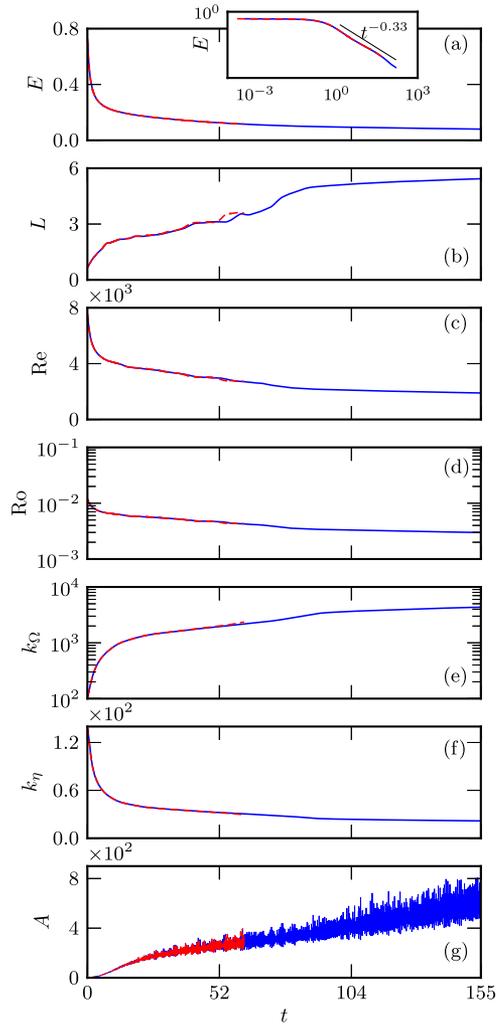


FIG. 1. For 1024^3 (red dashed) and 512^3 (blue) grid simulation with $\Omega = 16$, variation of characteristic system parameters with time t : (a) the total energy $E(t)$, (b) the integral length scale, L , (c) the Reynolds number, Re , (d) the Rossby number, Ro , (e) the Zeman wavenumber, k_{Ω} , (f) the Kolmogorov dissipation wavenumber, k_{η} , and (g) anisotropy of the system A . The inset in subfigure (a) is the log-log plot showing $E(t) \sim t^{-0.33}$.

Müller⁶³ and Baqui and Davidson.⁵¹ In Fig. 1(b), we plot the integral length scale defined by

$$L = 2\pi \frac{\int dk k^{-1} E(k)}{\int dk E(k)}. \quad (19)$$

We observe that beyond $t = 100$, $L \approx 5.43$, which is close to the box size (2π), thus signalling formation of large scale structures. As shown in Fig. 1(c), the Reynolds number $Re = U_0 L_0 / \nu$ is in the range of 2000–3000. For the computation of Re and Ro , we employ $L_0 = 2\pi$ (box size), and U_0 as the rms speed, which is given by $(2 \int_0^{k_{\max}} E(k) dk)^{1/2}$.

We compute the Rossby number Ro using Eq. (3) and plot its temporal variation in Fig. 1(d). The figure shows that Ro varies from 10^{-2} to 3×10^{-3} . Thus, the Rossby number at large times is quite small for our simulations. Hence, the Coriolis force dominates the nonlinear term of Eq. (1). We remark that some researchers⁵¹ report Ro based on the initial

TABLE I. Parameters of the direct numerical simulations (DNSs): List of total energy E ; integral length scale L ; Reynolds number Re ; Rossby number Ro ; Zeman wavenumber k_{Ω} ; Kolmogorov dissipation wavenumber k_{η} ; and anisotropy ratio A at $t = 49$, 148 for 512^3 grids and at $t = 49$ for 1024^3 grids.

Parameters	$t = 49$ ($N = 512^3$)	$t = 148$ ($N = 512^3$)	$t = 49$ ($N = 1024^3$)
E	0.13	0.08	0.13
L	3.11	5.40	3.10
Re	3003	1918	2989
Ro	0.005	0.003	0.005
k_{Ω}	1888	4300	1900
k_{η}	32	21	32
A	228	596	241

velocity of decaying turbulence. However our definition is based on the instantaneous velocity, thus Ro of Fig. 1(d) is that of instantaneous flow.

We also compute the Zeman wavenumber k_{Ω} using Eq. (4) and Kolmogorov's wavenumber k_{η} using $(\epsilon/\nu^3)^{1/4}$ and plot them in Figs. 1(e) and 1(f), respectively. We observe that in the asymptotic regime, $k_{\Omega} \gg 1$ indicating dominance of the Coriolis force. We observe that $k_{\Omega} > k_{\max}$, where $k_{\max} = N/2$ with N as the grid size. Hence the above estimate of k_{Ω} appears to be quite ambiguous. Note that the derivation of k_{Ω} using Eq. (4) assumes that for large k 's, the turbulence is isotropic and $E(k) \sim k^{-5/3}$, which is not the case for our simulations of strongly rotating turbulence (to be discussed in Sec. V). We show that strongly rotating turbulence makes the flow quasi-2D. Hence, strictly speaking, Eq. (4) cannot be employed to compute k_{Ω} .

In Fig. 1(g), we plot the anisotropy parameter

$$A = \frac{E_{\perp}}{2E_{\parallel}} \quad (20)$$

as a function of time. Here, $E_{\perp} = E_x + E_y$ and $E_{\parallel} = E_z$ [with $E_x = \int (u_x^2/2) d\mathbf{r}$, $E_y = \int (u_y^2/2) d\mathbf{r}$, and $E_z = \int (u_z^2/2) d\mathbf{r}$]. We observe that $A \gg 1$ indicating quasi-2D nature of the flow. We have also tabulated the values of E , L , Re , Ro , k_{Ω} , k_{η} , and A in Table I at $t = 49$, 148 for 512^3 grid simulations and at $t = 49$ for 1024^3 grid simulation, respectively.

In Sec. IV, we show that the strongly rotating flow is dominated by the columnar structures.

IV. COLUMNAR STRUCTURES AND ASSOCIATED FOURIER MODES

In Sec. III, we showed that the global parameters like the integral length scale and the anisotropy parameter indicate the presence of large scale structures. In this section, we describe these structures along with their associated Fourier modes.

We investigate the large scale structures by studying the flow profile in real space. In Figs. 2(a)–2(c), we exhibit the contour plots of the magnitude of the vorticity field, $|\omega|$, at $t = 49$, 98, and 148. These figures demonstrate the existence of strong vortical structures. To decipher the flow profiles of these columns, we take horizontal cross section of the flows at $z = \pi$ and present the density plots of ω_z superposed with the vector plot of $\mathbf{u}_{\perp} = u_x \hat{x} + u_y \hat{y}$. At $t = 49$, we observe

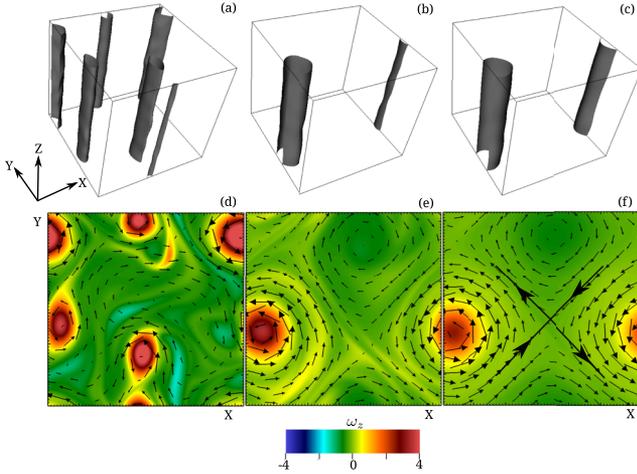


FIG. 2. For the rapidly rotating decaying turbulence on 512^3 grid, the top panels exhibit the isosurfaces of the magnitude of vorticity $|\omega|$ at (a) $t = 49$, (b) $t = 98$, and (c) $t = 148$. The bottom panels show the velocity vector plot superposed with the density plot for ω_z for the horizontal cross section at $z = \pi$ at (d) $t = 49$, (e) $t = 98$, and (f) $t = 148$. Note that (f) contains only one vortex that appears to be split in two due to the periodic boundary conditions.

four cyclonic vortices that have anticlockwise sense of rotation [see Fig. 2(d)]. Note the periodicity of the box. Subsequently these vortices merge and form a single cyclonic vortex, as shown in Figs. 2(c) and 2(f). In a periodic box, these vortices reside on a lattice along with weak anti-cyclonic vortices. These features are quite similar to those in 2D hydrodynamic turbulence.⁷⁰

The cyclonic and anti-cyclonic vortices are separated by a *saddle*, which is symbolised by a cross in Fig. 2(f). Morize *et al.*³³ and Staplehurst *et al.*³⁴ observed cyclonic/anti-cyclonic asymmetry in experiments, while van Bokhoven *et al.*⁷⁶ quantified the asymmetry between cyclonic and anti-cyclonic by studying skewness of vertical vorticity. Smith and Lee⁴⁶ argued that the cyclonic/anti-cyclonic asymmetry arises due to nonlinear interactions near resonance.

We remark that the size of the asymptotic (at large time) flow structures described above is proportional to the box size. We demonstrate this feature in Appendix A by simulating rotating turbulence in two boxes of sizes $(2\pi)^3$ and $(4\pi)^3$. These results show that the Fourier modes and their interactions are independent of the box size.

When we compare the flow structures of Figs. 2(e) and 2(f) with those found in two-dimensional Hamiltonian dynamics, the cross and the centres of the vortices of Fig. 2(f) correspond to the *saddle* and the *centres*, respectively. This similarity is due to the divergence-free condition of the velocity field that yields

$$\nabla \cdot \mathbf{u}_\perp = 0 \implies \frac{\partial \dot{x}}{\partial x} + \frac{\partial \dot{y}}{\partial y} = 0, \quad (21)$$

which is analogous to the equation for the conservation of phase space area of a two-dimensional Hamiltonian system.⁷⁷ Note that $u_z \ll u_\perp$, hence we treat our system as two-dimensional for the above discussion.

The emergence of large scale structures in the flow can be quantified using the energy contents of small wavenumber Fourier modes, which are listed in Table II. Evidently, in the

TABLE II. For rapidly rotating turbulence on 512^3 grid at $t = 4, 49, 148$, the energy contents of the dominant modes as percentage of the total energy. Note that $\mathbf{u}(-k_x, -k_y, 0)$, not listed in the table, has the same energy contents as that of $\mathbf{u}(k_x, k_y, 0)$. In the table, $E_{\text{mode}} = l u(\mathbf{k})^2/2$. For the modes $(0, 1, 0)$ and $(1, 0, 0)$, E_{mode}/E increases with time indicating strengthening of the vortical structures with time.

Mode (k_x, k_y, k_z)	E_{mode}/E (%)		
	$t = 4$	$t = 49$	$t = 148$
(0, 1, 0)	0.04	3.45	22.12
(1, 0, 0)	1.21	7.34	20.79
(1, 1, 0)	0.60	6.20	1.88
(-1, 1, 0)	0.32	1.75	1.66
(2, 1, 0)	0.65	1.60	0.47
(2, -1, 0)	1.08	2.17	0.44
(-2, 1, 0)	1.08	2.17	0.44
(1, 2, 0)	0.74	1.02	0.41
(1, -2, 0)	0.50	1.92	0.40
(2, -2, 0)	0.38	1.23	0.11
(-3, 0, 0)	0.68	0.36	0.10
(2, 2, 0)	0.52	0.96	0.10
(-1, -3, 0)	1.70	0.54	0.05
(3, 1, 0)	0.29	1.04	0.05
(-3, 1, 0)	2.80	0.19	0.05
(3, -2, 0)	0.45	0.19	0.02
(-2, -3, 0)	1.67	0.26	0.02
(-3, -2, 0)	0.89	0.04	0.02
Total (%)	13.084	32.43	49.13

asymptotic regime ($t = 148$), the Fourier mode $\mathbf{k} = (k_x, k_y, k_z) = (1, 0, 0)$ and $(0, 1, 0)$ are the most dominant modes, with the other strong Fourier modes being $(1, 1, 0)$ and $(-1, 1, 0)$. Note that we do not list the energies of $-\mathbf{k}$ modes because $\mathbf{u}(-\mathbf{k}) = \mathbf{u}^*(\mathbf{k})$. Hence, the energies $E(-1, 0, 0) = E(1, 0, 0)$ and $E(0, -1, 0) = E(0, 1, 0)$. When we add the energies of $(\pm 1, 0, 0)$ and $(0, \pm 1, 0)$, we observe that they contain approximately 80% of the total energy. These modes form a strong set of 2D vortices, as discussed in Appendix B and Fig. 8. The nonlinear dynamics of these modes are very interesting, and it may shed light on the cyclone-anticyclone antisymmetry in the presence of rotation. But these discussions are beyond the scope of this paper.

At $t = 148$, the sum the energies of the 18 dominant modes listed in Table II and those of their complex conjugate partners is approximately 98%. These modes lie within the sphere of radius 4. Hence, modes in the intermediate and small scales contain very small amount of energy. This result has a strong consequence on the energy and enstrophy spectra of the strongly rotating turbulence, which will be discussed in Sec. V.

In the following discussion, we argue why the small wavenumber modes become strong in rapidly rotating turbulence. The strong vortical structures of the flow indicate quasi-2D nature of the flow. This observation is reinforced by the fact that $A = E_\perp/(2E_\parallel) \gg 1$ [see Fig. 1(g)]. The flow becomes quasi-2D because of the Coriolis force that is active in the perpendicular plane (x, y) , as well as due to the inverse cascade of the kinetic energy from small scales (large k) to large scales (small k). This is in contrast to the emergence

of the Taylor columns in the linear limit, as predicted by the Taylor–Proudman theorem; the energy transfer is absent all together in the linear limit.

Let us now quantify the above observations using the mode-to-mode energy transfers and the energy flux. First we describe the energy transfers among the large scale Fourier modes. Dar *et al.*⁷⁸ and Verma¹⁴ showed that for an interacting triad of fluid flow $(\mathbf{k}', \mathbf{p}, \mathbf{q})$ that satisfies the relation $\mathbf{k}' + \mathbf{p} + \mathbf{q} = 0$, the rate of energy transfer from mode $\mathbf{u}(\mathbf{p})$ to mode $\mathbf{u}(\mathbf{k}')$ with mode $\mathbf{u}(\mathbf{q})$ acting as a mediator is

$$S(\mathbf{k}'|\mathbf{p}|\mathbf{q}) = -\text{Im}[(\mathbf{k}' \cdot \mathbf{u}(\mathbf{q}))(\mathbf{u}(\mathbf{p}) \cdot \mathbf{u}(\mathbf{k}'))]. \quad (22)$$

To investigate the growth of the large scale structures, we study the energy transfers among the small wavenumber modes listed in Table II. Figure 3 exhibits some of the dominant interacting triads involving small k Fourier modes. Note that a complete graph with N modes would contain approximately $N(N-1)/2$ edges, which is quite large for $N \sim 10$. Hence we show only some of the representative energy transfers. The numbers above the arrows represent the energy transfers. The most dominant energy transfer is in the triad $[\mathbf{k}' = (1, 1, 0), \mathbf{p} = (-1, 0, 0), \mathbf{q} = (0, -1, 0)]$ with the mode $\mathbf{u}(-1, 0, 0)$ supplying approximately 136×10^{-5} units of energy to the mode $\mathbf{u}(1, 1, 0)$. We observe that the dominant energy transfers are $(-1, 0, 0) \rightarrow (1, 1, 0) \rightarrow (0, -1, 0) \rightarrow (-1, 1, 0) \rightarrow (-1, 0, 0)$. Thus, the four modes $(1, 0, 0), (0, 1, 0), (1, 1, 0), (-1, 1, 0)$ play a critical role in the energy transfers in rotating turbulence. These issues will be studied in more detail in future. Figure 3 also exhibits other dominant energy transfers, but these transfers are order of magnitude smaller than those discussed above.

After the aforementioned discussion on the large scale structures and their associated Fourier modes, in Sec. V we present the energy and enstrophy spectra and fluxes of our system.

V. ENERGY AND ENSTROPHY FLUXES AND SPECTRA

The energy spectrum and flux provide valuable information about the flow. In this section, we compute these quantities and study their features.

We compute the energy flux at $t = 4, 49$, and 148 using the following formula:¹⁴

$$\Pi(k_0) = \sum_{k' > k_0} \sum_{p \leq k_0} S(\mathbf{k}'|\mathbf{p}|\mathbf{q}), \quad (23)$$

where k_0 is the radius of the wavenumber sphere from whom the flux is being computed, $S(\mathbf{k}'|\mathbf{p}|\mathbf{q})$ is defined in Eq. (22), and $\mathbf{k}' + \mathbf{p} + \mathbf{q} = 0$. We compute the energy flux at $t = 4, 49, 148$ using the 512^3 and 1024^3 grid data. These results are plotted in Figs. 4(a) and 4(c) for $t = 4$ (magenta), $t = 49$ (green), and 148 (red) of 512^3 grid simulation and for $t = 4$ (blue dashed) and 49 (cyan) of 1024^3 grid simulation. We observe a significant inverse cascade of kinetic energy at early stages. As time progresses, the kinetic energy flux becomes weaker and gets concentrated in the wavenumber band $k \in [1, 8]$. This feature is compatible with the strongly vortical quasi-2D structure of the flow. Ours is a decaying simulation, so we expect the inverse energy cascade regime to be narrower and the forward enstrophy cascade regime to be effective for a larger wavenumber range. For this reason, it is more appropriate to study the enstrophy spectrum and flux, for which we employ the data obtained from horizontal cross sections at $z = \pi/2, \pi$, and $3\pi/2$.

We compute the enstrophy flux of the 2D velocity field \mathbf{u}_\perp at the planes $z = \pi/2, \pi$, and $3\pi/2$ using

$$\Pi_\omega^{(2D)}(k_0) = \sum_{k' > k_0} \sum_{p \leq k_0} S^{\omega\omega}(\mathbf{k}'|\mathbf{p}|\mathbf{q}), \quad (24)$$

where

$$S^{\omega\omega}(\mathbf{k}'|\mathbf{p}|\mathbf{q}) = -\text{Im}[(\mathbf{k} \cdot \mathbf{u}_\perp(\mathbf{q}))\omega_z(\mathbf{p})\omega_z(\mathbf{k}')] \quad (25)$$

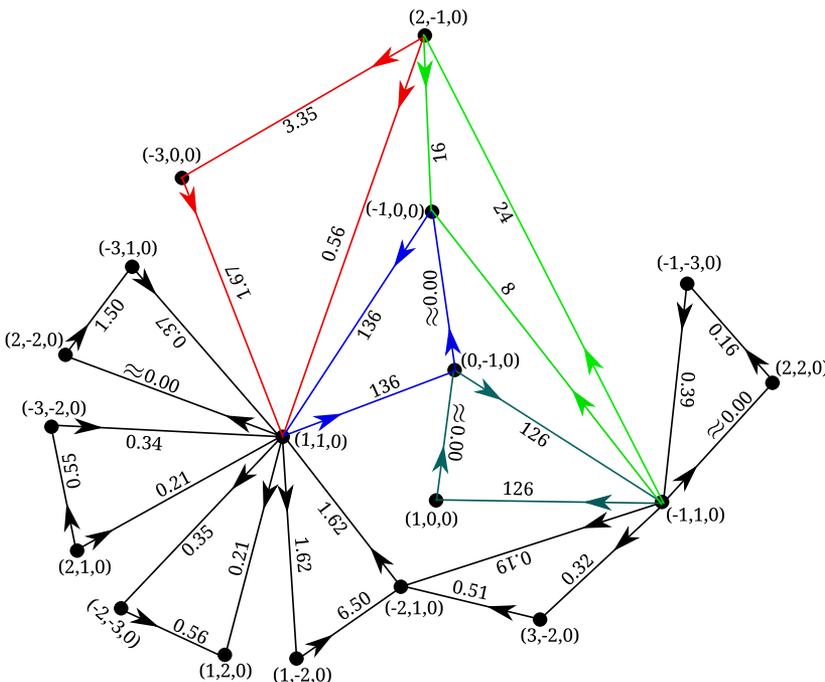


FIG. 3. For the rapidly rotating decaying simulation on 512^3 grid, the energy transfers $S(\mathbf{k}'|\mathbf{p}|\mathbf{q})$ for some of the dominant triads averaged over five eddy turnover times from $t = 148$ to 152 . The numbers above the arrows are $S(\mathbf{k}'|\mathbf{p}|\mathbf{q}) \times 10^5$ for convenient description. The most dominant energy transfers are in the triad $[(-1, 0, 0), (1, -1, 0), (0, -1, 0)]$.

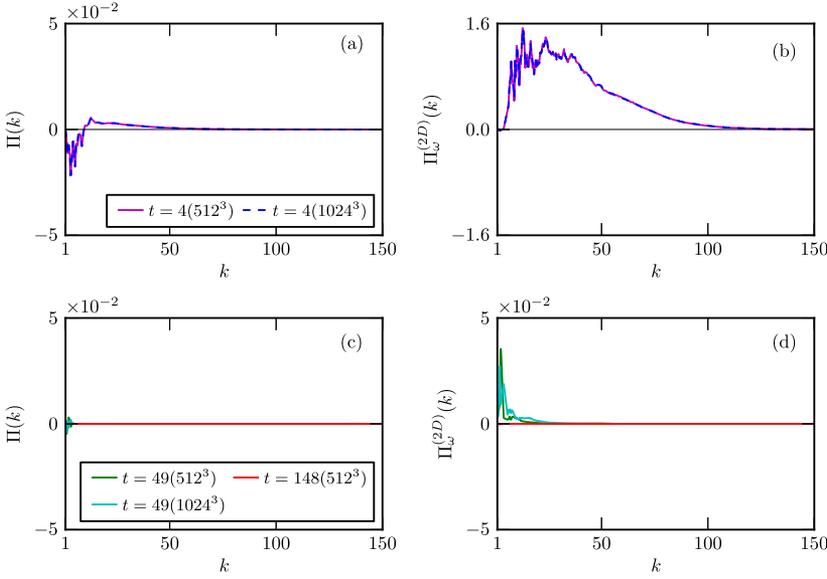


FIG. 4. Kinetic energy flux $\Pi(k)$ and enstrophy flux $\Pi_{\omega}^{(2D)}(k)$ for rapidly rotating decaying turbulence. (a) $\Pi(k)$ of 3D velocity field for 512^3 (magenta) and 1024^3 (blue dashed curve) grids at $t = 4$. (b) $\Pi_{\omega}^{(2D)}(k)$ for the 2D cross section at $z = \pi$. Figure (b) has the same color convention as (a). (c) $\Pi(k)$ at $t = 49$ (green), and $t = 148$ (red) for the grid resolutions of 512^3 and at $t = 49$ (cyan) for 1024^3 grids. (d) $\Pi_{\omega}^{(2D)}(k)$ for 2D cross section at $z = \pi$. Figure (d) has the same color convention as (c). Note that $\Pi_{\omega}^{(2D)}(k)$ is significant, but $\Pi(k)$ is negligible.

represents the enstrophy transfer from mode $\omega_z(\mathbf{p})$ to mode $\omega_z(\mathbf{k})$ with mode $\mathbf{u}_{\perp}(\mathbf{q})$ acting as a mediator. Note that $\mathbf{u}_{\perp} = u_x \hat{x} + u_y \hat{y}$ and $\omega_z(\mathbf{k}) = [i\mathbf{k} \times \mathbf{u}(\mathbf{k})]_z$. For the cross section at $z = \pi$, Figs. 4(b) and 4(d) illustrate the plots $\Pi_{\omega}^{(2D)}(k)$ vs. k at $t = 4, 49, 148$ for the grid resolution of 512^3 and at $t = 4, 49$ for the grid resolution of 1024^3 . For 512^3 and 1024^3 grids at the same time, the energy and enstrophy fluxes are equal, which is consistent with the fact that our results are grid-independent.

The enstrophy flux is positive definite, but it is not constant in a significant wavenumber band, in contrast to 2D

hydrodynamic turbulence for which $\Pi_{\omega}^{(2D)}(k)$ is constant in the inertial range and then it decreases after $k = k_d$.⁷⁰ The steepening of $\Pi_{\omega}^{(2D)}(k)$ in strongly rotating turbulence is due to the viscous effects, as in Eq. (14), and due to energy transfer from \mathbf{u}_{\perp} to u_z , analogous to that in quasi-static MHD;^{59,79} this is in contrast to constant $\Pi_{\omega}^{(2D)}(k)$ in hydrodynamic two-dimensional turbulence for $k > k_f$.

Now let us focus on the time frames $t = 49, 148$ when the coherent columnar structures are well developed and strong. For these times, in Fig. 5, we plot $kE_{\omega}^{(2D)}(k)$ and $\Pi_{\omega}^{(2D)}(k)$ vs. k

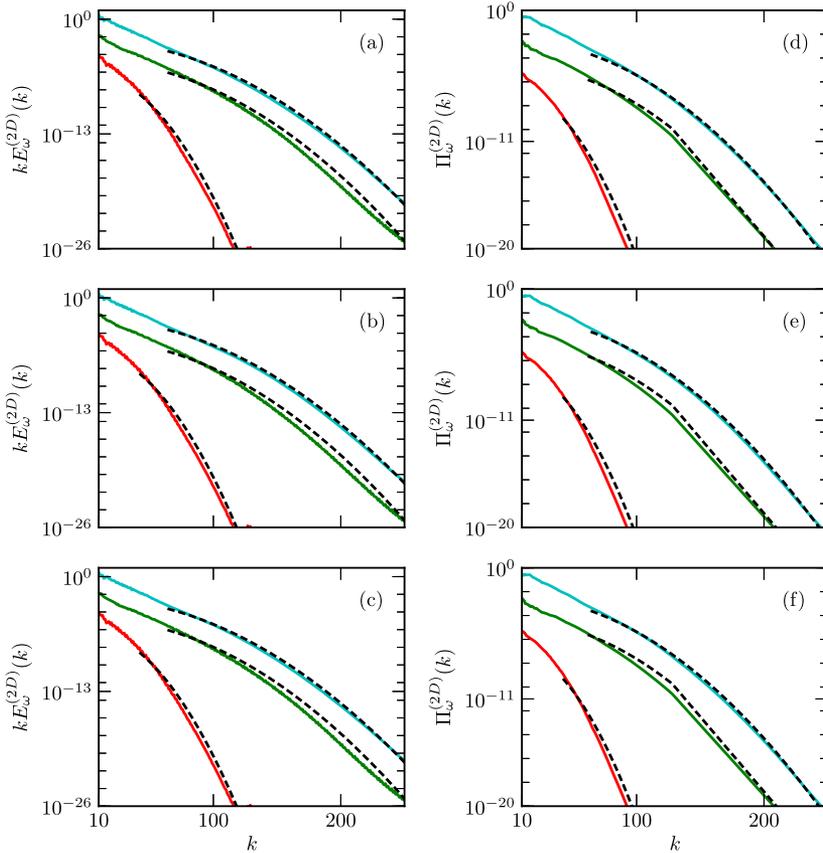


FIG. 5. For the rapidly rotating decaying simulation using \mathbf{u}_{\perp} at $z = \pi/2, \pi, 3\pi/2$. Two-dimensional enstrophy spectrum $kE_{\omega}^{(2D)}(k)$ for (a) $z = \pi/2$, (b) $z = \pi$, (c) $z = 3\pi/2$ at $t = 49$ (green) and $t = 148$ (red) for 512^3 grid simulation and at $t = 49$ (cyan) for the 1024^3 grid. The best fit curves using Eq. (13) with parameters of Table III are shown as black dashed curves. [(d)–(f)] The corresponding plots of the 2D enstrophy flux $\Pi_{\omega}^{(2D)}(k)$ following the same color convention as above. The best fit curves however are of the form of Eq. (14). Since $E_{\omega}^{(2D)}(k)$ and $\Pi_{\omega}^{(2D)}(k)$ of 512^3 and 1024^3 grids almost overlap each other, we multiply $E_{\omega}^{(2D)}(k)$ and $\Pi_{\omega}^{(2D)}(k)$ of 1024^3 grid with a factor 1/100 to differentiate the two plots.

TABLE III. For the rapidly rotating decaying turbulence, we take the instantaneous 2D velocity field on the horizontal cross sections at $z = \pi/2, \pi,$ and $3\pi/2$. List of the enstrophy dissipation rate, ϵ_ω , the enstrophy dissipation wavenumber, k_d , and constant C . These parameters are listed in the table at time $t = 49$ for 512^3 and 1024^3 and $t = 148$ for 512^3 .

Grid	t	z	k_d	$\epsilon_\omega \times 10^5$	$C \times 10^2$
512^3	49	$\pi/2$	6.0	4.8	(2.73 ± 0.04)
512^3	49	π	6.0	4.9	(2.76 ± 0.04)
512^3	49	$3\pi/2$	6.0	4.9	(2.74 ± 0.04)
1024^3	49	$\pi/2$	6.6	8.4	(3.00 ± 0.02)
1024^3	49	π	6.6	8.7	(3.03 ± 0.02)
1024^3	49	$3\pi/2$	6.6	8.2	(2.98 ± 0.02)
512^3	148	$\pi/2$	2.7	0.04	(2.45 ± 0.05)
512^3	148	π	2.7	0.04	(2.51 ± 0.05)
512^3	148	$3\pi/2$	2.7	0.04	(2.48 ± 0.05)

in semi-log scale for $z = \pi/2, \pi,$ and $3\pi/2$ planes. In rotating turbulence, $\Pi_\omega^{(2D)}(k)$ starts to decrease at small k itself because the enstrophy dissipation wavenumber, k_d , is quite small (see Table III). We compute the enstrophy dissipation wavenumber k_d using Eq. (15). It is difficult to estimate ϵ_ω of Eq. (13) because we do not have a band of wavenumbers where $\Pi_\omega^{(2D)}(k)$ is constant. In this paper, we compute ϵ_ω using Eq. (14) by identifying the neck of the wavenumber range from where $\exp(-C(k/k_d)^2)$ spectrum starts. If the wavenumber at the neck is k_* , then using Eq. (14),

$$\epsilon_\omega = \Pi_\omega^{(2D)}(k_*) \exp(C(k_*/k_d)^2). \quad (26)$$

Incidentally we observe that k_* is approximately twice of Kolmogorov's wavenumber, k_η . We compute C using linear regression analysis. The values of k_d , ϵ_ω , and C are tabulated in Table III. In Fig. 5, we plot the best fits to the enstrophy flux as dashed black curves [see Eq. (14)]. We observe that the above equation describes the numerical data very well. Also note that $\Pi_\omega^{(2D)}(k)$ of 512^3 and 1024^3 grids almost overlap on each other. Hence, to contrast the two plots, we multiply $E(k)$ and $\Pi_\omega^{(2D)}(k)$ of 1024^3 grid with a factor 1/100 to differentiate the two plots.

Motivated by the above observations, we model the enstrophy spectrum $E_\omega^{(2D)}(k)$ using Eq. (13). As shown in Fig. 5, we observe the numerical results to be in very good agreement with the model of Eq. (13). The red and the green

curves in Fig. 5 represent $kE_\omega^{(2D)}(k)$ and $\Pi_\omega^{(2D)}(k)$ for 512^3 resolution at $t = 49$ and $t = 148$, respectively, while the cyan curve represents the corresponding plots for 1024^3 resolution at $t = 49$. Also note that as expected, $E_\omega^{(2D)}(k)$ and $\Pi_\omega^{(2D)}(k)$ decrease with time due to their decaying nature. Thus we claim that the enstrophy spectrum and flux for the rapidly rotating decaying turbulence are described by Eqs. (13) and (14), respectively.

We now model the energy spectrum for the full flow. Using $\omega = \nabla \times \mathbf{u}$ and Eq. (13), we deduce that

$$\begin{aligned} E^{(2D)}(k) &= \frac{1}{k^2} E_\omega^{(2D)}(k) \\ &= C \epsilon_\omega^{2/3} k^{-3} \exp(-C(k/k_d)^2). \end{aligned} \quad (27)$$

For the cube (3D flow),

$$E_\perp(k) = E_x(k) + E_y(k). \quad (28)$$

Though $E_\perp(k)$ is not strictly equal to $\sum_{k_z} E^{(2D)}(k)$, we observe that $E_\perp(k)$ has similar scaling as $E^{(2D)}(k)$, i.e.,

$$E_\perp(k) = C'' \epsilon_\omega^{2/3} k^{-3} \exp(-C(k/k_d)^2), \quad (29)$$

where C'' is another constant. This is evident from Fig. 6(a), in which we plot $k^3 E_\perp(k)$ vs. k and the best fit curves as dashed black lines, $C \epsilon_\omega^{2/3} \exp(-C(k/k_d)^2)$, at $t = 49, 148$ for 512^3 and $t = 49$ for 1024^3 grid resolution. In the figure, the red and green curves represent the numerical spectra for the 512^3 resolution (at $t = 49, 148$), while the cyan curve is for the 1024^3 resolution (at $t = 49$). We observe that $C'' \approx C$ because $E_z(k) \ll E_\perp(k)$. In the plots of Fig. 6 too, at $t = 49$, the energy spectra of 512^3 and 1024^3 grids almost overlap with each other. Therefore, we multiply $E(k)$ of 1024^3 grids with 1/100 to contrast the two plots.

In rapidly rotating flows, the velocity component along the z direction is strongly suppressed, hence $E_z(k) \ll E_\perp(k)$. This observation is supported by the plot of the anisotropy parameter of Fig. 1(g). Hence

$$\begin{aligned} E(k) &= E_x(k) + E_y(k) + E_z(k) \approx E_\perp(k) \\ &\approx C \epsilon_\omega^{2/3} k^{-3} \exp(-C(k/k_d)^2). \end{aligned} \quad (30)$$

In Fig. 6(b), we plot $k^3 E(k)$ following the same convention as Fig. 6(a). We observe that Eq. (30) fits with the numerical data quite well with the same C as of Fig. 5. Thus, we claim that $E(k) \approx C \epsilon_\omega^{2/3} k^{-3} \exp(-C(k/k_d)^2)$ for the rapidly rotating decaying turbulence.

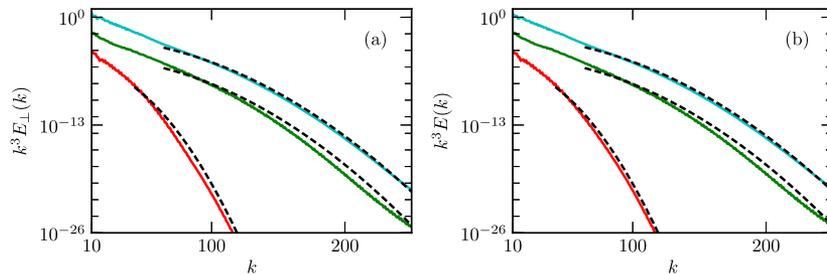


FIG. 6. For the rapidly rotating decaying turbulence, (a) normalized 3D energy spectrum $k^3 E_\perp(k)$ vs. k at $t = 49$ (green), 148 (red) for 512^3 grid simulation and at $t = 49$ (cyan) for 1024^3 grid. The best fit curves following Eq. (30) with ϵ_ω and C of Table III are plotted as black dashed curves. (b) The corresponding plots of $k^3 E(k)$ where $E(k) = E_\perp(k) + E_\parallel(k)$. Since $E(k)$ of 512^3 and 1024^3 grids almost overlap on each other, we multiply $E(k)$ of 1024^3 grid with a factor 1/100 to differentiate the two plots.

As described in Sec. I, for rapidly rotating turbulence, many researchers have reported dual spectrum for $E(k)$ with larger wavenumbers exhibiting Kolmogorov's spectrum. There is no unanimity on the spectral index for low wavenumber modes ($k < k_\Omega$) with researchers^{42–47,49–51} predicting the spectral exponents as -2 , -3 , $-5/3$. In fact, Morize *et al.*³³ argue that the spectrum steepens with the increase of rotation speed. For details refer to Sec. I. We did attempt to fit a power law ($E(k) \sim k^{-\alpha}$) with our data and observed that the spectral indices are -8 or lower for very small range of low wavenumbers. These observations yield stronger confidence in the model of Eqs. (13) and (14) that evidently spans over a much longer range of wavenumbers. We also remark that our model for $E(k)$ [Eq. (30)] performs better than those of Pao and Kraichnan, which are described by Eqs. (7) and (17), respectively. Note that the enstrophy and energy fluxes in Kraichnan's model are zero. See Appendix C for details.

We conclude in Sec. VI.

VI. CONCLUSIONS AND DISCUSSION

In this paper, we have performed numerical simulations of rapidly rotating decaying turbulence. In the asymptotic regime, the Rossby number of the flow is quite small, and the Reynolds number is quite large. Strong rotation causes quasi-two-dimensionalization of the flow and formation of large coherent columnar structures. Most of the kinetic energy is concentrated in these structures, and the fluctuations in the inertial and the dissipative ranges have a very small amount of energy leading to small Reynolds number for these fluctuations. Thus, rapidly rotating decaying turbulence has strong flow structures embedded in a sea of fluctuations of small magnitudes.

We have shown that the columnar structures are formed due to strong inverse cascade of energy. The vortex columns are quasi-2D with $u_z \ll u_\perp$, so we study the 2D energy and enstrophy spectra and fluxes of \mathbf{u}_\perp for various horizontal cross sections. We observe that the kinetic energy flux is quite small for $k > 8$, but the enstrophy flux, $\Pi_\omega^{(2D)}$ is significant in this regime. Furthermore, we deduce the expression for the enstrophy spectrum and the enstrophy flux for such flows as Eqs. (13) and (14), respectively. Since $u_z \ll u_\perp$, $E(k) \approx E_\omega^{(2D)}(k)/k^2$, thus $E(k) \sim C\epsilon_\omega^{2/3}k^{-3} \exp(-C(k/k_d)^2)$.

The anisotropy induced by intense rotation has strong similarities with those induced by strong mean magnetic field in MHD and quasi-static (QS) MHD turbulence.^{58,59} The flows in the rotating turbulence, as well as in MHD and QS MHD turbulence with strong mean magnetic field, are quasi-2D with $u_\perp \gg u_\parallel$. Also, the spectra of QS MHD and rotating turbulence for large k have exponential behaviour.⁵⁹ Further quantification of anisotropy in these systems would yield interesting insights.

The results presented in the paper indicate that the strong rotation induces strong vortical structures. Note that our simulations have very small Rossby number in comparison to the earlier simulations. Weak rotation (Rossby numbers of order 1) is likely to yield energy spectrum with power laws ($E(k) \sim k^{-\alpha}$). It will be interesting to make a comprehensive

study of variations of turbulence properties with the variation of Ro. However such a study is beyond the scope of this paper. Note that forced rotating turbulence yields different energy spectrum, as shown by earlier researchers.^{42–47,49–51} Also, kinetic helicity plays a major role in rotating flows. Thus, a detailed comparative study of the present work on strongly rotating turbulence with those on forced and helical rotating turbulence is in order; these studies will be carried out in future.

Many astrophysical, geophysical, and engineering flows involve strong rotation. For example,⁸⁰ Earth's outer core has $Ro \approx 10^{-6}$, and solar convection has $Ro \approx 10^{-2}$. The results discussed in this paper may be relevant to such systems. Note however that such systems typically involve thermal convection or magnetoconvection, hence they are more complex than what is discussed here.

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APPENDIX A: FINITE-SIZE EFFECTS

To verify the finite size scaling, we performed numerical simulations of rotating turbulence in two boxes of sizes $(2\pi)^3$ and $(4\pi)^3$. For the initial condition of these runs, we took our 512³ simulation data at $t = 98$. In the $(4\pi)^3$ box, $\mathbf{u}(\mathbf{x})$ outside the $(2\pi)^3$ box (central region) was set to zero at $t = 0$. The simulation was carried out until $t_f = 104$. In Figs. 7(a) and 7(b), we exhibit density plots of the vorticity field at $z = \pi$, i.e., $\omega_z(x, y, z = \pi)$, at $t = 100$. We observe that the size of the vortex in the $(4\pi)^3$ box is twice compared to that in $(2\pi)^3$ box. Hence, the large-scale vortex is indeed due to nonlinear effects, and it is independent of the box size. We also compute the integral length scales L for the two boxes and observe them to be approximately 5.1 and 11.5, respectively. See Fig. 7(c) for an illustration. Clearly, the integral length scale for $(4\pi)^3$ box is approximately double of that of $(2\pi)^3$ box. Thornber⁸¹ studied the impact of domain size and statistical error in decaying turbulence. Our preliminary studies on strongly rotating turbulence appear to show that its results are somewhat immune to system size. That is, the Fourier modes and their interactions are independent of the box size.

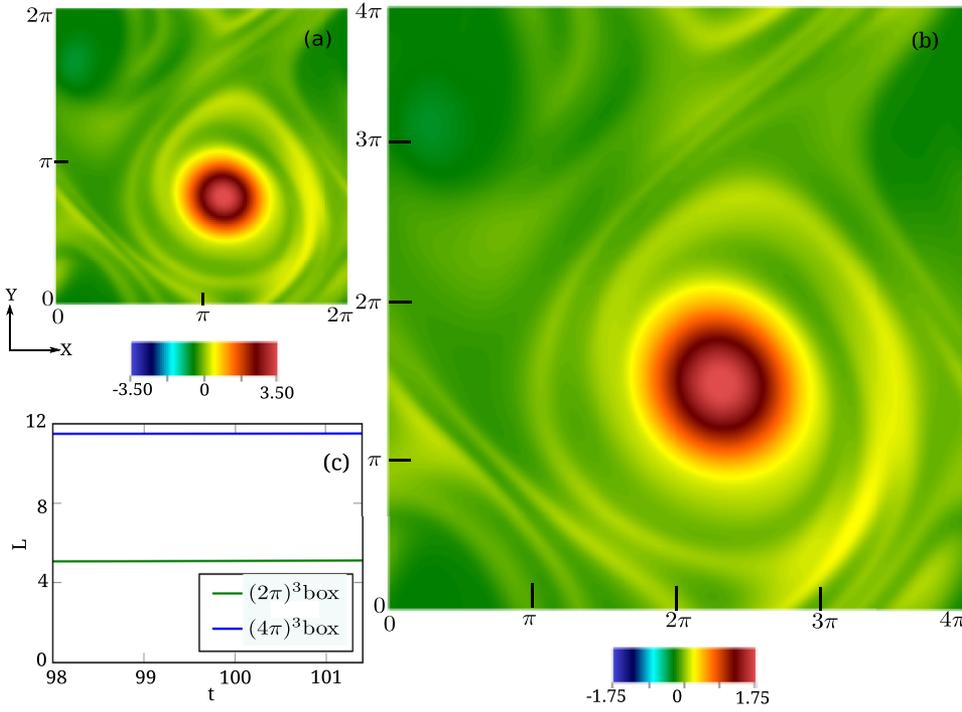


FIG. 7. For rapidly rotating turbulence, the density plots of ω_z at $z = \pi$ in boxes of sizes $(2\pi)^3$ (a) and $(4\pi)^3$ (b). For initial condition, we take the 512^2 data at $t = 49$ of Fig. 2. Clearly, the size of the vortex in $(4\pi)^3$ box is twice that of $(2\pi)^3$ box. (c) Plots of the time series of integral length scales L in $(2\pi)^3$ (green curve) and $(4\pi)^3$ (blue curve). Here $L_{4\pi} \approx 2L_{2\pi}$.

APPENDIX B: FOURIER MODES (1, 0, 0) AND (0, 1, 0) IN ROTATING TURBULENCE

In Sec. IV, we showed that $(1, 0, 0)$ and $(0, 1, 0)$ are the most dominant Fourier modes of strongly rotating turbulence. In this section, we discuss the 2D flow pattern when the system has only $(1, 0, 0)$ and $(0, 1, 0)$ Fourier modes. Here we choose the velocity field as

$$\mathbf{u}(x, y) = \hat{x} \sin y + \hat{y} \sin x. \quad (\text{B1})$$

In Fig. 8, we exhibit the velocity field along with the density plot of ω_z for the above field. The flow pattern is quite similar to that of Fig. 2(f), except that Fig. 8(a) shows cyclone-anticyclone symmetry, but Fig. 2(f) is asymmetric in cyclone-anticyclone pattern. Clearly, the cyclone-anticyclone asymmetry arises due to rotation.

In Appendix A, we described the results of numerical simulations of rapidly rotating turbulence. In Fig. 8(b), we exhibit the time series of energies of Fourier modes $(1, 0, 0)$, $(0, 1, 0)$, $(1, 1, 0)$, and $(-1, 1, 0)$, along with the total energy E_T . Clearly, the modes $(1, 0, 0)$, $(0, 1, 0)$ dominate other modes.

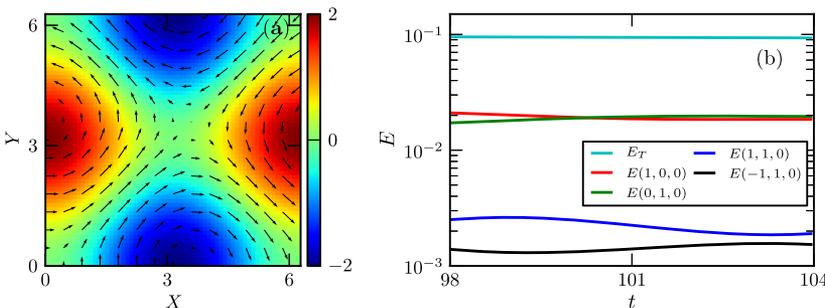


FIG. 8. (a) For the 2D velocity field $\mathbf{u}(x, y) = \hat{x} \sin y + \hat{y} \sin x$ that corresponds to the Fourier modes $(1, 0, 0)$ and $(0, 1, 0)$, the vector plot of the velocity field superposed with the density plot of the vorticity field ω_z . (b) For the rotating simulation of Appendix A, time series of the energies of the modes $(1, 0, 0)$, $(0, 1, 0)$, $(1, 1, 0)$, $(-1, 1, 0)$ and the total energy E_T .

APPENDIX C: COMPARISON BETWEEN OUR MODEL AND THAT OF PAO'S MODEL

In this paper, we argue that for strongly rotating decaying turbulence, Eq. (30) describes the energy spectrum in the intermediate and decaying range. This is the prediction of our model based on the variable enstrophy flux. However, it is important to compare it with the other models. In Fig. 9, we plot $E(k)$ for the numerical data of 1024^3 grid at $t = 49$. The plot also contains $E(k)$ predicted by Eq. (30) and that by Pao's model [see Eq. (7)]. Clearly, our model performs better than Pao's model.

We also remark that the predictions of Kraichnan's model given by Eq. (17) are not suitable for our simulations because $\text{Re} \gg 1$ for our flows. Kraichnan's model assumes that the nonlinearity is absent, which is not the case for our system. From Eq. (17) we deduce that the kinetic energy evolution in Kraichnan's model is

$$E(k, t) = E(k, 0) \exp(-2\nu k^2 t). \quad (\text{C1})$$

Thus, the evolution of $E(k)$ depends on the initial condition, and $E(k)$ quickly decays to zero.

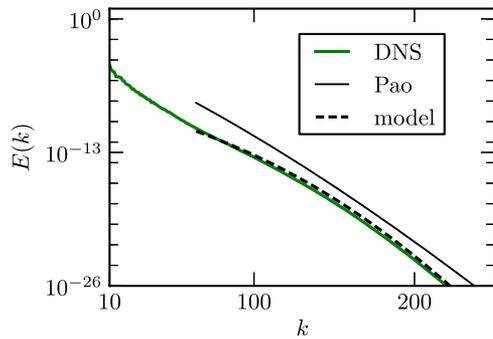


FIG. 9. Plot of $E(k)$ for the 3D velocity field of 1024^3 grid along with the predictions of our model and that of Pao. This plot corresponds to the green curve of Fig. 6.

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