Short circuits in the Corrsin–Obukhov cascade

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Experiments show that a blob of scalar released in the inertial range of scales of a turbulent medium is rapidly converted into a set of disjointed sheets whose spectrum exhibits a k^{-1} shape for wave numbers larger than the injection wave number $1/d$. The sheets diffusive uniformization onsets at a time $t_s=(d/u)\mathrm{f}(Sc)$ with $\mathrm{f}(Sc)\sim\ln(Sc)$ function of the injection time of the blob in the medium d/u, of the Schmidt number Sc, independently of the Reynolds number. This time is appreciably smaller than the time needed to cross the Kolmogorov cascade, which is "bypassed" by a strong and constant stretching rate acting at the injection scale. *© 2001 American Institute of Physics.*
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I. INTRODUCTION

It is a known fact that the time t_m necessary to achieve uniformity at the molecular level of a passive constituent released in a closed volume of linear scale L and stirred at a typical velocity $u(L)$ is

$$t_\mathrm{m}\sim\frac{L}{u(L)}, \tag{1}$$

with a prefactor depending on the detailed geometry of the mixing volume, the stirring device, and the nature of the dye being mixed. The above relationship is a rule of thumb in the engineering practice and is known to hold independently of the Reynolds number provided it is large enough, say larger than 10^4.[1,2] The apparent simplicity of (1) hides several difficulties. First, the fact that the mixing time t_m relates to the largest scales of motions in the flow may appear contradictory with the fact that the ultimate molecular uniformization of the scalar occurs at small scales compared to L. The mixing time (1) bears no trace, however, of straining motions and characteristic times associated with small scales. Consider, for example, a dye with a Schmidt number $Sc=\nu/D$ close to unity, where ν stands for the kinematic viscosity of the fluid and D for the molecular diffusion coefficient of the dye. The dye dissipation scale is, in that case, usually estimated as being identical to that of vorticity, namely the Kolmogorov scale $\eta=L\,Re^{-3/4}$, whose turnover time $t_\eta=L/u(L)\,Re^{-1/2}$ is appreciably smaller than t_m.

Second, it is clear that the intrinsic diffusive properties of the dye (i.e., its diffusion coefficient D) should appear in (1), the molecular uniformization requiring an infinite time in the limit $D\to 0$.

These difficulties are partly removed in the cascade description of turbulence stemming from Kolmogorov,[3] Obukhov,[4] Corrsin,[5] and Onsager.[6] Mixing of a passive scalar is in this frame represented as a succession of reduction of length processes, from the injection scale of the scalar blob to the scale where the thinning action of the stretch balances the spreading rate of the dye by molecular diffusion. These processes are assumed to be local in space (a blob of scale r is stretched at a rate corresponding to the scale r) and assumed to occur successively in time. From there, it is not difficult to estimate the time t_c for the blob to travel through the cascade from its injection size d to the scale where it is dissipated. Let γ_i be the stretching rate of scale r_i and t_i its associated persistence time; one obviously has $\gamma_i t_i=O(1)$ (see Ref. 7 for a discussion of the constant). After the first step of the cascade, the transverse size of the blob is $r_1=de^{-\gamma_1 t_1}$, after the second $r_2=r_1e^{-\gamma_2 t_2}$, and so on down to $r_n=r_{n-1}e^{-\gamma_n t_n}$, so that after n steps, the typical transverse size of the stretched blob is $r_n=de^{-n}$. Now, if the blob initially has a size lying in the inertial range of scales (that is, $\eta\ll d\ll L$), the time necessary to cross each step is $t_i\sim\epsilon^{-1/3}r_i^{2/3}$, with $\epsilon\sim u(L)^3/L$ according to the standard estimate.[3–6]

The time t_K for the blob to reach the viscous scale of the flow η, where the stretching times t_i become scale independent and are all equal to $t_\eta=(\nu/\epsilon)^{1/2}=L/u(L)\;Re^{-1/2}$, is $t_\mathrm{K}=\Sigma_{i=1}^n t_i\approx\epsilon^{-1/3}d^{2/3}$ for $n\gg 1$. The time t_K is dominated by the injection time scale. Assume that the Schmidt number of the dye is larger than unity. The thin sheets of scalar at the Kolmogorov scale η further dissipate at the Batchelor scale[8] $\eta_B=\sqrt{Dt_\eta}=\eta Sc^{-1/2}$.

The time needed to travel from η to η_B is $t_B=\Sigma_{i=1}^m t_i=\Sigma_{i=1}^m t_\eta=mt_\eta$, with $\eta_B=\eta e^{-m}$, and thus $t_B=\frac{1}{2}t_\eta\ln Sc$.

The total time to reach uniformity thus defines as "cascade time" for the scalar

$$t_c=t_\mathrm{K}+t_\mathrm{B}=\epsilon^{-1/3}d^{2/3}+\tfrac{1}{2}t_\eta\ln(Sc). \tag{2}$$

The cascade picture, at least as it has been originally conceived, assumes stationarity, meaning that the decay time of the scalar variance in the flow $C'^2(t)$ is long compared to all the time scales t_i involved in the cascade. In this limit, the rate at which the scalar variance $C'^2(t)$ is injected is the rate χ at which the variance is transferred to smaller scales and also the rate at which it is dissipated $2D\overline{(\nabla C)^2}$. The fundamental property of this stationary, forward cascade is that the reasoning can be reformulated for any mean squared concentration increment $\overline{c^2}(r)=\langle(C(x+r)-C(x))^2\rangle$ averaged over a volume of size r (with $r<d$), $\overline{c^2}(r)/\chi$ representing the time needed to reach the dissipation scale, starting from r. This time is precisely given by (2) for the special case r

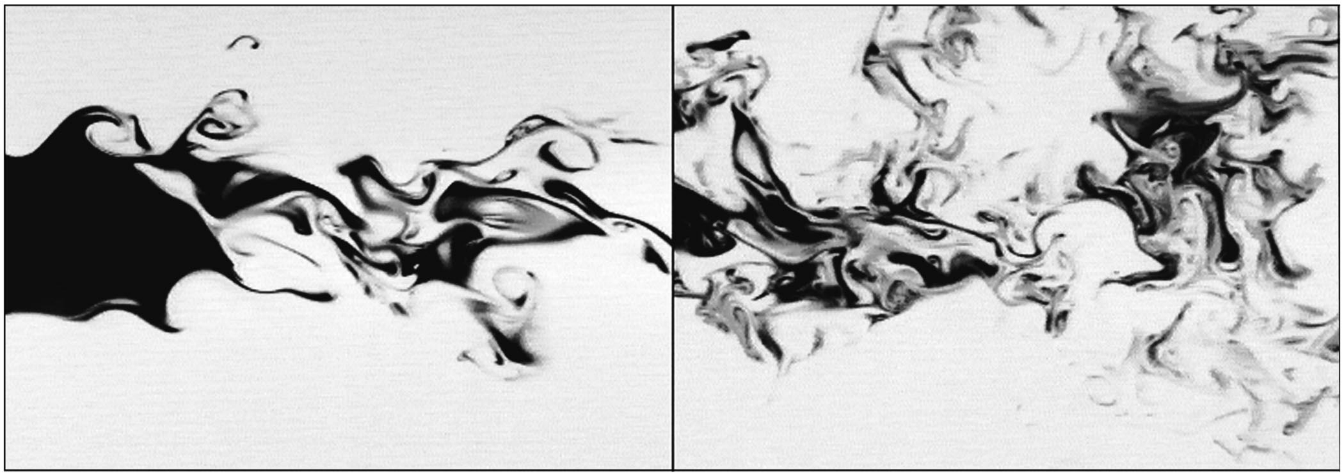


FIG. 1. Instantaneous planar cuts of the scalar field (disodium fluorescein in water, $Sc=2000$) downstream of the injection point illustrating how a scalar blob, initially compact and smooth, is progressively converted into disjointed sheets with broad fluctuations in thickness and concentration. $d=1$ cm. Left: the picture covers the region just downstream of the injection tube $0<x/d<4$. Right: further downstream $4<x/d<8$.

$=d$; that is, $\overline{c^2}(r)/\chi=t_c(r)=\epsilon^{-1/3}r^{2/3}+\frac{1}{2}t_\eta \ln(Sc)$. The Fourier spectrum $F(k)\sim\int e^{ikr}\overline{C^2}(r)\mathrm{d}r$ is thus

$$F(k)\sim\chi\Big(\epsilon^{-1/3}k^{-5/3}+\tfrac{1}{2}t_\eta k^{-1}\ln(Sc)\Big). \tag{3}$$

The relation (3) is a slightly altered reformulation of the Corrsin,[5,9] Obukhov,[4] Batchelor[8] theory. The spectrum interpolates from a $k^{-5/3}$ behavior in the inertial range (where t_c is dominated by $\epsilon^{-1/3}r^{2/3}$) to a k^{-1} dependence in the viscous range (where $\frac{1}{2}t_\eta \ln(Sc)$ dominates t_c). The crossover scale is $\eta(\ln(Sc))^{3/2}$, larger than η (this result is reminiscent of high Schmidt number experiments[10] where the spectrum has been observed to bend earlier than around $k\eta\approx1$). However, the spectrum of the passive scalar usually exhibits a dependence less steep than $-5/3$ in the inertial range,[10–14] the phenomenon being particularly clear in shear flows;[10–12] the existence of a -1 range at high wave numbers is not firmly established even at very large Schmidt numbers[10] so that it appears useful to re-examine the assumptions summed up above on a novel experimental basis.

II. EXPERIMENTS

The sequential cascade picture does not leave room for singular, efficient events which would directly connect injection scales to dissipative scales, thus participating majorly in the variance decay. We report in this paper on experiments suggesting that these events are the rule rather than the exception.

We follow the dilution of three types of scalars, disodium fluorescein ($Sc=2000$), temperature ($Sc=7$) in water, and temperature ($Sc=0.7$) in air, hastened by the high Reynolds number turbulence produced in the far field of a water jet. The experimental setup is described in Villermaux and Innocenti.[15] The scalar stream is injected continuously on the axis of the jet carrying the turbulence, through a tube whose diameter d is smaller than the local integral scale L (that is, $d/L=0.05$, 0.1, 0.16) and larger than the Kolmogorov scale $L\,Re^{-3/4}$ so that the scalar injection scale lies in the inertial range of scales. The injection point is located 30 diameters downstream of the exit of the main jet and the turbulence Reynolds number $Re=u'L/\nu$ is varied from 6000 to 45 000. u' denotes the local rms velocity ($u'/u\approx0.25$). The exit velocity of the injection tube is maintained constant and equal to the mean velocity of the main jet carrying the turbulence at the injection point in such a way that the tube behaves neither as a source, nor as a sink of momentum, in the mean.

Figure 1 shows how the scalar, initially confined in a compact region of space, is progressively fragmented into packets whose frontier gets more and more corrugated. These corrugations ultimately degenerate in thin elongated sheets, with broad fluctuations of transverse size and local concentration along the sheets. Fluctuations of concentrations and temperature are measured in the three different experiments on the axis of the jet by a fiber optics probe, a cold film thermometer, and a cold wire thermometer constructed at the laboratory, respectively. The resolution of the probes matches the Kolmogorov scale, in all cases; we relate space to time by $x=ut$ for x up to about one local integral scale, and we convert frequencies f into wave numbers k via Taylor hypothesis, i.e., $2\pi f=ku$. As can be seen in Fig. 2, the presence of the injecting tube does not perturb a region of more than two to three diameters d downstream of the injection point. The longitudinal velocity spectra recorded on the axis of the jet are essentially identical to their free stream analog (i.e., with the tube removed) at $x/d=4$.

In this configuration, Villermaux *et al.*[16] have measured the shape and the evolution as time progresses of the resulting one point concentration probability density function (PDF). The scalar concentration fluctuations PDF, $P(C/C_0)$, where C_0 denotes the injection concentration, rapidly exhibits a self-preserving shape, whose tail is an exponential with an argument increasing linearly in time t as

$$P\left(\frac{C}{C_0}\right)\sim\exp\left(-\frac{t}{t_s}\frac{C}{C_0}\right), \tag{4}$$

for $t\gg t_s$, with the mixing time t_s being found to be

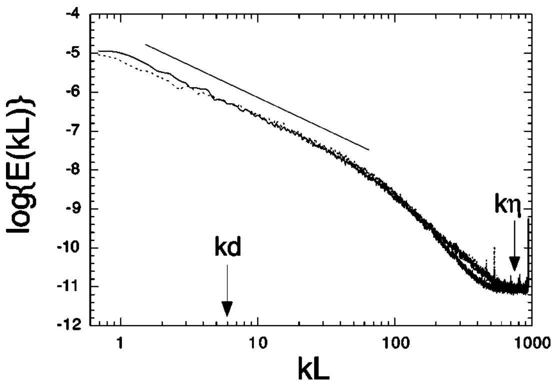


FIG. 2. Spectra of the longitudinal velocity fluctuations measured on the center line and 30 diameters downstream from the exit of the main jet. The local integral scale is $L=6$ cm. Dashed line: At $x/d=4$ downstream of the exit of the injection tube, with $d=1$ cm. Solid line: With the injection tube removed. The straight line has a $-5/3$ slope.

$$t_s=\frac{d}{u}f(Sc), \tag{5}$$

as shown in Fig. 3. The factor $f(Sc)$ is a slowly increasing function of the Schmidt number. A fit consistent with the data is $f(Sc)=0.12\ln(5Sc)$, but a power-law dependence of the form $f(Sc)\sim Sc^{0.2}$ is not inconsistent as well.

The mixing time t_s is both consistent with (1) and with the fact that the uniformization period increases with no bounds in the limit $D\to 0$.

The full PDF incorporates a Dirac delta at the origin $(C=0)$ accounting for the dispersion of the scalar support in the diluting medium so that the average concentration varies as $(ut/d)^{-2}$ (see Ref. 15). However, the leading contribution to the concentration fluctuations $C'^2(t)=\langle(C-\langle C\rangle)^2\rangle$ decay is given by the evolution of the PDF tail, and according to (4), the variance decays as $C'^2(t)\sim C_0^2(t/t_s)^{-3}$ for $t \gg t_s$. The use of three different injection diameters indicates that the characteristic time of the onset of the decay is actually proportional to d (Fig. 4), unlike t_{K}, or t_c, which scale like $d^{2/3}$ [see (2) and Ref. 9].

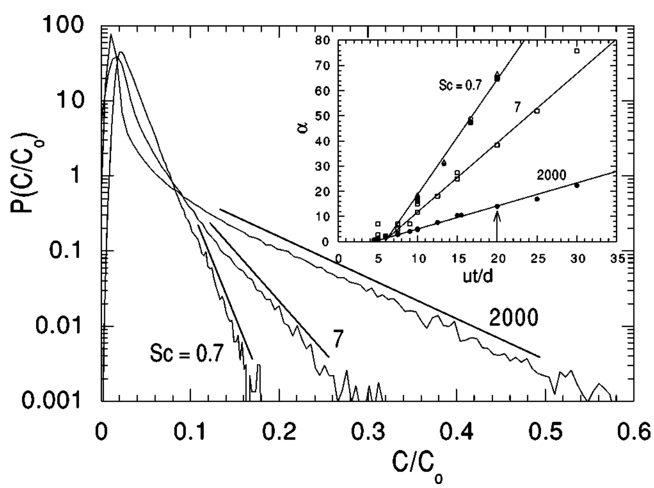


FIG. 3. Fluctuations PDFs, normalized by the initial concentration (temperature), recorded for the three Schmidt numbers 20 diameters d downstream from the injection point. The PDFs exhibit an exponential decay [Eq. (4)] of the form $P(C/C_0)\sim\exp(-\alpha C/C_0)$. Inset: The argument of the exponentials for different distances $x/d=ut/d$ and three Schmidt numbers is such that $\alpha=(ut/d)/[0.12\ln(5Sc)]$. ●: $Sc=2000$, $Re=6000$ and 12 000, $d/L=0.05$, 0.1, 0.6. □: $Sc=7$, $Re=6000$, $d/L=0.05$, 0.1, 0.16. ■: $Sc=0.7$, $d/L=0.08$, $Re=23\,000$; ○, △: $Sc=0.7$, $d/L=0.08$, $Re=45\,000$.

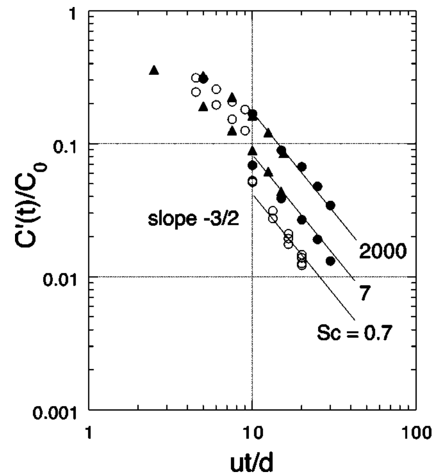


FIG. 4. Decay of the rms fluctuations $C'(t)$ downstream from the injection point for three different Schmidt numbers. ○, $d/L=0.16$; ▲, $d/L=0.1$; ●, $d/L=0.05$.

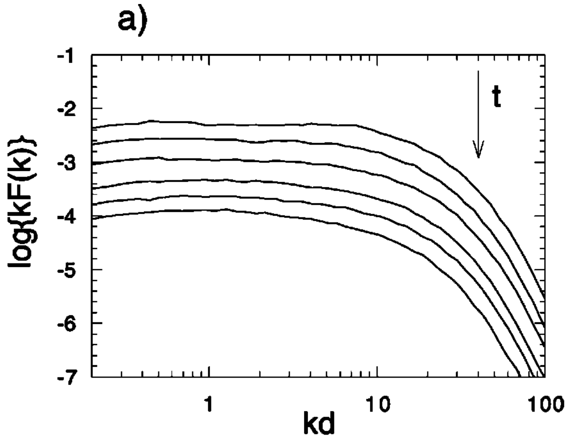
The power spectrum of the stationary scalar fluctuations signal recorded at different locations downstream of the injection point displays a clear k^{-1} range, between the injection scale d and the vicinity of the dissipation scale $dPe^{-1/2}$, when the latter can be resolved by the probe (the $Sc=0.7$ series of experiments) or down to the resolution scale of the probe otherwise (Fig. 5).[17] The spectra decay fairly self-similarly at large Schmidt number for increasing distances from the injection point, and tend to deviate slightly from the -1 slope at moderate, or order unity Schmidt numbers, by developing a more and more pronounced, but still slight curvature.

At the same location in the flow, the velocity spectrum displays a decay very close to $k^{-5/3}$ from $1/L$ to higher inertial range wave numbers (Fig. 2). As can be seen from Fig. 5, the ''k^{-1} shape'' is indeed the gross feature of the spectra. Their detailed shape is even not a pure power law, and their evolution reveals a more and more pronounced curvature as time elapses, particularly sensitive at low Schmidt numbers. It is not our aim to discuss this second (and probably important) aspect here; we emphasize that k^{-1}, or close-to-k^{-1}, is certainly different from $k^{-5/3}$, as expected from the cascade arguments summed up in Sec. I.

III. CASCADE BYPASS

The observations reported above suggest that the process of reduction of scale by sheet thinning in turbulent flows *does not* follow *step by step* the cascade prescribed by the pre-existing hierarchy of structures in the flow, even at large Reynolds numbers, but instead ''bypasses'' the cascade by a permanent, and efficient, stretching rate associated with the initial size of the blob and the maximal velocity in the flow.

We introduce below a mechanism which accounts for the continuous generation of scalar sheets by the small-scale (smaller than d, see Fig. 1) activity of the flow as an inter-



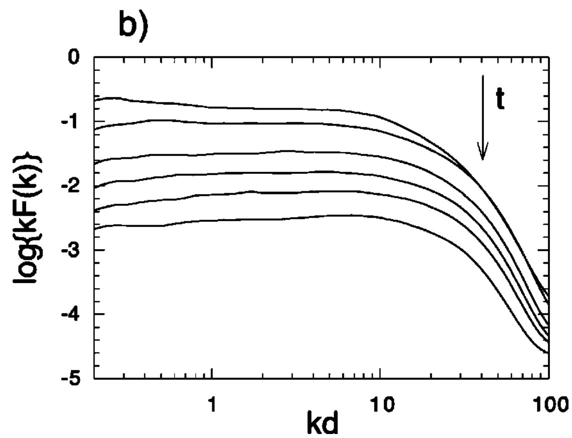


FIG. 5. Compensated power spectra $kF(k)$ of the scalar fluctuations at different successive times of the mixing process for the two extremes Schmidt numbers $(kd=d/r)$. Top: $Sc=0.7$, $d/L=0.16$, $L=20$ cm, $u=2$ m/s. The dissipation scale $dPe^{-1/2}$ corresponds here to a dimensionless wave number equal to $Pe^{1/2}\approx 55$, below the resolution wave number of the probe $k_c d=165$. In the direction of the arrow, $x=ut=10$, 13, 15, 18, 21 and 30 cm. Bottom: $Sc=2000$, $d/L=0.1$, $L=6$ cm, $u=0.4$ m/s. The probe resolution cutoff wave number is $k_c d=30$ and the dissipation scale in that case is not resolved $(Pe^{1/2}\approx 700)$. In the direction of the arrow, $x=ut$ $=1.5$, 3, 4.5, 6, 7.5 and 9 cm.

mediate step in the fragmentation process of a scalar blob in its route dissipation.

Let s_0 be the transverse size of a scalar sheet peeled off from the initial blob ($s_0<d$ *a priori*). The time $t_s(s_0)$ necessary to reduce the transverse thickness of this diffusing lamellae down to the equilibrium value $\sqrt{D/\gamma}$ where stretching-induced thinning balances diffusive spreading under the action of a permanent stretching rate γ is

$$t_s(s_0)\sim\frac{\gamma^{-1}}{2}\ln\left(\frac{s_0^2\gamma}{D}\right). \tag{6}$$

Equation (6) is reminiscent of the experimental mixing time (5) provided that $\gamma^{-1}/2\approx 0.12d/u$; that is, $\gamma\approx 4u/d$. The scale reduction process of the scalar sheets is mediated by a constant stretching rate, acting at the injection scale. The experiments reported in Fig. 3 also suggest that the mixing time is independent of the Reynolds number. This fact constrains the initial size s_0 of the scalar lamellae peeled off from the blob to be a decreasing function of the Reynolds number, precisely such that $s_0=\sqrt{5\nu/\gamma}\sim d/\sqrt{u'd/\nu}$. The size of the initial scalar lamellae scales like the Taylor scale, based on the injection diameter d.

Then, the dissipation scale $\sqrt{D/\gamma}=s_0 Sc^{-1/2}\sim d/\sqrt{u'd/D}=dPe^{-1/2}$ decreases like the square root of the Péclet number $Pe=u'd/D$ based on the injection diameter d. Note that according to the lines above, both the initial sheet thickness and the dissipation scale are expected to be independent of the turbulence integral scale L and are both proportional to the injection diameter d.

The ratio of the mixing time t_s given by (5) to the time t_c needed to travel through the ''Kolmogorov cascade'' (2), i.e.,

$$\frac{t_s}{t_c}=\frac{0.12\frac{d}{u}\ln(5Sc)}{\epsilon^{-1/3}d^{2/3}+\frac{L}{u}Re^{-1/2}\ln(Sc)}$$
$$\approx 0.12\frac{u'}{u}\left(\frac{d}{L}\right)^{1/3}\ln(5Sc), \tag{7}$$

is notably smaller than unity (of the order of 0.1 in the present experiments), even for large Schmidt numbers. From there originates the bypass.

The time t_s is the time needed to connect the initial peeled-off structures of size s_0 (these are produced in a very short time of the order of $\sqrt{u'd/\nu}$ times smaller than t_s), with the dissipation scale $dPe^{-1/2}$. The scalar field is indeed found to be composed mainly of thin sheets, although larger blobs do persist for all times. These larger scales originate in an incomplete peel off, or result from a subsequent diffusive reamalgamation of disjointed sheets.

For scales r larger than s_0 in the direction towards the inertial range of scales, up to d which is *a priori* the largest initial transverse size of a scalar lamellae, the sheet formation time $t_s(r)$ is somewhat larger than $t_s(s_0)$, within a logarithmic correction

$$t_s(r)\sim\gamma^{-1}\ln\left(Pe^{1/2}\frac{r}{d}\right). \tag{8}$$

Equations (6) and (8) are equivalent. Equation (8) is the version of (6) holding for *a priori* any scale $r<d$, thus formalizing the idea of the bypass for any scale up to the injection scale of the scalar d and therefore allowing us to make a prediction on the shape of the spectrum. Indeed, as soon as the global variance decay rate $-[1/C'^2(t)][dC'^2(t)/dt]$ $\sim 1/t$ is smaller than $1/t_s(r)$ for $Pe^{-1/2}<r/d<1$, the quasi-equilibrium limit for any mean squared concentration increment $\overline{c^2}(r)$ is

$$\frac{\overline{c^2}(r)}{t_s(r)}\sim\frac{C'^2(t)}{t_s(d)}, \tag{9}$$

with $t_s(d)$ being the largest sheet formation time of the few structures whose size has remained of order d at time t. The Fourier spectrum derives from (8) and (9) as

$$F(k)\sim\frac{C'^2(t)}{\ln(Pe^{1/2})}k^{-1}, \tag{10}$$

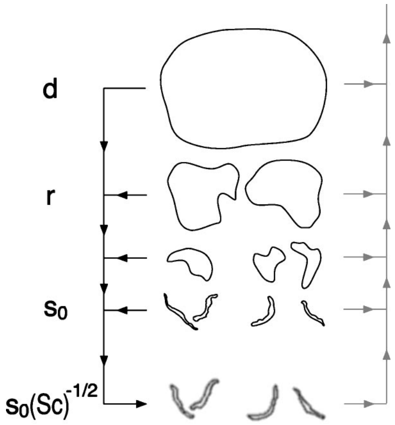


FIG. 6. Sketch of the route to dissipation by the ''cascade bypass.'' All scales are directly connected to the dissipation scale $s_0 Sc^{-1/2} = dPe^{-1/2}$ via the injection scale stretching rate $\gamma \sim u/d$. Possible diffusive reamalgamation of the sheets produces larger scales.

which holds for all wave numbers between the injection wave number $1/d$ and the dissipation wave number $Pe^{1/2}/d$.

In this picture, sketched in Fig. 6, the pre-existing hierarchy of scales in the flow is completely ignored, and this bypass also holds when a developed, nonintermittent spectrum prevails in the underlying velocity field, as for the present experiments. We do not interpret the intermittent (k^{-1}) character of the scalar spectrum as reflecting intermittency in the velocity spectrum.[18] On the contrary, these experiments rather suggest that the scalar distribution and the underlying displacement field are very loosely coupled to each other and that the mixing properties of a turbulent flow are very little sensitive to its detailed structure.[19,20]

The spectrum given in (10) decays in amplitude for increasing Reynolds number and Schmidt number ($F(k) \sim [C'^2(t)/\ln(Pe^{1/2})]k^{-1}$), but its integral $\int_{1/d}^{Pe^{1/2}/d} F(k)dk = C'^2(t)$ is actually conserved, independent of the Reynolds number, and with no spurious Schmidt number dependent factor other than that originally contained in $C'^2(t)$, which itself decays faster in time than the upper bound of the infinite Schmidt number limit.[21]

IV. CONCLUSION

We have presented a purposely designed transient experiment consisting of following the dilution of an initially segregated scalar blob in a sustained turbulent medium. The kinetics and spectral features of the mixture have been shown to be at odds with the sequential Corrsin–Obukhov cascade picture of scalar advection in turbulent flows.

The starting point of this work was exposed in Sec. I where it was recalled that the sequential Corrsin–Obukhov cascade picture has a direct consequence on (i) mixing times, and therefore (ii) on the spectral content of the mixture. This picture is confronted by an experiment designed to test it and the facts are presented in Sec. II: the mixing time of a blob of size d scales as $t_s \sim d/u' \ln(5Sc)$, as opposed to $L/u'(d/L)^{2/3}$ in the Corrsin–Obukhov vision and the spectrum of the mixture is k^{-1}, as opposed to $k^{-5/3}$, in the inertial range of scales, and for scales smaller than the injection scale d. Moreover, the mixing time t_s has been found to be *smaller* than the cascade time $L/u'(d/L)^{2/3}$. We have thus suggested that the process of mixing *does not* follow the sequential route expected from cascade arguments, but is on the contrary dominated by a more efficient (i.e., with a shorter time scale) mechanism, that we call bypass and for which we have imagined a possible scenario in Sec. III.

The k^{-1} shape is the ''natural'' spectral signature of an intermittent collection of disjointed sheets whose approach to uniformity is governed by a permanent stretching rate related to the injection scale of the scalar. We have interpreted this as a sign of nonlocal interactions bridging the injection scale of the blob d with the dissipation scale, which we anticipated to be given by $dPe^{-1/2}$, and we have underlined the spectral consequences of this point of view.

The permanence of the injection features of the scalar in the flow is manifested by the Reynolds number independence of the mixing time $t_s \sim d/u' \ln(5Sc)$ and by the fact that the scalar dissipation scale depends solely on its initial segregated scale d, and on its molecular properties through the Péclet number $u'd/D$.

Molecular viscosity does, however, play a role. The ''cascade bypass'' scenario we have imagined in this paper involves, as an intermediate step, the formation of scalar lamellae peeled off from the initial segregated blob, whose size s_0 scales like the Taylor scale based on d as $s_0 \sim d/\sqrt{u'd/\nu}$. These structures originate from the destabilization of internal shear layers consecutive to the relative motion of the blob with its surrounding environment. The destabilization of these shear layers is made possible as soon as the Reynolds number based on their own thickness $u's_0/\nu = \sqrt{u'd/\nu}$ is larger than about 150,[22,23] otherwise the instability is smoothed out by viscous spreading. The injection Reynolds number $u'd/\nu$ has thus to be large enough, that is typically larger than 10^4, for this ''mixing transition'' to occur.[2] Provided this condition is fulfilled, that is, provided the Reynolds number is *high enough*, the mixing time becomes *independent* of the Reynolds number as we have explained in Sec. II.

The persistence of the injection conditions of the scalar goes even further. It not only determines the global decay rate of the inhomogeneities through the mixing time, but also the whole statistics of the concentration distribution, as shown in Fig. 2. This observation has already been made in shear flows dominated by a mean scalar gradient[24,25] where the skewness of the temperature (concentration) derivative parallel to the mean gradient was found to be of order unity, and very weakly dependent on the Reynolds number.

The disjointed scalar sheet distribution resulting from a punctual injection in a flow is the ''quantum'' of scalar mixing. The interaction between several distributed sources and the amalgamation of the sheets as they diffuse, which occurs all the more early when the Schmidt number is low (and is probably a phenomenon sensitive to the Reynolds number) is the next ingredient to examine for understanding how the scalar field finally adapts to the hierarchical structure of the stirring field in the uniform, well-mixed limit.[10,12,26]

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