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Key Points:

- Mean evaporation rate was shown to scale with the 3/4 power law of the friction velocity and the square-root of molecular diffusivity
- The scaling was derived from surface renewal theory with ad hoc specification of contact probability between eddies and the surface
- A new approach that links these scaling results with the universal structure of the turbulent cascade is derived

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A Kolmogorov-Brutsaert structure function model for evaporation into a turbulent atmosphere

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Abstract In 1965, Brutsaert proposed a model that predicted mean evaporation rate $\bar{E}$ from rough surfaces to scale with the 3/4 power law of the friction velocity (u_*) and the square-root of molecular diffusivity (D_m) for water vapor. In arriving at these results, a number of assumptions were made regarding the surface renewal rate describing the contact durations between eddies and the evaporating surface, the diffusional mass process from the surface into eddies, and the cascade of turbulent kinetic energy sustaining the eddy renewal process itself. The working hypothesis explored here is that $\bar{E} \sim \sqrt{D_m} u_*^{3/4}$ is a direct outcome of the Kolmogorov scaling for inertial subrange eddies modified to include viscous cutoff thereby bypassing the need for a surface renewal assumption. It is demonstrated that Brutsaert's model for $\bar{E}$ may be more general than its original derivation implied.

Plain Language Summary The movement of water vapor molecules from rough surfaces such as soils by eddies into the atmosphere is of primary significance to a plethora of applications including hydrological and meteorological forecasting, irrigation planning, energy partitioning, and subsequent growth of the atmospheric boundary layer, to name a few. In 1965, W. Brutsaert proposed a general model that links the movement of water vapor molecules from the surface to the molecular diffusivity of water vapor in air and the wind-induced shear stress at the surface. The derivation considered air packets sweeping down and coming in contact with the wet surface. When in contact with the surface, these packets become enriched with water vapor molecules during a finite contact duration, after which these packets are ejected from the surface. Brutsaert made key restrictive assumptions about the statistical properties of the contact duration of these packets with the surface to arrive at the final form of the evaporation equation. The work here demonstrates that the same result can be derived by assuming a turnover velocity of these air packets to follow a universal form based on a widely accepted theory of turbulent flows proposed by A. N. Kolmogorov.

1. Introduction

In his acceptance letter of the Rumford Medal in 1881, Gibbs declared that "One of the principal objects of theoretical research is to find the point of view from which the subject appears in the greatest simplicity." Guided by this quotation, the subject of evaporation into the atmosphere from rough surfaces by turbulence offered in *Brutsaert* [1965] is reexamined. The work in *Brutsaert* [1965] considered the evaporation rate from a rough surface characterized by a momentum surface roughness height z_o in terms of molecular diffusion transmitting water vapor into eddies experiencing random contact times of duration τ with the surface. To proceed further, *Brutsaert* [1965] assumed that the distributional properties of τ, labeled $p(\tau)$, are exponential as originally proposed by *Danckwerts* [1951]. The novel outcome of the *Brutsaert* [1965] renewal analysis is that the mean evaporation rate $\bar{E}$ scales with the 3/4 power law of the friction velocity u_* (instead of a linear scaling) and the square-root of molecular diffusivity (D_m) of water vapor. Furthermore, *Brutsaert* [1965] assumed that the mean turbulent kinetic energy (TKE) dissipation rate ϵ describing the transfer of energy along the turbulent energy cascade is external to the evaporating surface and given by the mechanical production of TKE. This assumption holds for stationary and planar-homogeneous turbulence with no mean vertical velocity and when the sum of transport terms by turbulence and by pressure redistribution can be ignored. With these assumptions, the TKE budget at $z = z_o$ reduces to

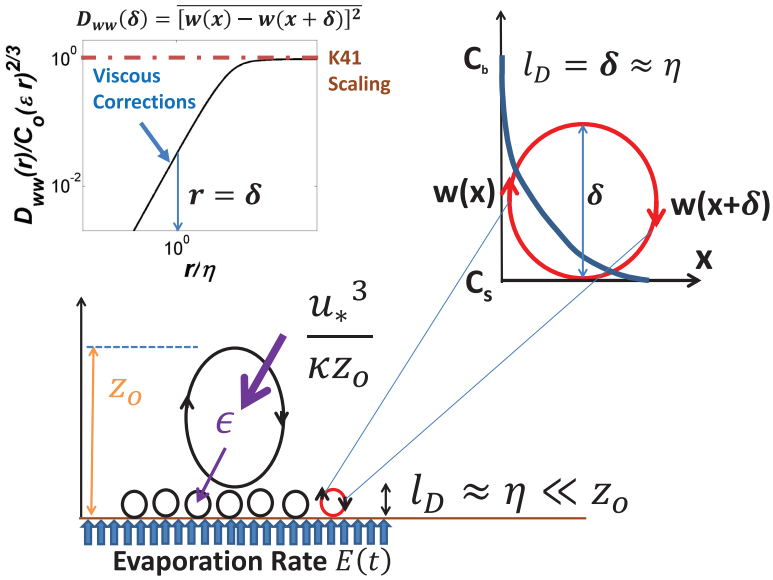

Figure 1. Schematic illustrating the interfacial transfer of water vapor into a turbulent atmosphere. A disturbed viscous sublayer of thickness (l_D) exists above a rough surface characterized by a momentum roughness length z_o that scales with the protrusion height of the roughness elements. Attached eddies are mechanically produced at $z = z_o$ and they transfer energy at rate ϵ to eddies within the disturbed viscous sublayer, where ultimately this energy is dissipated by viscosity. The figure here only shows the production and dissipation stages without showing the intermediate eddy sizes in-between for clarity. The mean water vapor concentration profile $\overline{C(z)}$ is characterized by a mean surface concentration C_s and then rapidly decreases with increasing height z within the disturbed viscous sublayer to a background value C_b. For $l_D < z < z_o$, $C(z)$ does not deviate appreciably from the background C_b. The turnover velocity of these eddies within the disturbed viscous sublayer is described by the structure function $D_{ww}(\delta) = \overline{[w(x+\delta) - w(x)]^2}$. The main novelty here is to show that when $D_{ww}(r)$ scales with K41, the final outcome in *Brutsaert* [1965] is recovered without requiring a surface renewal scheme.

$$\epsilon = u_*^2 \frac{u_*}{\kappa z_o}. \tag{1}$$

The complete analysis in *Brutsaert* [1965] yields

$$\bar{E} = A\sqrt{D_m}\frac{u_*^{3/4}}{(\nu\kappa z_o)^{1/4}}\overline{\Delta C} = A\left[\frac{D_m}{Sc\kappa z_o}\right]^{1/4} u_*^{3/4}\overline{\Delta C}, \tag{2}$$

where overbar is averaging to be defined later, A is a proportionality constant, ν is the kinematic viscosity of air, κ is the von Kármán constant, $Sc = \nu/D_m$ is the molecular Schmidt number, and $\overline{\Delta C}$ is the mean water vapor concentration difference between the evaporating surface and some distance away where the water vapor concentration does not vary appreciably with height z, shown in Figure 1. Equation (2) is supported by a large corpus of experiments as discussed in *Brutsaert* [1965] and later in *Brutsaert* [1975a]. This result spawned a number of studies on interfacial transfer of scalars [*Brutsaert*, 1975a, 1975b], evaporation from porous media at single and multiple pore scales [*Haghighi and Or*, 2013], bulk evaporation from bare soil surfaces [*Ishihara et al.*, 1992; *Or et al.*, 2013], as well as isotopic fractionation in hydrological applications [*Gat*, 1996; *Merlivat and Coantic*, 1975; *Hondzo*, 1998]. It also correctly foreshadowed the much discussed "universal" scaling of liquid transfer coefficients of sparingly soluble gases $(\sim (\nu\epsilon)^{1/4})$ in air-sea exchange studies [*Garbe et al.*, 2014; *Lorke and Peeters*, 2006; *Garbe et al.*, 2004; *Kermani and Shen*, 2009; *Soloviev and Schlüssel*, 1994; *Zappa et al.*, 2007]. This 1/4 scaling was first reported in the experiments by *Calderbank and Moo-Young* [1961] and its prediction is commonly attributed to another renewal model [*Lamont and Scott*, 1970], though as will be demonstrated here, it may be inferred from the model in *Brutsaert* [1965].

It is shown that the model outcome in *Brutsaert* [1965] can be succinctly derived using Kolmogorov's well-established inertial subrange scaling of the turbulent vertical velocity structure function [*Kolmogorov*, 1941] subject to viscous-cutoff corrections [*Batchelor*, 1951]. No arbitrary specification of random renewal or assumptions about $p(\tau)$ is necessary to arrive at equation (2). It is further demonstrated that even when a stochastic renewal process is enforced, the internal length scale of turbulence given by the Kolmogorov microscale $\eta = (\nu^3/\epsilon)^{1/4}$ imposes severe constraints on the mean contact time $\bar{\tau}$ making the final outcome

in *Brutsaert* [1965] insensitive to the specification of $p(\tau)$. Hence, the theoretical outcome in *Brutsaert* [1965] may be far more robust to assumptions invoked in its original derivation, the main conjecture explored here.

2. Theory

2.1. Background and Definitions

In the roughness sublayer (RSL) above a rough-wall boundary, a viscous sublayer whose thickness is much smaller than z_o is disturbed by turbulent flow characterized by a total stress τ_o and a squared shear or friction velocity (i.e., a total kinematic surface stress) $u_*^2 = \tau_o/\rho$, where ρ is the mean air density. The τ_o here is a combination of viscous and turbulent components and is assumed to be constant with distance z from the boundary at distances commensurate to or larger than z_o. Fully rough flow conditions require that $z_o u_*/\nu > 2$ to ensure a z_o larger than the viscous sublayer thickness as discussed in Figure 1. A sufficiently large separation between z_o and η ensures that eddies produced within the roughness sublayer have enough cascading steps to "forget" the eddy-generation mechanism as small scales are approached.

Transport of a scalar entity such as water vapor with an *Sc* near unity is conducted by molecular diffusion into the smallest eddies in contact with the surface disturbing the thin viscous sublayer as shown in Figure 1. While the primary focus here is on water vapor, the results derived apply to any scalar whose *Sc* is near unity, which is the case for many gases in the atmosphere [*Massman*, 1998].

2.2. The Surface Renewal (SR) Reviewed

To simplify the SR analysis in *Brutsaert* [1965], only the case where air parcels from a reasonably well-mixed region characterized by C_b are swept down and make full contact with the rough surface. The surface concentration is characterized by a near-constant C_s as shown in Figure 1. When air parcels arrive and are then in full contact with the surface, molecular diffusion governs the instantaneous water vapor flux $E(t)$ between the surface and these air parcels as shown in Figure 2. A model of maximum simplicity for the instantaneous

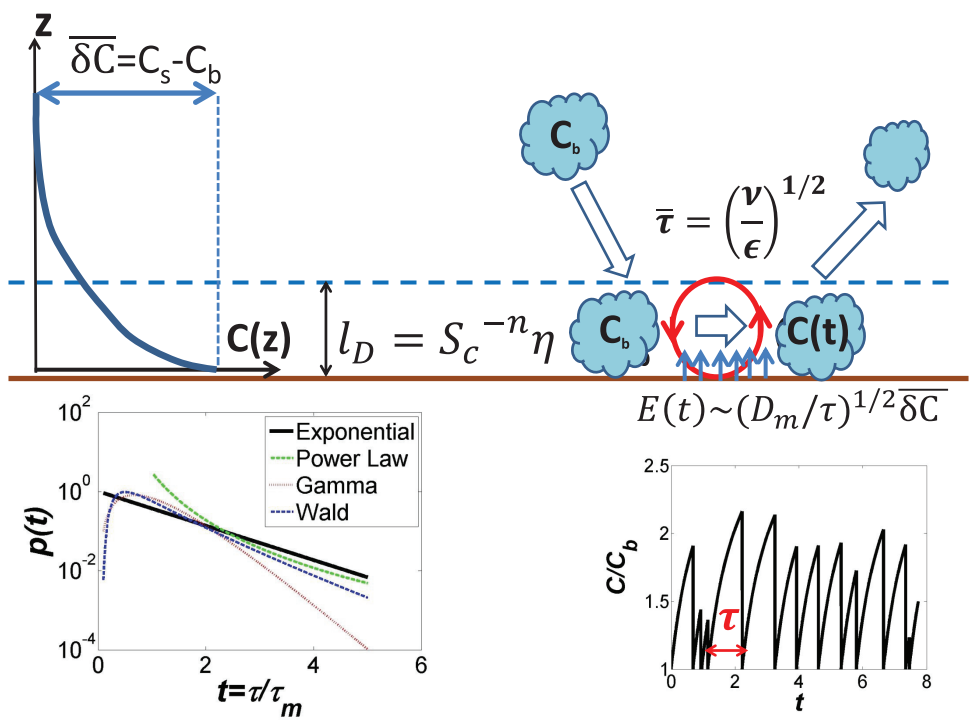

Figure 2. The surface renewal scheme assumes air parcels with water vapor concentration $C_b \approx \overline{C_b}$ first sweep to the surface, make contact with surface for duration $\tau > 0$ where their concentration increases due to an instantaneous $E(t)$, then are ejected out and followed by another sweeping event. The mean contact duration $\bar{\tau}$ is given by the Kolmogorov time scale. The concentration time pattern from such a renewal scheme is also shown for an exponential $p(\tau)$. The simplifying assumption here is that the renewal scheme does not alter the surface concentration appreciably. Other distributional properties of $p(\tau)$ assumed in the text are also featured and compared in the figure.

flux is $E \sim (D_m/\tau)^{1/2}\Delta C$ [*Brutsaert*, 1965; *Thomson and Silver*, 1972; *Soloviev*, 2007; *Kermani and Shen*, 2009] with $\Delta C = C_s(t) - C_b$ and $\tau > 0$ is, as before, the contact duration between air parcels and the surface where water vapor exchange occurs. A plausibility argument to this formulation may be offered by noting that Fickian diffusion from the surface into air parcels leads to $E(t,z) = -D_m dC/dz$. Replacing dC/dz by $\Delta C/l_D(t)$, where $l_D(t)$ is the distance separating C_s from C_b leads to $E(t) = [D_m/l_D(t)]\Delta C$. When the diffusion distance is set to $l_D(\tau) \sim \sqrt{D_m \tau}$, $E \sim (D_m/\tau)^{1/2}\Delta C$ is recovered. Turbulent eddies generate random τ characterized by $p(\tau)$ satisfying the normalizing condition $\int_0^\infty p(\tau)d\tau = 1$. Because replacing the air parcels in contact with the surface is associated with the SR process itself, $p(\tau)$ reflects the renewal occurrence after time τ. Infrared camera measurements of surface temperature fluctuations suggest the space-time renewal process can be reasonably captured by $p(\tau)$ [*Haghighi and Or*, 2015; *Garbe et al.*, 2004]. Hence, for each finite contact duration τ with the surface by air parcels, it is possible to formulate an instantaneous transfer coefficient $k_T(\tau) \sim \sqrt{D_m/\tau}$ [*Thomson and Silver*, 1972; *Soloviev*, 2007; *Kermani and Shen*, 2009] defined only when $\tau > 0$. The nonlinear $k_T(\tau) - \tau$ curve is assumed to be not altered by the renewal scheme. Only a variable τ drawn from $p(\tau)$ is to be used in the determination of k_T. To clarify the effect of $p(\tau)$ on SR and the mean evaporation rate, local interactions between $\Delta C(t)$ and the transfer coefficient are also ignored. This simplification amounts to setting $E(t) = k_T(t)\Delta C(t)$, $\overline{E(t)} \approx \overline{k_T(t)}\,\overline{\Delta C(t)}$. Such simplification may be unrealistic as surface concentration is likely to be impacted by the SR process and is not considered here though is considered in other studies [*Haghighi and Or*, 2013; *Garbe et al.*, 2004; *Soloviev and Schlüssel*, 1994]. Despite such simplification, the "naive" approach here recovers all the essential features of the *Brutsaert* [1965] model. It also offers a reasonable analytical foresight as the mean transfer coefficient controlling $\overline{E(t)}$ will only vary with the distributional properties of $p(\tau)$. Hence, only model differences in the selection of $p(\tau)$ will translate to differences in k_T. The shape of $p(\tau)$ can be readily linked to mean contact time and mean k_T when defined as

$$\bar{\tau} = \int_0^\infty \tau p(\tau)d\tau; \overline{k_T} = \int_0^\infty k_T(\tau)p(\tau)d\tau, \tag{3}$$

where overbar represents ensemble averaging and $\overline{k_T(\tau)} = \overline{E(t)}/\Delta C$ when assuming $\Delta C \approx \overline{\Delta C}$ as earlier noted. For $k_T(\tau) \sim \sqrt{D_m/\tau}$, $\overline{k_T(\tau)}$ may not be identical to $k_T(\bar{\tau})$ due to the nonlinear dependence of k_T on τ. However, when $\overline{k_T(\tau)} = k_T(\bar{\tau})$ and $\bar{\tau}$ is set to the Kolmogorov time scale $(\nu/\epsilon)^{1/2}$, then $k_T(\bar{\tau}) \sim \sqrt{D_m}(\epsilon/\nu)^{1/4}$. With $\epsilon = u_*^3(\kappa z_o)^{-1}$ as noted earlier, the two main results in *Brutsaert* [1965] are recovered as

$$\overline{E(t)} \sim \sqrt{D_m}\frac{u_*^{3/4}}{(\nu\kappa z_o)^{1/4}}\Delta C. \tag{4}$$

Exploring how different $p(\tau)$ representation impact the relation between $\overline{k_T(\tau)}$ and $\bar{\tau}$ when $k_T(\tau) \sim \sqrt{D_m/\tau}$ is examined. The motivation for this examination is to assess the robustness of $\overline{E(t)} \sim \sqrt{D_m}u_*^{3/4}\Delta C$ to assumptions made about the $p(\tau)$ model. Alternative models that solve the full diffusion equation can replace the simplified expression $k_T(\tau) \sim \sqrt{D_m/\tau}$ though they are not pursued further here. Various models for $p(\tau)$ are now considered and their shapes are compared in Figure 2. Exponential and gamma distributed $p(\tau)$ are routinely employed in SR analysis as discussed elsewhere [*Danckwerts*, 1951; *Perlmutter*, 1961; *Harriott*, 1962; *Bullin and Dukler*, 1972] with the gamma distribution receiving significant experimental support for heat and gas exchanges from soils [*Haghighi and Or*, 2013; *Or et al.*, 2013]. Other models have also been suggested or inferred from high-resolution infrared measurements [*Garbe et al.*, 2004] and direct numerical simulation runs [*Kermani and Shen*, 2009] for air-water interfaces. Power law distributions have been experimentally reported for bursting durations of sensible heat above bare soils [*Katul et al.*, 1994] and for momentum near smooth walls [*Rao et al.*, 1971] whereas Wald (or inverse-Gaussian) distributions naturally emerge when describing crossing statistics of Brownian air parcels with finite drift [*Katul et al.*, 2005]. Stretched exponential (i.e., power law distribution censored by an exponential cutoff) have been used in persistence studies of scalar concentration fluctuations in the roughness sublayer [*Cava and Katul*, 2009] and those distributions can be linked to $p(\tau)$ using approaches discussed elsewhere [*Laherrere and Sornette*, 1998]. Their overall features resemble Wald distributions and are not treated separately here. It is shown next that all these $p(\tau)$ models, characterized by two parameters a and b, do not differ appreciably in predicting $\overline{k_T(\tau)}$ when constrained by a mean contact duration commensurate with the Kolmogorov time microscale.

1. *Exponential*. When assuming $p(\tau)=a\exp(-b\tau)$ as used by *Brutsaert* [1965], then enforcing $\int_0^\infty p(\tau)d\tau=1$ leads to $a/b=1$ or $a=b$. Employing the definitions in equation (3), it directly follows that $\bar{\tau}=ab^{-2}=b^{-1}$ and $\overline{k_T}\sim\sqrt{\pi}D_m^{1/2}b^{1/2}$ or $\overline{k_T(\tau)}=k_T(\bar{\tau})$ with no additional simplifications. Upon setting $\bar{\tau}=b^{-1}=(\nu/\epsilon)^{1/2}$, it also follows that $\overline{k_T(\tau)}\sim D_m^{1/2}(\epsilon/\nu)^{1/4}$. With $\epsilon=u_*^3(\kappa z_o)^{-1}$, $\overline{E(t)}\sim\sqrt{D_m}u_*^{3/4}(\nu\kappa z_o)^{-1/4}\Delta C$ is recovered.
2. *Power law (Pareto or Zipf)*. When assuming $p(\tau)=ba^b/\tau^{b+1}$, then equation (3) results in $\bar{\tau}=(ab)/(b-1)$ for $\tau\geq a$ and $\overline{k_T}\sim D_m^{1/2}\,2b/[\sqrt{a}(1+2b)]$. If b is sufficiently large, the mean and mode of contact times converge to a and $\overline{k_T}\sim D_m^{1/2}\,1/[\sqrt{a}]$ (i.e., resulting in $\overline{k_T(\tau)}=k_T(\bar{\tau})$). Setting $\bar{\tau}=a=(\nu/\epsilon)^{1/2}$ as before recovers $\overline{k_T(\tau)}\sim D_m^{1/2}(\epsilon/\nu)^{1/4}$. With $\epsilon=u_*^3(\kappa z_o)^{-1}$, $\overline{E(t)}\sim\sqrt{D_m}u_*^{3/4}(\nu\kappa z_o)^{-1/4}\Delta C$ is again recovered.
3. *Gamma*. This is a common distribution in SR analysis and is given by $p(\tau)=[\tau^{a-1}/(\Gamma(\mathrm{a})\mathrm{b}^{\mathrm{a}})]\exp(-\tau/b)$, where $\Gamma(.)$ is the Gamma function. From equation (3), $\bar{\tau}=ab$ because $\Gamma(a+1)=a\Gamma(a)$. It also leads to $\overline{k_T}\sim D_m^{1/2}\Gamma(a-1/2)/\left[\sqrt{b}\,\Gamma(a)\right]$. For sufficiently large a, $\Gamma(a+\alpha)/[\Gamma(a)a^\alpha]\approx 1$, $\overline{k_T}\sim D_m^{1/2}/\sqrt{ba}$ and with the mean contact time $\bar{\tau}=ab=(\nu/\epsilon)^{1/2}$ (i.e., resulting in $\overline{k_T(\tau)}=k_T(\bar{\tau})$), $\overline{k_T(\tau)}\sim D_m^{1/2}(\epsilon/\nu)^{1/4}$. With $\epsilon=u_*^3(\kappa z_o)^{-1}$, $\overline{E(t)}\sim\sqrt{D_m}u_*^{3/4}(\nu\kappa z_o)^{-1/4}\Delta C$ is, once again, recovered. For small a, this argument fails though a sensitivity analysis shows that $a>3$ may be treated as sufficiently large for $\Gamma(a+\alpha)/[\Gamma(a)a^\alpha]\approx 1$.
4. *Inverse-Gaussian (or Wald)*. When assuming $p(\tau)=\sqrt{b/2\pi\tau^3}\exp\left[-b(\tau-a)^2/(2a^2\tau)\right]$, then equation (3) results in $\bar{\tau}=a$ and $\overline{k_T}\sim D_m^{1/2}\,a^{-1}\exp(b/a)\sqrt{2b/\pi}K_\alpha(s)$, where $K_\alpha(s)$ is a modified Bessel function of the second type of order α and $s=b/a$. For an arbitrary α,

$$K_\alpha(s)=\sqrt{\frac{\pi}{2s}}\exp(-s)\left[1-\frac{4\alpha^2-1}{8s}+\ldots\right], \tag{5}$$

with higher-order terms decaying rapidly with increasing s. For large s, the leading order term is $K_\alpha(s)\approx\sqrt{\pi/2}\,\exp(-s)$. This expression is only exact for $\alpha=1/2$. With this simplification for $K_\alpha(s)$, $\overline{k_T}\sim D_m^{1/2}a^{-1/2}$. Setting $\bar{\tau}=a=(\nu/\epsilon)^{1/2}$ once again recovers $\overline{k_T(\tau)}\sim D_m^{1/2}(\epsilon/\nu)^{1/4}$. With $\epsilon=u_*^3(\kappa z_o)^{-1}$, $\overline{E(t)}\sim\sqrt{D_m}u_*^{3/4}(\nu\kappa z_o)^{-1/4}\Delta C$ emerges.

To sum up, setting $\bar{\tau}=(\nu/\epsilon)^{1/2}$ leads to $\overline{k_T(\tau)}\sim D_m^{1/2}(\epsilon/\nu)^{1/4}$ for various $p(\tau)$ models suggesting that the precise details of the renewal scheme encoded in $p(\tau)$ are not as significant for the main result by *Brutsaert* [1965] to hold. The analysis here also confirms earlier conclusions [*Perlmutter*, 1961; *Koppel et al.*, 1966; *Seo and Lee*, 1988] that the precise shape of $p(\tau)$ may be less consequential to bulk transfer coefficients when using SR provided some constraints on mean contact durations are established.

3. Beyond Surface Renewal: A Structure Function Approach

The $E(t)$ is related to a mean concentration difference between the interface and the bulk layer situated at some δ from the surface as shown in Figure 1. For this setup, $\bar{E}=\omega_{T.O.}\overline{\Delta C}$. Here the transfer velocity may be interpreted as an effective turnover turbulent velocity $|\omega_{T.O.}|$ with magnitude $\sqrt{\omega_{T.O.}^2}$ for eddies of arbitrary size δ as shown in Figure 1. This exchange process requires that scalar quantities move coherently with $\omega_{T.O.}$ and do not appreciably diffuse in or out of eddies of size δ during a turnover. It will be shown later that this assumption requires Sc to be commensurate to or greater than unity. The simplest model is to relate $\omega_{T.O.}^2$ to ϵ and eddy size using a Kolmogorov inertial subrange scaling (or K41) for locally homogeneous and isotropic turbulence [*Kolmogorov*, 1941]. That is

$$\omega_{T.O.}^2=D_{ww}(r)=\overline{[w(x+\delta)-w(x)]^2}=C_o(\epsilon\delta)^{2/3}, \tag{6}$$

where C_o is the Kolmogorov constant for the vertical velocity, w is the turbulent vertical velocity, and x is an arbitrary position as shown in Figure 1. This conceptual approach successfully reproduced many features linking eddy sizes to bulk turbulent flow properties [*Gioia and Bombardelli*, 2002; *Gioia et al.*, 2010; *Katul et al.*, 2011; *Manes et al.*, 2012; *Konings et al.*, 2012]. This success is surprising as near solid or porous boundaries, the flow is not locally homogeneous and isotropic (e.g., eddies near smooth solid surfaces exhibit worm-like topology) and deviations from K41 scaling are to be expected. Notwithstanding this critique, invoking K41 scaling is a logical first step to be conducted as shown in prior studies flow over rough surfaces [*Gioia and Bombardelli*, 2002]. Setting δ to $l_D=(\nu^3/\epsilon)^{1/4}$, the bulk transfer velocity can be interpreted as an effective eddy turnover velocity resulting in $k_T=\sqrt{\omega_{T.O.}^2}\sim\sqrt{C_o}\,(\nu\epsilon)^{1/4}$. This outcome explains why such transfer velocity appears to be related to $(\nu\epsilon)^{1/4}$ across different types of surfaces or interfaces [*Zappa et al.*, 2007]. The quantity $(\nu\epsilon)^{1/4}=v_{Kol}$ is simply the Kolmogorov velocity defined by the Reynolds number

$v_{Kol}\eta/\nu=1$. Unsurprisingly, any dimensional consideration for k_L invoking only Kolmogorov variables (i.e., ν and ϵ) describing the "hand-shake" between the viscous sublayer and the overlying turbulent layer must yield v_{Kol} as the "master" variable.

With the TKE production assumed to be balancing dissipation again (i.e., $\epsilon=u_*^3(\kappa z_o)^{-1}$) and only accounting for K41 scaling, the structure function approach directly leads to

$$\bar{E} \sim \left[\frac{\nu}{\kappa z_o}\right]^{1/4} u_*^{3/4}\overline{\Delta C} \approx \left[\frac{D_m^2}{\nu\kappa z_o}\right]^{1/4} u_*^{3/4}\overline{\Delta C}. \tag{7}$$

Equations (2) (right-end equality) and (7) converge when Sc is set to unity (i.e., $\nu=D_m$) in equation (2). An amendment to equation (7) may also be derived by explicitly including viscous corrections to K41 scaling, which are expected to be significant for scales commensurate to the Kolmogorov microscale (the case here). Specifically, the viscous corrections to a K41 structure function are given by

$$\frac{\overline{[w(x+r)-w(x)]^2}}{C_o(\epsilon r)^{2/3}} = 1-\frac{1}{\zeta}\mathrm{Daw_F}(\zeta), \tag{8}$$

where $\zeta=\theta(r/\eta)^{2/3}$, $\theta=(10C_k')^{-1/2}$, $C_k'=2$, and the Dawson function is given by

$$\mathrm{Daw}_F(\zeta)=\exp(-\zeta^2)\int_0^{\zeta}\exp(p^2)dp \approx \zeta-\frac{2}{3}\zeta^3+\ldots \tag{9}$$

This expression presented in Figure 1 was derived from approximations to the von Kármán-Howarth equation [*Kármán and Howarth*, 1938] (VKH) without resorting to a constant skewness assumption [*Obukhov*, 1949; *Panchev*, 1971] across scales. It was shown elsewhere [*Katul et al.*, 2015] that this approximate solution to the VKH equation recovers several features about transitions from K41 scaling to viscous range, including the spectral bottleneck in the turbulent kinetic energy spectrum [*Hill*, 1978; *Herring et al.*, 1982; *Saddoughi and Veeravalli*, 1994; *Lamorgese et al.*, 2005; *Meyers and Meneveau*, 2008; *Ishihara et al.*, 2009; *Donzis and Sreenivasan*, 2010; *Frisch et al.*, 2013]. It also agrees with the Batchelor correction to the K41 scaling in the viscous subrange [*Batchelor*, 1951]. Setting $r=\delta=\eta$ and considering only a two-term expansion of the Dawson function does not alter the final outcome in equation (7). The expected viscous corrections to K41 scaling for $D_{ww}(r)$ here simply do not alter the $(\epsilon)^{1/4}$ scaling for k_T. Because Sc is near unity for many gases in the atmosphere [*Massman*, 1998], equation (7) is the sought result.

4. Discussion and Study Limitation

In arriving at the results in equation (2), a number of assumptions were made in the original derivation of *Brutsaert* [1965] regarding the surface renewal rate describing the contact durations between eddies and the evaporating surface, the diffusional mass transfer process from the surface into the eddies, and the cascade of turbulent kinetic energy sustaining the eddy renewal process itself from production to dissipation. The structure function approach proposed here recovers all the scaling results pertinent to equation (2) without any explicit consideration of (i) the renewal process and (ii) the precise diffusional mass transfer of water vapor from the surface into the eddies. Despite the maximum simplicity of this approach relative to the original derivation in *Brutsaert* [1965], the proposed structure function approach suffers from a number of limitations. To begin with, it is known that near roughness elements, production, and dissipation of TKE are not in balance and production may be up to 1.7 times larger than ϵ [*Pope*, 2000] even for very high Reynolds numbers and close to the surface at $zu_*/\nu \approx 10$ [*McColl et al.*, 2016]. However, this imbalance may not be as detrimental given the subunity exponent dependency of k_T on ϵ. Another assumption is that the TKE production rate scales with z_o, which ignores large eddy contributions near the surface as discussed elsewhere [*Katul and Manes*, 2014]. In the scalar transport calculations, it is assumed that K41 scaling with the usual viscous cutoff describes the structure function of the vertical velocity without much regard to the shape of the production regime. Near-wall measurements of the structure functions above smooth and rough walls suggest that K41 scaling may be plausible at $zu_*/\nu \approx 15$ (the limit of the measurements) for smooth wall [*Poggi et al.*, 2003] and immediately above rough surfaces composed of cylinder type roughness elements above rough walls [*Poggi et al.*, 2004]. Studies in the RSL above tall forests (porous roughness elements) confirm an extensive K41 scaling near the forest-atmosphere interface [*Katul et al.*, 1998] for

turbulent vertical velocity. Hence, adopting K41 scaling laws to the structure function of vertical velocity may not be "too far" off as an operational assumption for modeling $\bar{E}$.

Two interconnected results emerge from the structure function analysis. The nonlinear relation between k_T and u_* and the well-known dissimilarity between momentum and water vapor roughness heights. The above analysis makes clear that when $\epsilon \sim u_*^3(z_o)^{-1}$, the turnover velocity $|\omega_{T.O.}| \sim u_*(\delta/z_o)^{1/3}$. If $\delta \sim z_o$, then $k_T(=|\omega_{T.O.}|)$ must scale linearly with u_* not $u_*^{3/4}$. Any nonlinearity in the $k_T - u_*$ relation can then be attributed to a dependence of $(\delta/z_o)^{1/3}$ on u_*. The z_o is expected to scale with the mean protrusion height of roughness elements and arises here in the production of TKE term within the roughness sublayer. It can be interpreted as the length scale most effective at TKE generation in the RSL. The δ mainly reflects eddy sizes transporting water vapor from the surface after these eddies cascaded from their parent length scale z_o. These eddies are shown to scale with η and hence their $u_*^{-3/4}$ dependence. It directly follows that for a constant z_o, $(\delta/z_o)^{1/3}$ must scale with $u_*^{-1/4}$ and this result alone explains why k_T scales with $u_*^{3/4}$ instead of u_*. Moreover, a (δ/z_o) that varies with u_* also explains the origin of variations in water vapor roughness height z_{ov} with z_o and u_*, a topic that was considered later on in *Brutsaert* [1975b] and is not repeated here. For practical measurements and simulations of evaporation, it is the variation of z_{ov}/z_o with $u_* z_o/\nu$ that bundles much of the transfer processes in this interfacial sublayer covered by the formulations developed here.

The structure function derivation proposed here assumes $Sc \approx 1$, which is reasonable for several gases in the atmosphere, including water vapor, CO_2, CH_4, and air temperature. When Sc deviates appreciably from unity, a number of issues must be confronted that are beyond the immediate scope of this work. Few of them are highlighted for Sc values comparable to those reported in the atmosphere for many gases (0.6–0.7). The first issue is that l_D differs from the Kolmogorov microscale η when Sc is sufficiently different from unity as the effective eddy sizes in the vicinity of η may no longer reflect "scalar-transporting" eddies. The second is that the shape of the scalar concentration structure function may deviate appreciably from its value when $Sc = 1$, especially in its decay rate as r/l_D approaches unity due to the action of molecular diffusion [*Batchelor*, 1959]. A common approach to accommodate the first issue is to assume $l_D = Sc^{-n}\eta$ [*Hondzo*, 1998; *Jähne and Haußecker*, 1998; *Lorke and Peeters*, 2006]. It has been known for some time that n depends on whether $Sc \gg 1$ or $Sc \ll 1$. In the limit of very large or very small Sc, two length scales have been proposed to link l_D to η [*Hill*, 1978]: the Batchelor length scale $l_B = Sc^{-1/2}\eta$ when $Sc \gg 1$ [*Batchelor*, 1959] and the Corrsin length scale $l_C = Sc^{-3/4}\eta$ when $Sc \ll 1$ [*Corrsin*, 1951]. The l_B and l_C are derived from controls over the diffusive time t_m representing a diffusive scale $l_D \sim \sqrt{D_m t_m}$. If $Sc \gg 1$, then ν still restricts the diffusion time t_m. Hence, t_m is given by the Kolmogorov time scale $t_m = (\nu/\epsilon)^{1/2}$. This estimate of t_m results in $l_D \sim \sqrt{(\nu/Sc)(\nu/\epsilon)^{1/2}}$, which recovers $l_D = l_B = Sc^{-1/2}\eta$. This argument assumes that the velocity straining rate on turbulent scalar concentration fluctuations is approximately uniform for scales commensurate to or smaller than η. When $Sc \ll 1$, then molecular diffusion is far more restrictive than fluid viscosity and $t_m = (D_m/\epsilon)^{1/2}$ instead of $(\nu/\epsilon)^{1/2}$ so that $l_D \sim \sqrt{D_m(D_m/\epsilon)^{1/2}}$ recovering $l_D = l_C = Sc^{-3/4}\eta$. Hence, for $Sc \geq 1$, $l_D \geq \eta$ and ν still restrict the diffusion time t_m. When setting $l_D = Sc^{-n}\eta$ and assuming $Sc \geq 1$, the structure function approach with viscous corrections yield

$$\bar{E} \sim D_m{}^n \nu^{1/2-n} \left[\frac{1}{\nu\kappa z_o}\right]^{1/4} u_*^{3/4}\overline{\Delta C}. \tag{10}$$

Equation (10) is identical to Brutsaert's result when $n = 1/2$ (or $Sc \geq 1$). When $Sc \ll 1$, equation (10) no longer holds because molecular diffusion dampens scalar concentration fluctuations well before the attainment of the scale at which $l_D \sim \eta$. Stated differently, during an eddy turnover time commensurate to its Kolmogorov time scale value, scalar mass diffusion can be sufficiently large due to a large D_m or small Sc so that the scalar does not move coherently along l_D. The obvious question remains how close to unity does Sc need to be when $Sc < 1$ so that equation (7) still hold (derived when $Sc = 1$). To answer this question, it is necessary to explore how Sc modifies the molecular-diffusion corrections to K41 scaling in the scalar concentration structure function in the vicinity of the larger of all scales here (i.e., $l_D = l_C = Sc^{-3/4}\eta$). Recall that $l_C \geq l_B \geq \eta$ when $Sc \leq 1$. Using approximations and assumptions analogous to those used to arrive at the VKH equation for the vertical velocity structure function, the scalar structure function (written for air temperature) is given by [*Monin and Yaglom*, 1975]

$$D_{TTU}(r)-2D_m\frac{dD_{TT}(r)}{dr}=-\frac{4}{3}\epsilon_T r, \tag{11}$$

where $D_{TT}(r)=\overline{[T(x+r)-T(x)]^2}$ is the temperature structure function (typical of a scalar with Schmidt number close to 0.7), $D_{TTU}(r)$ is the mixed velocity-temperature structure function (note that $D_{UUU}(r)$ is the triple longitudinal velocity structure function that describes the vortex stretching effects in the VKH equation), and ϵ_T is the temperature variance dissipation rate. As a closure model, assume the skewness F is given by

$$F=-\frac{D_{TTU}(r)}{D_{TT}(r)\sqrt{D_{uu}(r)}} \tag{12}$$

and is constant independent of r or Sc. Its value is $F=0.4$ determined from measurements in the atmospheric surface layer as reviewed elsewhere [*Katul et al.*, 1997]. This condition already presumes that Sc effects (or l_C) are not contaminating the range of scales where $D_{TTU}(r)$ is large compared to the molecular-diffusion term on the left-hand side of equation (11). The $D_{TTU}(r)$ dominates in the inertial subrange but rapidly decays at very small scales. For analytical tractability, assume K41 scaling also holds for $D_{UU}(r)=C_{o,u}\epsilon^{2/3}r^{2/3}$ throughout (i.e., no bottleneck or viscous corrections to the longitudinal velocity structure function), where $C_{o,u}$ is the Kolmogorov constant for the longitudinal velocity component. It is to be noted that $D_{uu}(r)$ or $D_{ww}(r)$ are independent of Sc and do not introduce any Schmidt number corrections to equation (11). Replacing this approximation into equation (11) and simplifying results in

$$D_m\frac{dD_{TT}(r)}{dr}+aD_{TT}(r)r^{1/3}=br, \tag{13}$$

where

$$a=\frac{F\sqrt{C_{o,u}}\epsilon^{1/3}}{2D_m};\ b=\frac{2\epsilon_T}{3D_m}. \tag{14}$$

Solving equation (13) subject to $D_{TT}(0)=0$ yields

$$\frac{D_{TT}(r)}{ba^{-1}r^{2/3}}=1-\frac{1}{\zeta_T}\mathrm{Daw_F}(\zeta_T), \tag{15}$$

where $\zeta_T=r^{2/3}(\sqrt{3a}/2)$ and $ba^{-1}r^{2/3}$ is the K41 scaling for the air temperature structure function given by $C_T\epsilon_T\epsilon^{-1/3}r^{2/3}$, and C_T is the Kolmogorov-Obukhov-Corrsin constant related to F and the Kolmogorov constant via $C_T=(4/3)C_{o,u}^{-1/2}F^{-1}$. For large r, ζ_T is large and

$$1-\frac{1}{\zeta_T}\mathrm{Daw_F}(\zeta)\approx 1. \tag{16}$$

$D_{TTU}(r)$ is much larger than the molecular-diffusion term and K41 scaling describes $D_{TT}(r)$ with no Sc dependency. For sufficiently small r, a two-term series expansion of $\mathrm{Daw_F}$ becomes plausible and yields

$$1-\frac{1}{\zeta_T}\mathrm{Daw_F}(\zeta_T)\approx\frac{2}{3}\zeta_T^2. \tag{17}$$

The dependence of this solution on Sc in the vicinity of $r=l_C=Sc^{-3/4}\eta$ (i.e., the largest scale where D_m effects are expected) is compared to the case when $Sc=1$. To proceed further, it is noted that $\epsilon^{-1/3}=\nu\eta^{-4/3}$, $\zeta_T=\beta_T Sc^{1/2}(r/\eta)^{2/3}$, and $\beta_T=(3/2)^{1/2}2^{-1}(C_oF)^{1/4}$, which when inserted into equation (17) lead to a molecular-diffusion correction to the K41 scaling of $D_{TT}(l_C)$ given by

$$\frac{2}{3}\zeta_T^2\approx\frac{2}{3}Sc(Sc^{-3/4})^{4/3}. \tag{18}$$

The result here suggests that molecular corrections to K41 scaling in the temperature structure function do not vary with Sc provided F is constant independent of Sc (and r). The shape of $D_{TT}(l_C)$ at $Sc=0.7$ and $Sc=1$ given by equation (18) is, for all purposes, self-similar. Moreover, the shape of $D_{TT}(r)$ is similar to $D_{ww}(r)$ when assuming $Sc=1$. A logical follow-up question to ask is whether there is experimental support for the similarity between $D_{TT}(r)$ and $D_{ww}(r)$ shapes when $Sc=0.7$. Measured air temperature spectra do exhibit K41 scaling laws followed by "bottlenecks" at the crossover from inertial to molecular dissipation regimes at scales $r/\eta\approx 0.1$ [*Hill*, 1978] commensurate with those reported for vertical velocity and turbulent kinetic

energy [*Pope*, 2000; *Saddoughi and Veeravalli*, 1994] consistent with the analysis here. Hence, equation (7) may still hold given the self-similarity in shapes of the temperature and vertical velocity structure functions (even when $Sc = 0.7$). Besides Sc, n may be sensitive to other factors such as surface state, wind conditions, and associated roughness [*Csanady*, 1990; *Jähne and Haußecker*, 1998; *Esters et al.*, 2016] as discussed in many air-sea exchange studies. In particular, *Jähne and Haußecker* [1998] demonstrated variability in n with $n = 3/4$ measured for low wind speed conditions whereas n dropped to 1/2 with increasing wind speed over water surfaces (with small waves) for the same soluble gas.

5. Conclusion

The work in *Brutsaert* [1965] demonstrated that the mean evaporation rate scales with the 3/4 power law of the friction velocity and the square-root of molecular diffusivity. The two aforementioned results were recovered here using a simpler surface renewal analysis than the one offered by *Brutsaert* [1965]. Key to all surface renewal schemes is the assumption about the probabilistic nature of the contact durations of eddies with the surface. The two results in *Brutsaert* [1965] were shown to be robust to this assumption provided the mean contact duration remains proportional to the Kolmogorov time microscale. A new approach was proposed that links eddy turnover velocity to the structure function of the vertical velocity that is adequate when the molecular Schmidt number is near or larger than unity. The main advantage of the new approach is that the structure function exhibits universal scaling laws (e.g., K41) that are modified by viscous cutoff for small scales. However, these viscous-cutoff modifications can be derived from the von Kármán-Howarth equation [*Kármán and Howarth*, 1938]. Hence, the new approach ameliorates the ad hoc nature of the surface renewal scheme used in *Brutsaert* [1965]. It directly results in a mean evaporation rate equation that is identical to the one derived in *Brutsaert* [1965] when imposing an effective eddy size commensurate with the Kolmogorov microscale. It can be concluded that the findings in *Brutsaert* [1965] are more general than originally proposed provided the finest eddies at the surface-atmosphere interface scale with the Kolmogorov length (for the structure function approach) or time (for the surface renewal approach) scales.

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