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Nonequilibrium thermodynamics of circulation regimes in optically-thin, dry atmospheres

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Abstract

An extensive analysis of an optically-thin, dry atmosphere at different values of the thermal Rossby number $\mathcal{Ro}$ and of the Taylor number $\mathcal{F}_f$ is performed with a general circulation model by varying the rotation rate Ω and the surface drag τ in a wide parametric range. By using nonequilibrium thermodynamics diagnostics such as material entropy production, efficiency, meridional heat transport and kinetic energy dissipation we characterize in a new way the different circulation regimes. Baroclinic circulations feature high mechanical dissipation, meridional heat transport, material entropy production and are fairly efficient in converting heat into mechanical work. The thermal dissipation associated with the sensible heat flux is found to depend mainly on the surface properties, almost independent from the rotation rate and very low for quasi-barotropic circulations and regimes approaching equatorial super-rotation. Slowly rotating, axisymmetric circulations have the highest meridional heat transport. At high rotation rates and intermediate-high drag, atmospheric circulations are zonostrophic with very low mechanical dissipation, meridional heat transport and efficiency. When τ is interpreted as a tunable parameter associated with the turbulent boundary layer transfer of momentum and sensible heat, our results confirm the possibility of using the Maximum Entropy Production Principle as a tuning guideline in the range of values of Ω . This study suggests the effectiveness of using fundamental nonequilibrium thermodynamics for investigating the properties of planetary atmospheres and extends our knowledge of the thermodynamics of the atmospheric circulation regimes.

Keywords: Circulation regimes, nonequilibrium thermodynamics,
terrestrial planetary atmospheres, baroclinic instability, entropy production
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1. Introduction

In the last two decades, more than 700 planets outside the solar system (exoplanets) have been discovered (Udry and Santos, 2007), and the Kepler Space Telescope has recently located over 2,000 exoplanet candidates (Borucki et al., 2011). The study of exoplanets and their climates is in its early stage and it is quickly developing (Seager and Deming, 2010). Observational data are still poor and difficult to obtain, particularly for those planets – super-Earths (Charbonneau et al., 2009) – that might be capable of sustaining liquid water and thus potentially suitable for life. Nevertheless, the discovery of exoplanets is extending the scope of planetary sciences towards the study of the so-called “exoclimates” (Heng, 2012; Burrows et al., 1997; Heng et al., 2011a; Showman et al., 2009; Joshi, 2003; Merlis and Schneider, 2010; Lewis et al., 2010; Pierrehumbert, 2010; Thrastarson and Cho, 2011; Rauscher and Menou, 2012; Dobbs-Dixon et al., 2012). Exoplanets and their atmospheres are in general capable of supporting a broad set of circulation regimes since they are characterized by a range of physical (atmospheric composition, rotation rate, dimension, surface) and orbital (obliquity, eccentricity, distance from the parental star, spectral type of the parental star, presence or not of phase locking) parameters even wider than that of Solar System planets (Williams and Pollard, 2002). Planetary science aims at predicting and classifying in a concise but comprehensive way exoclimates once the main orbital and physical parameters are known.

Recently Read (2011) noted that the large variety of circulation regimes may be better understood by adopting the fluid-dynamical method of similarity, i.e. by defining a set of dimensionless numbers that fully characterise the planetary circulations. Two climate states that share the same set of dimensionless numbers are dynamically equivalent and so the statistical properties of one can be mapped onto those of the other. Obviously the set of parameters is fairly large, and one of the main objectives of planetary science is to understand what is the minimal number of dimensionless parameters needed to define virtually equivalent circulations (Wang, 2012; Showman et al., 2010). In this study we focus on the impact of two parameters, the rotation rate Ω

and on the surface turbulent exchange rate τ , on the atmospheric circulation of an Earth-like dry atmosphere. The choice of such parameters naturally leads to the definition of two dimensionless numbers, the thermal Rossby number $\mathcal{Ro}$ and the Taylor (frictional) number $\mathcal{F}_f$ (Read, 2011).

Over the last three decades, the effect of the planetary rotation on atmospheric circulation has been investigated in some details with the aid of general circulation models (Hunt, 1979; Williams, 1988a,b; Navarra and Boccaletti, 2002; Genio and Suozzo, 1987; Geisler et al., 1983; Read, 2011; Vallis and Farneti, 2009). Variations in the value of Ω impacts directly the size of the baroclinic waves and the extent of the Hadley cell, which are the main features of the large-scale Earth atmospheric circulation. The size of the baroclinic disturbances, being proportional to the Rossby deformation radius (Eady, 1949), scales as $1/\Omega$. The latitudinal extent of the Hadley cell also scales as $1/\Omega$ (Held and Hou, 1980). Numerical simulations of slowly rotating Earth-like planets and of Solar System planets like Venus and Titan (Clancy et al., 2007; Hourdin et al., 1995) have shown the presence of one poleward-extended Hadley cell in each hemisphere and the weakening or complete disappearing of the midlatitude baroclinic disturbances. On the other hand, at fast rotation rates the emergence of multiple cells in the meridional circulation and multiple jets in the zonal circulation has been observed both in numerical simulations (Williams, 1988a, 1978) and observations (e.g. Jupiter).

The dynamical effects of the solid lower boundary of terrestrial planets on the atmospheric circulation is also quite important in order to understand planetary circulations and has not been fully addressed yet (Showman et al., 2010). The characteristics of the surface have been recognised as a key factor in shaping Earth’s atmospheric circulation (James, 1994; James and Gray, 1986), although this topic has received less attention than that related to Ω . The surface of a terrestrial planet, due to its roughness, affects the turbulent flow within the planetary boundary and thus the exchange of momentum and energy between the surface and the atmosphere (Arya, 1988). It has been shown (James and Gray, 1986; James, 1987; Kleidon et al., 2003) that the reduction of the surface drag leads to strong horizontal barotropic shears in the zonal mean flow. By using a two-level quasi-geostrophic model, James (1987) showed that the growth rate of the most unstable baroclinic modes is reduced considerably by the strong horizontal wind shears. This is related to the general fact that the linearised baroclinic instability equations obey the Squire’s theorem (Kundu and Cohen, 2004). The role of drag has received

71 some attention in the exoplanets context (Rauscher and Menou, 2012) but,
72 to the authors' knowledge, has not been systematically investigated so far
73 for rotation rates which are different from the Earth's. In this study we
74 investigate the combined effect of rotation speed and surface roughness on
75 the dynamics, linking it to the nonequilibrium thermodynamics of the system.

76 Thermodynamics provide a way for characterizing concisely a complex
77 physical system, bringing together comprehensive but minimal physical in-
78 formation. The atmosphere of a planet is an example of a nonequilibrium
79 system (Gallavotti, 2006; DeGroot and Mazur, 1984; Kleidon, 2009), and its
80 general circulation redistributes energy in order to compensate for the ra-
81 diative differential heating between hot and cold regions. The atmospheric
82 circulation therefore is fuelled by the conversion of available potential energy
83 due to large temperature gradients into kinetic energy. The atmosphere, in
84 other terms, produces mechanical work, acting as a heat engine (Lorenz,
85 1967; Peixoto et al., 1991; Johnson, 2000; Lucarini, 2009). It seems therefore
86 natural to adopt nonequilibrium thermodynamics as a general framework for
87 studying exoclimates. Such an approach has been, for example, applied in
88 Lucarini et al. (2010) and Boschi et al. (2012) for studying the bistability of
89 an Earth-like planet. Furthermore, thermodynamical disequilibrium drives
90 a variety of irreversible processes, from frictional dissipation to chemical re-
91 actions. The irreversibility of climatic processes is quantified by the mate-
92 rial entropy production (Goody, 2000; Kleidon and Lorenz, 2005; Kleidon,
93 2009). The interest in studying climate material entropy production largely
94 stemmed from the proposal of the maximum entropy production principle
95 (MEPP) by Paltridge (Paltridge, 1975, 1978, 2001), who suggested that the
96 climate adjusts in such a way as to maximize the material entropy produc-
97 tion. In its weak form, the MEPP suggests to use the entropy production
98 as a target function to be maximized when tuning an empirical or uncertain
99 parameter of a model (Kleidon et al., 2003; Kunz et al., 2008). Whereas the
100 theoretical foundations of MEPP are still unclear (Dewar, 2005; Grinstein
101 and Linsker, 2007; Goody, 2007), such a conjecture has also been proposed
102 as a way to estimate the meridional heat transport of other planets, such
103 as Mars and Titan (Lorenz et al., 2001; Jupp and Cox, 2010) and poten-
104 tially to exoplanets too, and has stimulated the re-examination of climatic
105 dissipative processes (Peixoto et al., 1991; Goody, 2000; Pauluis and Held,
106 2002a,b; Kleidon and Lorenz, 2005; Fraedrich and Lunkeit, 2008; Pascale
107 et al., 2011a).

108 In this study we perform a large ensemble of numerical simulations with

an Earth-like general circulation model for many different values of Ω and τ in order to compute the dissipative properties ζ (where ζ is any dissipative function, e.g. material entropy production) of circulations of dry atmospheres at different thermal Rossby and Taylor numbers, $\zeta(\mathcal{R}o, \mathcal{F}_f)$. We relate, for the first time, the properties of $\zeta(\mathcal{R}o, \mathcal{F}_f)$ to the different circulation regimes and extend our knowledge on the global thermodynamic properties of rotating fluids. We anticipate that particular regimes (e.g. baroclinic, zonostrophic, super-rotation) are effectively characterized in terms of their thermodynamic properties. We conclude with a brief analysis of how effectively the MEPP can be used to infer the optimal value for an uncertain or empirical parameter, in this case exactly the time scale controlling the exchange of momentum and energy between free atmosphere and the surface.

The paper is organized as follows, In Section 2 we will shortly discuss the dimensionless parameters relevant for this study. In Section 3 the model and the experimental setup are presented. The characterization of different dynamical regimes is the subject of Section 4 whereas in Section 5 the thermodynamical properties of the circulation regimes are analysed. In Section 6 the main conclusions are summarized.

2. Parametric range of general circulations and dimensionless numbers

The role of the rotation rate in planetary circulations has been first investigated in laboratory experiments with a thermally driven rotating annulus (Hide, 1953, 1969; Hide and Mason, 1975; Read et al., 1998; Read, 2001; Wordsworth et al., 2008; Hide, 2010). The system consists of a fluid confined between coaxial cylinders maintained at two different temperatures and rotating at an angular velocity Ω . When the basic parameters Ω and ΔT (temperature difference between the inner and outer cylinder) are varied, a wide variety of flow patterns is observed. Different dynamical regimes can be identified if results are grouped with respect to two dimensionless parameters, the *thermal Rossby number*:

$$\mathcal{R}o = \frac{g\alpha D\Delta T}{\Omega^2 L^2}, \quad (1)$$

and the *Taylor number*:

$$\mathcal{T}a = \frac{4\Omega^2 L^5}{\nu^2 D}, \quad (2)$$

140 in which L is the channel width, D its depth, ν the kinematic viscosity
 141 of the fluid, α its volumetric expansion coefficient, and g the gravitational
 142 acceleration.

143 Read (2011) has extended the definition of the thermal Rossby number
 144 and of the Taylor number to the case of atmospheric circulations. The anal-
 145 ogous of the thermal Rossby number is defined as:

$$\mathcal{R}_o = \frac{R\Delta\theta_h}{\Omega^2 a^2}, \quad (3)$$

146 where a is the planet’s radius, R the specific gas constant and $\Delta\theta_h$ the hor-
 147 izontal (potential) temperature contrast between equator and poles. A dif-
 148 ference between the definitions in eq. (1) and eq. (3) is that $\Delta\theta_h$ is not
 149 fixed externally but rather determined by the circulation itself. In the fol-
 150 lowing we will take $\Delta\theta_h = \Delta\theta_{hE}$, as done for example in Mitchell and Vallis
 151 (2010), where θ_{hE} is the radiative-convective equilibrium potential temper-
 152 ature, since this is externally determined by the incoming stellar radiative
 153 energy and thus a more objective quantity to describe the horizontal differ-
 154 ential driver for the circulation. A Taylor number can be defined analogously
 155 to the case of the rotating annulus as:

$$\mathcal{F}_f = 4\Omega^2 \tau_f^2 \quad (4)$$

156 in which τ_f is the typical timescale for kinetic energy dissipation. We note
 157 that $\mathcal{F}_f \propto (\tau_f/\tau_{rot})^2$, where $\tau_{rot} = 2\pi/\Omega$, i.e. $\mathcal{F}_f$ is proportional to the
 158 ratio of (the squares of) the typical timescales associated with turbulent
 159 dissipation of kinetic energy and rotation. For planets with a solid core, τ_f is
 160 the surface drag timescale and is in general determined by the characteristics
 161 of the surface. The use of (3) and (4) has been proved to be very useful in
 162 classifying atmospheric circulation (Wang, 2012).

163 3. Model and experimental setup

164 3.1. The Planet Simulator

165 Numerical simulations have been performed with the Planet Simulator
 166 (PlaSim), a general circulation model of intermediate complexity (Fraedrich
 167 et al., 2005). The model is freely available at www.mi.uni-hamburg.de/plasim.
 168 PlaSim is a fast running model and it is therefore suitable for large-ensemble
 169 numerical experiments. Moreover, a full set of thermodynamic diagnostics

170 is available, thus making it well suited for this work (Fraedrich and Lunkeit,
171 2008; Lucarini et al., 2010).

172 The atmospheric dynamic core uses the primitive equations, which are
173 solved using a spectral transform method (Eliassen et al., 1970; Orszag, 1970).
174 Interaction between radiation and atmosphere is dealt with using simple but
175 realistic longwave (Sasamori, 1968) and shortwave (Lacis and Hansen, 1974)
176 radiative schemes. In particular the incoming solar flux F_{SW}^{toa} at the top of
177 the atmosphere (TOA) is

$$F_{SW}^{toa} = S_0 \cos Z \quad (5)$$

178 where S_0 is the solar constant (1365 W m^{-2}) and Z the zenith angle, which
179 is in general a function on the latitude, time of the year and time of the
180 day, and it is computed following Berger (1978). All simulations have been
181 performed with orbital parameter – obliquity, eccentricity, distance from the
182 Sun, typical of Earth. Other sub-grid scale parametrisations include interac-
183 tive clouds (Stephens, 1978; Stephens et al., 1982; Slingo and Slingo, 1991),
184 moist (Kuo, 1965, 1974) and dry convection, large scale precipitation, bound-
185 ary layer fluxes and vertical and horizontal diffusion (Louis, 1979; Louis et al.,
186 1981; Laursen and Eliassen, 1989). More information can be found in PlaSim
187 reference manual, freely available at www.mi.uni-hamburg.de/Downloads-un.245.0.html.
188

189 In all simulations the lower boundary is a flat surface with prescribed
190 albedo and heat capacity (see Table 1). This is implemented with a shallow
191 energy-conserving slab-ocean model with an areal heat capacity ($C_{slab} = 10^7$
192 $\text{J K}^{-1} \text{m}^{-2}$) comparable to that chosen in Frierson et al. (2006) and Heng et al.
193 (2011b). In this way we avoid fixed surface temperature and have a simple
194 but energetically consistent climate model. The surface temperature evolves
195 in time according to $C_{slab} \dot{T}_s = F_{SW}^{surf} + F_{LW}^- = \sigma T_s^4 - F_T$ (F_{SW}^{surf} net solar
196 radiation at the surface, F_{LW}^- downward longwave radiation at the surface,
197 F_T surface sensible heat flux). We set the depth of the mixed layer to 5 m
198 in order to have an areal heat capacity ($C_{slab} = 10^7 \text{ J K}^{-1} \text{m}^{-2}$) comparable
199 to that chosen in Frierson et al. (2006) and Heng et al. (2011b). We have
200 checked our result at $C_{slab} = 10^8 \text{ J K}^{-1} \text{m}^{-2}$ too, finding little effects on the
201 circulations and on the global thermodynamical properties. Simulations are
202 performed at T42 spectral resolution ($2.8^\circ \times 2.8^\circ$) with ten levels (T42/10LEV
203 in the following).

204 In this study we consider dry atmospheres. Dry atmospheres are relevant
205 for planetary (e.g. Mars) and paleoclimatological (e.g. Snowball Earth) stud-

ies and, moreover, allow us to avoid the role of phase transitions associated with condensing substances, simplifying the problem and making neater the connection between dynamics and thermodynamics of the system. Such configuration is obtained by switching off the surface evaporation module and starting from a dry atmospheric condition. Water vapour is consequently not inserted within the atmosphere, which remains dry for all timesteps.

3.2. The strength of the turbulent surface exchanges

In order to have a wide and controlled variation in $\mathcal{F}_f$ (Eq. 4), we simplify the representation of the surface fluxes. In PlaSim the temperature tendency of the first atmospheric layer (of thickness dz) due to the turbulent sensible heat flux, $(\partial T/\partial t)_{shf}$, is computed as:

$$\left(\frac{\partial T}{\partial t}\right)_{shf} = -\frac{F_T}{\rho c_p dz} = \frac{\gamma_h |\mathbf{u}|}{dz} (T_s - \xi T) = \frac{T_s - \xi T}{\tau_h(\mathbf{x}, t)}, \quad (6)$$

in which $F_T = \gamma_h |\mathbf{u}| (T_s - \xi T)$ is the surface sensible heat flux, $\gamma_h = (k/\ln(z/z_0))^2 f(Ri, z_0)$ is the heat transfer coefficients (z is height from the surface, k is the von-Karman parameter, z_0 is the surface roughness, and f is an empirical function dependent on stability (as expressed by the Richardson number Ri) and surface roughness), ξ is the Exner factor (for more details see Louis, 1979; Lunkeit et al., 2010). The parameter τ_h has time dimension and in a standard run is a function of space and time, $\tau_h(x, y, z, t) = dz/(\gamma_h(x, y, t)|\mathbf{u}(x, y, t)|)$ but remains of the same order of magnitude. Since we are interested in variations of orders of magnitude in τ_h , we substitute the locally computed τ_h with a fixed (in space and time) time scale τ_h as:

$$\left(\frac{\partial T}{\partial t}\right)_{shf} = -\frac{\xi T - T_s}{\tau_h}. \quad (7)$$

Similarly to eq. (6), for the wind tendency due to the surface stress, $(\partial \mathbf{u}/\partial t)_{stress}$, we have:

$$\left(\frac{\partial \mathbf{u}}{\partial t}\right)_{stress} = -\frac{\mathbf{u}}{\tau_m(\mathbf{x}, t)}. \quad (8)$$

with $\tau_m(x, y, z, t) = dz/(\gamma_m(x, y, t)|\mathbf{u}(x, y, t)|)$ and the drag coefficient γ_D defined similarly to γ_h . Again we substitute the locally compute $\tau_m(\mathbf{x}, t)$ with a fixed (in space and time) drag timescale τ_m (Rayleigh friction timescale). Generally the drag and heat transfer coefficients γ_D and γ_h – and therefore

the time constants τ_m and τ_h – have similar magnitude. This is particularly true in the case of neutral flows, for which $\gamma_D = \gamma_h$ is indeed a very good approximation (Arya, 1988; Louis, 1979). For non-neutral flows, γ_h and γ_D are different but still of the same order of magnitude, as can be seen in Fig. 11.6 of Arya (1988). On the base of this and since in this study we are going to explore a wide parametric range, we assume for the sake of simplicity:

$$\tau_m = \tau_h = \tau. \quad (9)$$

Experiments are performed for $\Omega^* = \Omega/\Omega_E = 1/10, 1/5, 1/2, 1, 2, 4, 8$, where Ω_E is the Earth rotation rate. For each value of Ω^* we run the model with $\tau = 2700, 3600, 10800, 21600, 43200, 86400, (86400 \times 3), (86400 \times 10), (86400 \times 30), (86400 \times 100), (86400 \times 500)$ seconds, that is from 45 minutes (model timestep for $\Omega/\Omega_E \leq 1$) to 500 days. Simulations with very large τ are representative of an atmosphere with no solid lower boundary (James, 1994; Menou and Rauscher, 2009; Heng et al., 2011b).

Let us note that as Ω increases, the typical size of the baroclinic disturbances L_c decreases as (Eady, 1949)

$$L_c = 2.4\pi L_R, \quad (10)$$

with the Rossby deformation radius $L_R = NH/f$ (James, 1994; Williams, 1988a), N the buoyancy frequency, H the height scale and $f = 2\Omega \sin \varphi$ the Coriolis parameter. For our dry-atmosphere simulations an order-of-magnitude estimate at the midlatitudes for $\Omega^* = 8$ leads to $\Delta\theta \approx 110$ K, $\bar{\theta} \approx 240$ K (see, e.g., Fig.3(h)), $\Delta z = 9$ km, $N \approx (g/\bar{\theta}(\Delta\theta/\Delta z))^{1/2} \approx 2 \times 10^{-2}$ s⁻¹ and therefore to $L_R \sim 200$ Km. This implies that T42 simulations (spatial resolution about 250 Km) should be able to capture at least the largest eddies at $\Omega^* = 8$ and more than adequate for $\Omega^* \leq 4$.

4. Circulation regimes at different $\mathcal{Ro}$ and $\mathcal{F}_f$

The diagram in Fig. 1(b) shows the dimensionless space $(\mathcal{F}_f, \mathcal{Ro})$. The over-plotted bullet points represent numerical experiments performed at $\Omega^* = 0.1$ (circles, denoted as “slow rotation”), $\Omega^* = 1$ (squares, “intermediate rotation”) and $\Omega^* = 8$ (triangles, “fast rotation”) for strong, intermediate and weak drag condition (τ equal to 45 minutes, 1 day and 500 days respectively) whose mean meridional and zonal circulations are shown in Fig. 2 and Fig. 3

and delimit the portion of the $(\mathcal{F}_f, \mathcal{R}o)$ space covered by the numerical simulation performed in this study. We have over-plotted the corresponding values of Ω^* (horizontal dot-dashed lines) and τ (dotted lines) in order to highlight the connection between the dimensionless numbers and the physical parameters Ω^* and τ . Note that Ω^* and $\mathcal{R}o$ as well as τ and $\mathcal{F}_f$ point in opposite directions. In order to help to set the stage for the reader to understand the results in the following and make it easier to interpret the montage of figures (3) and (2), we anticipate the main characteristics of the simulated circulations:

1. At high thermal Rossby number ($\mathcal{R}o \geq 8$), the decrease of the surface drag controls the transition from counter- to super-rotating (SR in Fig.1(a)) equatorial flow. Super-rotation is approached for $\mathcal{F}_f \geq 10^4$;
2. At intermediate rotation speed ($1 \leq \mathcal{R}o \leq 0.01$), strong drag ($\mathcal{F}_f \leq 10$) is associated with axisymmetric circulations (AR in Fig. 1(a)). The decrease of τ leads to the appearance of the indirect Ferrel cell for $10 \leq \mathcal{F}_f \leq 10^5$ characterized by baroclinic activity (BC in Fig.1(a)); further decrease of the surface drag ($\mathcal{F}_f \geq 10^5$) leads to the emergence of a barotropic flow (BT in Fig.1(a)) characterised by a large reduction in the vertical shears of the zonal wind and the complete disappearing of the Ferrel cell;
3. For fast rotations ($\mathcal{R}o \leq 10^{-3}$) the increase the of Taylor frictional number ($\mathcal{F}_f > 10^4$) leads to the appearance of a multi-jet, zonostrophic flow (ZN in Fig. 1(a)) for $\mathcal{F}_f > 10^3$ ($\tau > 6$ hours)

Boundaries between the different regimes are schematically sketched in Fig. 1(a). In the following we give a detailed description of the different regimes.

4.1. Slow rotation ($\mathcal{R}o = 8$)

Fig. 2(a), 2(b), 2(c) and Fig. 3(a), 3(b), 3(c) show the slow rotation rate ($\mathcal{R}o = 8$). Such circulations are dominated by one Hadley cell in each hemisphere which extends northward up to the poles (this regime is denoted AS in the Fig.1(a)). This is a general consequence of the conservation of angular momentum and in agreement with the theory of the Hadley circulation of Held and Hou (1980). The temperature features almost no latitudinal dependence, especially in the middle atmosphere. This is typical of slowly rotating planets (Williams, 1988a; Navarra and Boccaletti, 2002), and is due

297 to the strong Hadley cell circulation. It is interesting to note the effect
 298 of the surface drag on shaping the Hadley circulation. By comparing Fig.
 299 2(e) to Fig. 2(b) ($\mathcal{R}o, 10^{-1} \rightarrow 8$; $\mathcal{F}_f, 10^3 \rightarrow 10$) and Fig.2(c) to Fig.2(f)
 300 ($\mathcal{R}o, 10^{-1} \rightarrow 8$; $\mathcal{F}_f, 10^7 \rightarrow 10^5$) we note a decrease of the counter-rotating
 301 westward upper-level equatorial jet approaching the beginning of the equa-
 302 torial super-rotation (for example compare Fig 2(c) to Fig. 13 of Heng and
 303 Vogt, 2011). Equatorial super-rotation is indeed expected to take place when
 304 $\mathcal{R}o \gg 1$ (Mitchell and Vallis, 2010). Therefore simulations with $\Omega^* < 1/10$
 305 and moderate or high drag are needed in order to obtain fully super-rotating
 306 atmospheric circulations (as is the case, for example, for Venus to Titan).

307 4.2. Intermediate rotation ($\mathcal{R}o = 0.08$)

308 In the medium rotation case ($\mathcal{R}o = 0.08$), we have atmospheric circu-
 309 lations characterized by strong eastward zonal jets at about $50 - 60^\circ$ and
 310 by a thermally direct (Hadley) and indirect (Ferrel) meridional cell (Fig. 2
 311 (d,e,f) and Fig. 3 (d,e,f)). The general circulation is considerably affected
 312 by the different surface properties. In particular we note that at large $\mathcal{F}_f$,
 313 the flow develops strong barotropic horizontal shears, as first discussed by
 314 James and Gray (1986). Note that, as we are considering a dry optically-thin
 315 atmosphere, none of the three circulations shown in Fig. (2(d)-2(f)) is close
 316 to the one we observe on Earth (e.g. Peixoto and Oort (1992)) but rather
 317 similar to that of Mars (Lewis et al., 2010).

318 The effect of the surface drag is particularly evident in the meridional
 319 circulation, which is largely modified by the surface properties. A clear
 320 thermally direct-indirect cell structure emerges in the intermediate cases
 321 $\mathcal{F}_f \sim 10^2$ ($\tau \sim 1$ day), with the boundaries of the Hadley cell at about
 322 40° . The intensity and the extent of the indirect cell is greatly reduced
 323 in the high drag ($\mathcal{F}_f \leq 10^{-1}$) case, when the baroclinic waves are largely
 324 suppressed and the flow tend to become axisymmetric. The Ferrel cell is
 325 instead completely suppressed in the low drag ($\mathcal{F}_f \geq 10^5$) case, where the
 326 flow becomes barotropic. The large impact of the surface properties on the
 327 meridional circulation is related to their impact on the baroclinic distur-
 328 bances (James and Gray, 1986), which normally develop at the edge of the
 329 thermally direct (Hadley) and indirect (Ferrel) cells. The Ferrel cell is re-
 330 lated to the presence of eddy momentum convergence, a key ingredient of
 331 baroclinic disturbances (Holton, 2004), and its disappearance points out the
 332 suppression or weakening of the midlatitude disturbances. In the presence of
 333 weak surface drag, zonal winds tend to have high values at the surface which

334 remain fairly constant with height but change sign at the midlatitudes from
 335 westward to eastward going from the equator to the poles (e.g. Fig. 2(f))
 336 thus generating a strong horizontal shear. Such strong horizontal shears in-
 337 hibit the growth of baroclinic waves, as demonstrated in (James, 1987). On
 338 the other hand, with a surface characterized by a high drag, baroclinicity
 339 is suppressed too, because the system frictional dissipation is too high and
 340 kinetic energy is rapidly extracted not giving eddies the chance to grow and
 341 develop (Kleidon et al., 2003).

342 Let us also note in Fig. 3 the presence of shallow cells embedded close
 343 to the surface embedded in a larger one. This is a characteristic of optically-
 344 thin atmospheres of rocky planets in which the solid lower boundary with low
 345 thermal inertia respond very quickly to diurnal and seasonal solar heating
 346 (Caballero et al., 2008). Similar features are indeed observed in Mars circu-
 347 lation (see e.g. figure 2 of Lewis et al., 2010). Such shallow cells disappears
 348 in fact in the additional runs we have performs at $C_{slab} = 10^8 \text{ J K}^{-1} \text{ m}^{-2}$ (not
 349 shown) and have very little effect on the thermodynamic properties we are
 350 going to discuss in the following sections.

351 4.3. Fast rotation ($Ro = 10^{-3}$)

352 Finally, in the fast rotation runs ($Ro = 10^{-3}$) we observe multiple jets
 353 (Fig. 2(g)-(i)) and multiple meridional cells (Fig. 3(g)-(i)) in agreement with
 354 previous studies (Hunt, 1979; Williams, 1988a) and with the scaling of the
 355 Rossby deformation radius (eq. 10). The decrease of L_R with the rotation
 356 rate makes baroclinic waves less and less efficient in the poleward heat trans-
 357 porting process and reduction of the meridional temperature contrast. The
 358 temperature field in fact shows larger contrast in the meridional and vertical
 359 profile, and the thermal structures tend to be in radiative-convective equi-
 360 librium. The effect of τ is mainly observed in the zonal wind profiles (Fig.
 361 2 (g,h,i)) and in the meridional stream function (Fig. 3(g,h,i)). Multi-jet,
 362 zonostrophic flow (Wang, 2012) emerges as the surface drag decreases for
 363 $\mathcal{F}_f > 10^3$, as can be seen in Fig. 3(i).

364 5. Thermodynamic analysis

365 5.1. Thermodynamic diagnostics

366 The general circulation is the result of the conversion of the available
 367 potential energy generated by radiative differential heating into mechanical
 368 work (winds), as first shown by Lorenz (1955, 1960, 1967). For an atmosphere

in a statistical steady state, the rate of generation of available potential energy, G , the rate of conversion into kinetic energy, W , and the rate of dissipation of kinetic energy through the turbulent cascade (and ultimately via viscous dissipation), D , have to be equal when averaged over long time periods (e.g. a year or longer), $\overline{G} = \overline{W} = \overline{D}$ ($\overline{(\cdot)}$ denotes the time mean). They are therefore equivalent ways of measuring the strength of the Lorenz energy cycle (Lorenz, 1955).

The energy cycle introduced by Lorenz has been set onto a thermodynamic framework through the consideration of the effective Carnot engine describing the ability of the atmosphere to perform work (Johnson, 2000; Adams and Rennó, 2005; Lucarini, 2009). The atmosphere is seen as a heat engine which generates mechanical work at average rate $\overline{W}$ from the differential heating due to radiative and material (e.g. latent heat release) diabatic processes. If $\dot{Q}^+$ and $\dot{Q}^-$ are the local positive and negative diabatic heating rate (i.e. $\dot{Q}^+ = \dot{Q}$ where $\dot{Q} > 0$ and $\dot{Q}^+ = 0$ where $\dot{Q} < 0$ and similarly for $\dot{Q}^-$) with

$$\Phi^\pm = \int \dot{Q}^\pm \rho dV, \quad (11)$$

we have that $\overline{\Phi^+} + \overline{\Phi^-} = \overline{W} \geq 0$. Moreover, one can define an efficiency η :

$$\eta = \frac{\overline{\Phi^+} + \overline{\Phi^-}}{\overline{\Phi^+}} \quad (12)$$

which gives us an indication about the capability of the general circulation of generating kinetic energy given the net heating input Φ^+ . From Eq. (7) it follows that

$$\overline{W} = \eta \overline{\Phi^+} \quad (13)$$

in full analogy with the definition of efficiency of a heat engine (Fermi, 1956). Such a quantity has been proved to be particularly relevant in marking the climatic shifts between the present day climates and the Snowball Earth (Lucarini et al., 2010; Boschi et al., 2012)

Dissipation, and therefore irreversibility, is ubiquitous in planetary atmospheres and, more generally, in nonequilibrium systems. The kinetic energy of the atmospheric flow is ultimately transferred through a turbulent cascade to smaller scales where it is then dissipated into heat by friction due to viscosity. Thermal dissipation due to sensible heat fluxes between the surface and lower atmosphere is another irreversible process which may take place

399 in planetary atmospheres. Planets whose atmospheres allow phase transi-
 400 tions of one or more of their chemical substances (e.g. water on Earth or
 401 methane on Titan) also experience further irreversible processes as evapo-
 402 ration/condensation and diffusion (Goody, 2000; Pauluis and Held, 2002b).
 403 Irreversible processes are associated with a positive-defined material entropy
 404 production (Peixoto et al., 1991; DeGroot and Mazur, 1984; Kondepudi and
 405 Prigogine, 1998; Fraedrich and Lunkeit, 2008; Kleidon, 2009). General dis-
 406 cussions about the entropy budget of the climate system and about how to
 407 estimate it from climate models can be found in Peixoto et al. (1991), Goody
 408 (2000), Kleidon and Lorenz (2005), Kleidon (2009), Pascale et al. (2011a),
 409 Pascale et al. (2011b), Lucarini et al. (2011). For a climate with a dry at-
 410 mosphere the material entropy production is due to two kinds of processes:
 411 dissipation of kinetic energy and sensible heat fluxes. If ϵ^2 is the local rate
 412 of kinetic energy dissipation such that $D = \int \epsilon^2 \rho dV$, the entropy production
 413 associated with it reads:

$$\dot{S}_{kedis} = \int \frac{\epsilon^2}{T} \rho dV. \quad (14)$$

414 In PlaSim the dissipation of kinetic energy is due to: (i) turbulent stresses in
 415 the surface boundary layer (which accounts for more than 50% of the overall
 416 dissipation) and, gravity wave drag, implemented as a Rayleigh friction at
 417 the highest level with a timescale of 50 days, which we define as D_{phys} .
 418 Such contribution to the total mechanical dissipation is diagnosed in the
 419 model as $1/2 \int \rho dz (\mathbf{v}_a^2 - \mathbf{v}_b^2)$ where $\mathbf{v}_b$ and $\mathbf{v}_a$ is the velocity before and
 420 after the application of the boundary layer scheme and Rayleigh friction;
 421 (ii) numerical dissipation due to numerical diffusion of momentum (Johnson,
 422 1997), which we call D_{num} . More precisely, in PlaSim horizontal diffusion
 423 is implemented by a 8th order hyperdiffusion term applied to the vertical
 424 component of the relative vorticity $\zeta = \mathbf{k} \cdot (\nabla \times \mathbf{v})$ and horizontal wind
 425 divergence $\delta = \nabla_h \cdot \mathbf{v}$, $\kappa \nabla^8(\zeta, \delta)$, where κ is a coefficient of numerical diffusion
 426 – the prognostic equations for the horizontal velocity are transformed into
 427 equations for ζ and δ , for more details on PlaSim dynamical core see Lunkeit
 428 et al. (2010) –. Although it is hard to interpret D_{num} as representative of
 429 small scale dissipative processes (Jablonowski and Williamson, 2011) – the
 430 hyperdiffusion schemes do not usually match the symmetry requirements of
 431 the stress tensor needed to ensure the conservation of the angular momentum
 432 (Becker, 2001) – these contributions are produced by the model and will be
 433 taken into account in order to be consistent with the model itself (Johnson,

1997; Egger, 1999; Woollings and Thuburn, 2006). The total dissipation of kinetic energy of the model is therefore $D = D_{phys} + D_{num}$.

Sensible heat in PlaSim is associated with turbulent surface fluxes F_T driven by the temperature difference existing between the lowermost part of the atmosphere and the surface and with numerical vertical and horizontal diffusion (of the same kind of that used for momentum) and dry convection. The material entropy production associated with F_T is:

$$\dot{S}_F = \int F_T \left(\frac{1}{T_a} - \frac{1}{T_S} \right) dA, \quad (15)$$

where T_a is the temperature of the first atmospheric level (where F_T is absorbed thus heating it) and T_S the surface temperature. The material entropy production associated therefore to sensible heat is the sum of the material entropy production due to surface turbulent fluxes, $\dot{S}_{sens}$ and to the other sources of sensible heat (diffusion and dry convection), $\dot{S}_{sens}$, and it reads

$$\dot{S}_{sens} = \dot{S}_F + \dot{S}_{diff}. \quad (16)$$

The total material entropy production of the system is therefore:

$$\dot{S}_{mat} = \dot{S}_{kediss} + \dot{S}_{sens}. \quad (17)$$

The ratio

$$\alpha = \dot{S}_{sens} / \dot{S}_{kediss} \quad (18)$$

is a measure of the degree of irreversibility of the system, which is zero if all the production of entropy is due to the unavoidable dissipation of the mechanical energy (Lucarini et al., 2010). The parameter α introduced above is related to the Bejan number $\mathcal{Be}$ as $\mathcal{Be} = \alpha + 1$ (Paoletti et al., 1989). Systems with large α are instead characterized by high thermal dissipation relatively to the mechanical viscous dissipation and therefore by a higher degree of irreversibility.

5.2. Dissipative properties of circulation regimes

In this section we analyse the dissipative properties of the different circulations described in Sec. 4 as the parameters Ω and τ , and consequently $\mathcal{Ro}$ and $\mathcal{F}_f$, are varied. Sensitivity studies of dissipative properties have been proposed first by Kunz et al. (2008) and then used extensively in Pascale et al. (2011b) and Boschi et al. (2012) as an insightful way to assess the models' tuning and their thermodynamical properties. In the following, we plot quantities in the (Ω^*, τ) plane for practical purposes, and we overplot the values of $\log_{10} \mathcal{Ro}$ and $\log_{10} \mathcal{F}_f$ (Fig. 4 to Fig. 11).

464 *Kinetic energy dissipation and meridional heat transport.* In Fig. 4, the
465 results of the numerical simulations show that for $10^{-2} < \mathcal{Ro} < 1$ and
466 $1 < \mathcal{F}_f < 10^3$ there is the highest total dissipation of kinetic energy, D . We
467 observe a non-trivial dependence on Ω and τ . The most intense dissipation
468 is centered around $\mathcal{Ro} \approx 0.1$ and $\mathcal{F}_f \approx 10^2$ ($\tau = 12$ hours and $\Omega^* = 1$), with
469 $D \approx 0.45 \text{ W m}^{-2}$. This is mainly associated with the dissipation of kinetic
470 energy in the boundary layer, as can be seen in Fig. 5 where D_{phys} is shown.
471 On the base of the discussion in Section 4, we can speculate that at low val-
472 ues of Ω , the baroclinic eddies become larger than the size of the exoplanet
473 (see equation (10) and related discussion) and thus do not develop; at high
474 values of Ω they become too small, convert inefficiently available potential
475 energy into kinetic energy (Hunt, 1979), and dissipate quickly. Furthermore,
476 the surface properties have a dramatic impact on the circulation, as shown
477 also by James and Gray (1986), because the growth rate of the most unsta-
478 ble baroclinic waves is strongly inhibited by horizontal shears (James, 1987)
479 observed, for example, in Fig. 2(e). This explains the drop of D at high
480 $\mathcal{F}_f$ and intermediate $\mathcal{Ro}$. On the other hand, strong drag leads to kinetic
481 energy extraction early in the development of baroclinic eddies. Therefore,
482 the optimal situation is expected for intermediate values of Ω and surface
483 drag. Our results are in agreement with those of Kleidon et al. (2003, 2006),
484 who considered the case $\Omega^* = 1$ only.

485 Moving on to fastly rotating planets, there is a significant decrease of D
486 at low thermal Rossby number ($\mathcal{Ro} < 10^{-2}$) for any value of $\mathcal{F}_f$ (zonostrophic
487 flow, ZN). The strength of the Lorenz energy cycle therefore tends to become
488 more insensitive to the surface properties. Interestingly, also circulations of
489 slowly rotating planets with low drag ($\mathcal{Ro} > 1$, $\mathcal{F}_f > 10^4$, corresponding
490 with the super-rotation regime, see Fig.1(b)) have very weak kinetic energy
491 dissipation. The dissipation rate remains high for slow rotation and for strong
492 drag ($\mathcal{F}_f \leq 0.1$, $\mathcal{Ro} \geq 10$, AS circulations, Fig.1(a)). This is consistent with
493 the fact that in the low rotation, axisymmetric circulations, baroclinicity is
494 mostly absent, and the dissipation of kinetic energy is simply related to the
495 strength of the surface drag, which extracts kinetic energy from the mean
496 flow, thus causing very weak winds near the surface.

497 The meridional heat transport (Peixoto et al., 1991) is in general a very
498 important quantity in planetary atmospheres (Lorenz et al., 2001) and it is
499 associated with the radiative imbalance between high and low temperature
500 regions. The zonal mean of the meridional heat transport $T(\vartheta)$ is worked
501 out at each latitude ϑ by integrating the longitudinally averaged top-of-the-

atmosphere (TOA) radiation budget (Lucarini and Ragone, 2011). A scalar index, MHT , of the meridional heat transport is then defined as half of the difference of the values of the poleward heat transport in the two hemispheres at 30° latitude,

$$MHT = 1/2(Tr(\pi/3) - Tr(-\pi/3)). \quad (19)$$

MHT thus represents the net heat flowing out of the equatorial region through zonal walls placed at 30° .

Overall we observe that the meridional heat transport increases with Ro , in agreement with the results found in Vallis and Farneti (2009). This general feature is due to the inefficiency of the too small baroclinic eddies at high Ω in transporting heat (eq. 10).

Furthermore, it is evident that for intermediate rotation rates ($1/5 \leq \Omega^* \leq 2$) MHT peaks at $\tau \approx 1$ day (≈ 1 PW), that is in the region of baroclinic circulations (Fig. 1(b) and 1(a)), coinciding with the maximum in dissipation (Fig. 4). It is well known in fact that midlatitude eddies constitute a very important mechanism of meridional heat transport (Lorenz, 1967; James, 1994). This is also clear from the zonal mean of the transient eddy flux $\overline{v'T'}$ (not shown), which reaches the highest values $\approx 8 \text{ K ms}^{-1}$ at 900 hPa and 50 N/S for the values of τ maximizing D , compared to 0.5 K ms^{-1} for $\tau = 45$ min (at 700 hPa and 60 N/S) and 4 K ms^{-1} for $\tau = 500$ days (at 1000 hPa and 50 N/S). Just for the sake of comparison, let us note that for earth's circulation $\overline{v'T'}|_{max} \approx 15 \text{ K ms}^{-1}$ at 850 hPa and 50 N/S (e.g. James, 1994). In the slow rotation region ($Ro \approx 10$) we have the largest heat transport (≈ 1.5 PW) at high drag (τ of few hours), which may be explained by lower wind velocities in the lower branch of the Hadley cell (equatorwards motion).

Efficiency and material entropy production. The efficiency diagram (Fig. 7) shows that the highest value of η lay in the intermediate rotation range with values of $\approx 3\%$ in correspondence of the baroclinic and axisymmetric circulations. At low rotations, the high-drag circulations ($\mathcal{F}_f < 1$) are the most efficient. Interestingly, we note that circulations tending toward equatorial super-rotation have a quite substantial drop in efficiency which reduces to $\approx 1\%$. At low Ro the thermodynamic efficiency drops below 1% because of the drastic drop in D associated with the weakening of the Lorenz energy cycle, therefore zonostrophic flows are very inefficient circulation regimes in terms of converting heat into mechanical work. Let us note that although we are dealing with a dry atmosphere, and therefore very different from a moist one (in which the magnitude of the heat losses and gain is much higher,

for example the latent heat gives a positive heating contribution of ~ 80 W m^{-2}), η has comparable values (see e.g. Lucarini et al., 2010) and does not generally exceeds 3%.

The material entropy production terms (eq. (14, 16 and 17)) are shown in Fig. 8-10. Fig. 8 shows the contribution due to thermal dissipation $\dot{S}_{sens}$ (15). This is dominated by $\dot{S}_F$, which accounts for almost 2/3 of $\dot{S}_{sens}$ and is almost independent from Ro , having its highest values for $\tau \sim 3$ days. Such a pattern is explained by a trade-off mechanism between the sensible heat flux, which decreases with τ independently at any Ro (not shown), and the temperature difference between the surface and the near-surface atmosphere, which increases with τ since, due to eq. (7), surface and atmospheres tend to be more decoupled. The entropy production associated with the dissipation of kinetic energy, $\dot{S}_{kedis}$ (Fig. 9) closely follows the pattern of D (Fig. 4) as evident from its own definition (eq. (14)).

The total material entropy production (17) is the sum of the two, so its properties are determined mainly by $\dot{S}_{sens}$ which is generally larger than $\dot{S}_{kedis}$ ($\sim 1 - 2$ times in the at low-intermediate rotation rates, as can be seen in Fig. 11 where the irreversibility parameter α is shown, and up to 10 times for fast rotating planets). The region of highest material entropy production ($\approx 3.5 \text{ mW m}^{-2} \text{ K}^{-1}$) is observed for $0.1 \leq Ro \leq 0.01$ and $10^2 \leq \mathcal{F}_f \leq 10^3$, and generally the whole region of the diagram in Fig. 1(b) with $0.5 \text{ day} \leq \tau \leq 5 \text{ days}$ have large material entropy production. Overall, the material entropy production tends to be fairly low ($\approx 1.5 \text{ mW m}^{-2} \text{ K}^{-1}$) for fast rotation speeds (e.g. $Ro \sim 10^{-3}$) where we have very low values of $\dot{S}_{sens}$ and lower values of $\dot{S}_{kedis}$. Let us note that the portion of the diagram corresponding to super-rotating fluids (SR in Fig. 1(a)) is characterized by very low mechanical and thermal dissipation and therefore very low material entropy production. In this respect super-rotating flows are quite interesting since such circulations are also characterized by very low efficiency. In other terms they seem to have a behavior close to inviscid, non-dissipative fluids (for which $D = 0$ and $\dot{S}_{sens} = 0$ by definition). Mitchell and Vallis (2010) also pointed out some peculiar dynamical properties of super-rotating flows, as for example the fact that the equatorial, strong eastwards jet, once established, do not need eddy-forcing to be maintained. Interestingly, these results make clear that there is no obvious correspondence between the presence of large amount of kinetic energy in the atmosphere and the presence of an intense Lorenz energy cycle to support its generation. This matter has been hotly debated in a rather different scientific context, where the possibility of extracting

576 massive amounts of energy from the atmospheric circulation by wind turbines
 577 is discussed (Miller et al., 2011).

578 A schematic diagram summarising the main thermodynamical properties
 579 discussed so far for the different circulation regimes is shown in Fig. 1(b):

- 580 1. Baroclinic regime (BC): high D , high η , relatively high MHT ;
- 581 2. Super-rotation (SR): low D , low η , low $\dot{S}_{mat}$;
- 582 3. Zonostrophic flow (ZN): low D , low MHT , low η ;
- 583 4. Axisymmetric flow (AS): high MHT and D for $Ro > 1$, high η for
 584 $1 < Ro < 0.1$, low D , MHT and η for $Ro < 0.01$.

585 5.3. Implications for the Maximum Entropy Production Principle

586 In this section we briefly describe our results in the context of the Max-
 587 imum Entropy Production Principle (MEPP, Paltridge, 1975, 1978, 2001),
 588 as this conjecture has gained some momentum also in the planetary science
 589 community (Lorenz et al., 2001; Taylor, 2010). MEPP has been used as
 590 a closure condition for climatic toy-models (Lorenz et al., 2001) or simple
 591 energy balance climate models (e.g. Paltridge, 1975) in order to determine
 592 dynamical quantities as the meridional heat transport. A further, possible
 593 application was shown by Kleidon et al. (2003) and Kunz et al. (2008), who
 594 suggested to use MEPP as a guide for tuning sub-grid motion parameters of
 595 PUMA, an atmospheric general circulation models (Fraedrich et al., 2005).
 596 For example, let us consider the Rayleigh drag constant τ (eq. 6 and following
 597 discussion) depends on the drag coefficient γ_h which in turn depends on both
 598 surface roughness and dynamical quantities. Therefore different values of τ
 599 can be thought of associated with either different surface properties (as done
 600 in the rest of the paper) or to different strengths of the turbulent transfer in
 601 the planetary boundary layer. Following the second interpretation, Kleidon
 602 et al. (2003) showed that the value of τ giving the most realistic atmospheric
 603 state was that maximizing the entropy production of the system. However,
 604 one major criticism that MEPP has encountered is that it does not take into
 605 account the effects of the rotation speed (Rodgers, 1976; Goody, 2007; Jupp
 606 and Cox, 2010). This was related to the criticisms on whether one could
 607 use MEPP to infer the meridional energy transport. In this study we are
 608 in a position to have a broader look on the results of Kleidon et al. (2003)
 609 since a more detailed diagnostics for the dissipative properties and a larger

610 dynamical range for atmospheric circulations are available. Of course our
 611 aim is not, and we do not claim, to prove or disprove MEPP, for which a
 612 rigorous demonstration is still missing (Dewar, 2005; Grinstein and Linsker,
 613 2007).

614 In order to test MEPP, we perform control runs in which the full bound-
 615 ary layer scheme (Louis, 1979; Louis et al., 1981) is employed without the
 616 simplification of Sect. 3.2 (so τ is not prescribed but dynamically determined
 617 depending on the winds and vertical stability). In the following we shall re-
 618 fer to them and to quantities evaluated for such simulations with the label
 619 “BLS” (boundary layer scheme). In BLS simulations the drag coefficient is
 620 consistently determined at each timestep and each grid-point according to the
 621 Monin-Obukhov theory (e.g. Arya, 1988) and not prescribed as a constant
 622 parameter. Since this set up employs a more refined and realistic represen-
 623 tation of the boundary layer physics, we consider it as our “reality” towards
 624 which comparing simulations in which the rougher, tunable τ -scheme is used.
 625 Zonal means of the BLS simulations are shown in Fig. 13 – cross sections of
 626 temperature and zonal winds – and in Fig.14 – meridional streamfunctions –
 627 for simulations for $\Omega^* = 1/10, 1, 8$ respectively. For each Ω^* , we consider τ as
 628 a tunable parameter and select the value $\tau_{max}(\Omega^*)$ maximising $\dot{S}_{mat}$ (which
 629 can be easily visualized in Fig. 10). Furthermore, we take into account also
 630 $\dot{S}_{kedis}$ (Fig. 9), so that we can be informative also on the maximum dissipa-
 631 tion principle (Lorenz, 1967; Ozawa et al., 2003; Schulman, 1977; Pascale
 632 et al., 2011b). We denote with $\tilde{\tau}_{max}(\Omega^*)$ the values of τ maximising $\dot{S}_{kedis}$.
 633 As can be seen in Fig. 9-10, τ_{max} and $\tilde{\tau}_{max}$ differ mostly for $\Omega^* \leq 1/2$ (where
 634 the maximum dissipation steady states occur for τ of few hours) whereas
 635 they are mostly the same (1 day) for $\Omega^* > 2$ days ($\tau \approx 1$ day).

636 In Fig. 12(a) and 12(b) we compare $\dot{S}_{mat}(\Omega^*; \tau_{max})$ and $\dot{S}_{kedis}(\Omega^*; \tau_{max})$
 637 (dashed line) with $\dot{S}_{mat}^{BLS}(\Omega^*)$ and $\dot{S}_{kedis}^{BLS}(\Omega^*)$ respectively (continuous lines).
 638 On the same diagrams we also show the same quantities for $\tau = 0.1 \tau_{max}(\Omega^*)$
 639 (dotted line) and $\tau = 10 \tau_{max}(\Omega^*)$ (dotted-dashed line) in order to provide
 640 an indication of the sensitivity of $\dot{S}_{mat}$ and $\dot{S}_{kedis}$ with respect to τ_{max} . The
 641 MEPP estimate of $\dot{S}_{mat}$ slightly overestimate the values obtained in controls
 642 runs ($\leq 5\%$) but, impressively, captures fairly well the dependence on Ω^* .
 643 Similarly, the values of $\dot{S}_{kedis}$ obtained for τ_{max} compare relatively well with
 644 the ones obtained in the controls runs. Circulations corresponding to τ_{max}
 645 are indeed fairly similar to BLS circulations, as can be seen by comparing
 646 Fig. 13(a,b,c) with Fig. 2(b,e,h) and Fig. 14(a,b,c) with Fig. 3(b,e,h).

647 When the values of $\tilde{\tau}_{max}(\Omega^*)$ associated with the maximum of $\dot{S}_{kedis}$ is

648 instead taken into account (Fig. 12(c)-12(d)), we observe that $\dot{S}_{mat}(\Omega^*, \tilde{\tau}_{max})$
649 provides again a quite good estimate of $\dot{S}_{mat}^{BLS}$, with a slight underestimate
650 ($\approx 9\%$) for $\Omega^* < 1/2$, due to the fact that for such values of the rotation
651 rate $\tilde{\tau}_{max}$ bends towards smaller τ where $\dot{S}_{mat}$ tends to decrease (Fig. 10).
652 More unsatisfactory is $\dot{S}_{kediss}(\Omega^*, \tilde{\tau}_{max})$ again for $\Omega^* < 1/2$, with a difference
653 of about 16% with respect to $\dot{S}_{kediss}^{BLS}$.

654 In the end, both maximum entropy production and maximum dissipa-
655 tion principle provide fairly reasonable estimates of $\dot{S}_{kediss}^{BLS}$ and $\dot{S}_{mat}^{BLS}$, with
656 the maximum entropy production one having better skills at low Ω^* . The
657 quasi-equivalence of the the two methods is due to the fact that, for such
658 simulations, both $\dot{S}_{mat}$ and $\dot{S}_{kediss}$ have their maxima in the (Ω^*, τ) almost in
659 the same regions. These results seem to confirm, in a relatively large range
660 of dynamical regimes, the possibility of using MEPP in its weak form, as a a
661 guide for tuning sub grid parameters associated with turbulent motions, as
662 indicated by Kleidon et al. (2003).

663 6. Conclusions

664 Stimulated by the ongoing development of exoplanet sciences, in this
665 study we have investigated the nonequilibrium thermodynamic properties
666 (kinetic energy dissipation, material entropy production, efficiency, merid-
667 ional heat transport) of optically-thin, non-condensing planetary atmospheres
668 at different values of the thermal Rossby number $\mathcal{Ro}$ and the Taylor number
669 $\mathcal{F}_f$ through a systematic variation of the rotation rate Ω and surface drag
670 time constant τ . The most relevant achievement of this study has been the
671 characterization of the nonequilibrium properties of the different circulation
672 regimes (axisymmetric, super-rotation, baroclinic, barotropic, zonostrophic)
673 obtained with numerical simulations with some interesting connection to the
674 Maximum Entropy Production Principle (MEPP).

675 Slowly rotating planets ($\mathcal{Ro} > 1$) circulation are mostly Hadley cell-
676 dominated but tend to equator; super-rotation for $\mathcal{F}_f > 10^5$. For interme-
677 diate rotation rates ($1 < \mathcal{Ro} < 0.01$) an axisymmetric ($\mathcal{F}_f < 10$), baroclinic
678 ($10 < \mathcal{F}_f < 10^5$) and barotropic ($\mathcal{F}_f > 10^5$) regime are found. At high
679 rotation rates ($\mathcal{Ro} < 0.01$) circulations are characterized by multiple jets
680 (zonostrophic) for $\mathcal{F}_f > 10^4$.

681 The baroclinic regime has high values of D and MHT since midlatitude
682 baroclinic waves provide a very effective way to convert available potential
683 energy into mechanical kinetic energy and transport energy from low to high

latitudes. Such mechanism is inhibited by strong barotropic shears characterizing the barotropic regime and therefore both D and MHT experience lower values. The axisymmetric regime has different thermodynamic properties depending on the value of $\mathcal{Ro}$ at which it is realised. For $\mathcal{Ro} > 1$, a very intense Hadley cell develops associated with high MHT and D ; for $1 < \mathcal{Ro} < 0.1$ such quantities are weaker but circulations are more efficient in converting heat into mechanical work (high η); at faster rotation speeds ($\mathcal{Ro} < 0.01$) a dramatic drop in D , MHT and η is observed. A very interesting case is that of circulation approaching equatorial super-rotation ($\mathcal{Ro} \leq 10$, $\mathcal{F}_f > 10^5$), for which low D , low η , low $\dot{S}_{mat}$ occurs, thus showing a behavior close to inviscid, non-dissipative fluids (for which $D = 0$ and $\dot{S}_{sens} = 0$ by definition). Zonostrophic flows low, typical of fast rotating, low surface drag planets, have a very weak atmospheric energy cycle (low D), are very inefficient in converting potential energy into work and have very low meridional heat transport MHT , therefore showing a temperature profile close to the radiative-convective equilibrium (which by definition has $MHT = 0$).

The thermal dissipation $\dot{S}_{sens}$ is instead fairly insensitive to $\mathcal{Ro}$ and is determined mainly by the timeconstant τ , due to a trade-off mechanism between the temperature difference and the heat flux.

Moreover, we have shown that the possibility of applying MEPP in its weak form, e.g. as a tool for providing guidance in tuning subgrid scale, seems to work relatively well in the range of values of the rotation rate considered in this study, thus extending the results obtained by Kleidon et al. (2003) when considering the terrestrial rotation rate only. Interestingly, there is broad agreement between what prescribed by applying MEPP and the maximum dissipation principle.

This is a first preliminary study for a special case of dry atmosphere. The presence of the hydrological cycle has a huge effect on the circulation and on the energetics and would be definitely worth investigating. Another issue is the role of the surface heat capacity, which would also deserve a systematic investigation. Furthermore, thermodynamic and dynamical properties of slowly rotating planets, e.g. from $\Omega^* = 1/10$ up to phase-locked planets, are still poorly known and would deserve more investigation too.

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Table 1: Parameters and symbols list

parameter/symbol	explanation	value
Ω_E	Earth's rotation rate	$7.29 \cdot 10^{-5} \text{ rad}^{-1}$
c_d	specific heat of dry air	$1004 \text{ J kg}^{-1} \text{ K}^{-1}$
c_{pw}	specific heat of mixed layer model	$4180 \text{ J kg}^{-1} \text{ K}^{-1}$
g	gravitational acceleration	9.81 m s^{-2}
ρ_w	ocean water density	1030 kg m^3
h_{ml}	mixed layer depth	5 m
C_{slab}	slab-ocean areal heat capacity	$10^{-7} \text{ J K}^{-1} \text{ m}^{-2}$
α_s	surface albedo	0.2
S_0	solar constant	1365 W m^{-2}
a	planet's radius	6300 km
$\mathcal{R}o$	thermal Rossby number	
$\mathcal{F}_f$	"frictional" Taylor number	
ASR	absorbed stellar radiation at TOA	
OLR	outgoing long wave radiation at TOA	
F_T	surface sensible heat flux	
F_{SW}^{toa}		
F_{SW}^{surf}		
F_{LW}^-		
γ_h	heat transfer coefficient	
γ_D	drag coefficient	
MHT	meridional heat transport index	
L_R	Rossby deformation radius	
N	buoyancy frequency	
α	irreversibility parameter	

Figures' captions

- Figure 1

1(a) Schematic diagram of the $(\mathcal{F}_f, \mathcal{R}o)$ parametric space spanned in this study. Overplotted are the values of Ω^* (dashed-dotted) and τ (dotted). We have schematically sketched the boundaries between different circulation regimes found for dry PlaSim on the base of the circulations (AS, axisymmetric; BC, baroclinic; BT, barotropic; ZN, zonostrophic; SR, super-rotation). Circles, pentagons and triangles represent the simulations performed with $\Omega^* = 0.1, 1, 8$ respectively (see Fig.3 and 2). 1(b) The same regime diagram is summarizing schematically the properties of kinetic energy dissipation (continuous line, high and low D), meridional energy transport (dotted-dashed line, high MHT), thermal material entropy production (dotted line, high and low $\dot{S}_{sense}$), efficiency (dashed line, high and low η).
- Figure 2

Zonal winds and temperature for $\Omega^* = 1/10$ ($\tau = 2700s$ (a), 1 day (b), 500 days (c)), $\Omega^* = 1$ ($\tau = 2700s$ (d), 1 days (e), 500 days (f)), $\Omega^* = 8$ ($\tau = 2700s$ (g), 1 days (h), 500 days (i)).
- Figure 3

As in Fig.3 but for the meridional mass streamfunction (units 10^9 Kgs^{-1}).
- Figure 4

Total kinetic energy dissipation; overplotted (as in all the following plots) are the values of $\log_{10} \mathcal{R}o$ (dashed) and $\log_{10} \mathcal{F}_f$ (dotted).
- Figure 5

Contribution to the total kinetic energy dissipation due to parametrizations representing boundary layer stresses and gravity wave drag, D_{phys} .
- Figure 6

Atmospheric meridional energy transport index MHT .
- Figure 7

Carnot efficiency η .

- 1063 • Figure 8
- 1064 Entropy production associated with surface sensible heat flux. Units
- 1065 in $10^{-3} \text{ W m}^{-2} \text{ K}^{-1}$.

- 1066 • Figure 9
- 1067 Material entropy production associated with dissipation of kinetic en-
- 1068 ergy. Units in $10^{-3} \text{ W m}^{-2} \text{ K}^{-1}$.

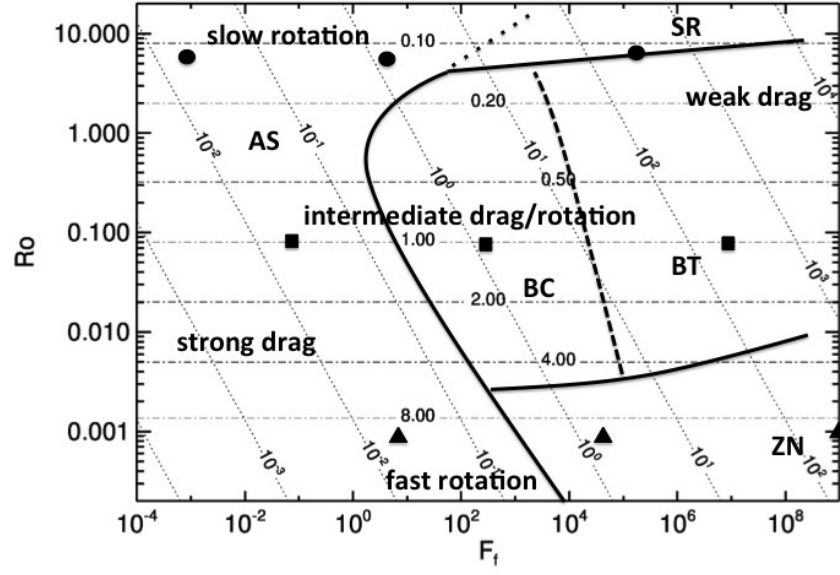
- 1069 • Figure 10
- 1070 Total material entropy production. Units in $10^{-3} \text{ W m}^{-2} \text{ K}^{-1}$.

- 1071 • Figure 11
- 1072 Irreversibility parameter α .

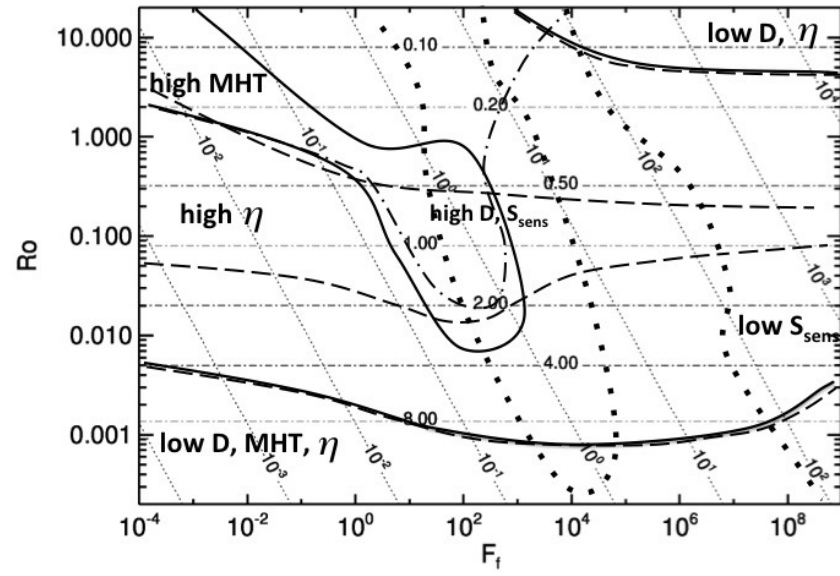
- 1073 • Figure 12
- 1074 $\dot{S}_{mat}$ (12(a)) and $\dot{S}_{kediss}$ (12(b)) for the control runs BLS (continuous
- 1075 line), for $\tau_{max}(\Omega^*)$ maximizing $\dot{S}_{mat}$ (dashed) and for $\tau = 0.1 \tau_{max}$ (dot-
- 1076 ted) and $\tau = 10 \tau_{max}$ (dotted-dashed) days. 12(c)-12(d) Same as in Fig.
- 1077 12(a) and 12(b) but for $\tilde{\tau}_{max}$ maximising $\dot{S}_{kediss}$.

- 1078 • Figure 13
- 1079 Zonal winds and temperature for $\Omega^* = 1/10$ (a), $\Omega^* = 1$ (b) and $\Omega^* = 8$
- 1080 for the BLS simulations.

- 1081 • Figure 14
- 1082 Meridional streamfunction for $\Omega^* = 1/10$ (a), $\Omega^* = 1$ (b) and $\Omega^* = 8$
- 1083 for the BLS simulations.

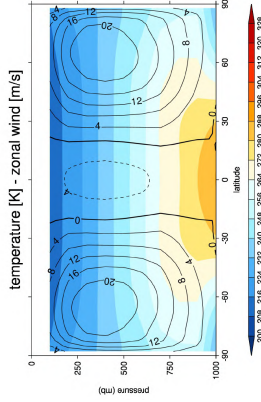


(a)

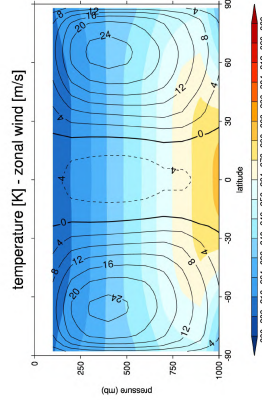


(b)

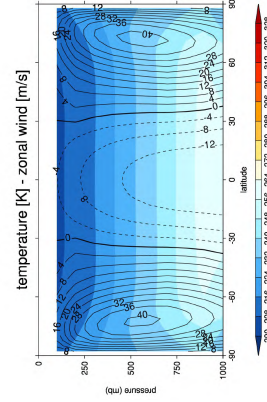
Figure 1:



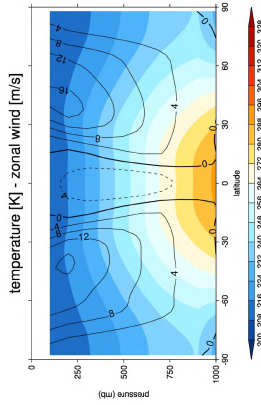
(a) $\mathcal{F}_f = 1.5 \times 10^{-3}$, $\mathcal{R}o = 8$



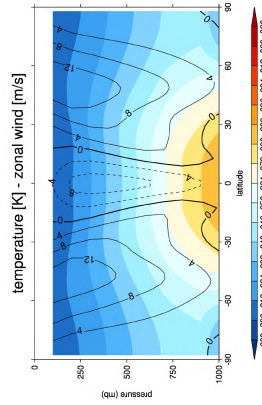
(b) $\mathcal{F}_f = 1$, $\mathcal{R}o = 8$



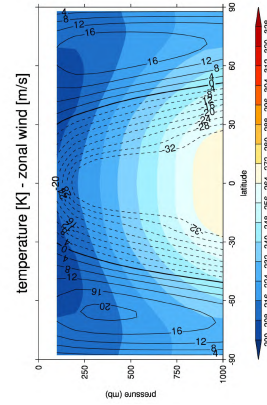
(c) $\mathcal{F}_f = 4 \times 10^5$, $\mathcal{R}o = 8$



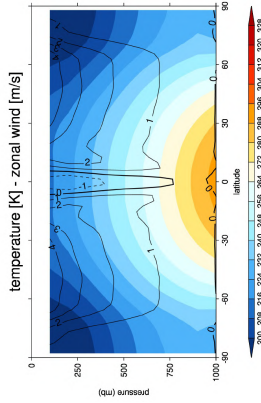
(d) $\mathcal{F}_f = 10^{-1}$, $\mathcal{R}o = 0.08$



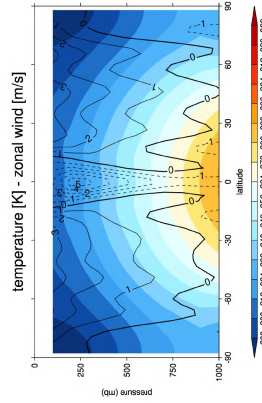
(e) $\mathcal{F}_f = 10^2$, $\mathcal{R}o = 0.08$



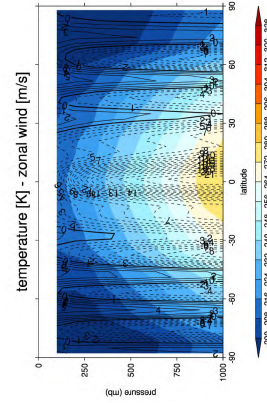
(f) $\mathcal{F}_f = 4 \times 10^5$, $\mathcal{R}o = 0.08$



(g) $\mathcal{F}_f = 10$, $\mathcal{R}o = 10^{-3}$

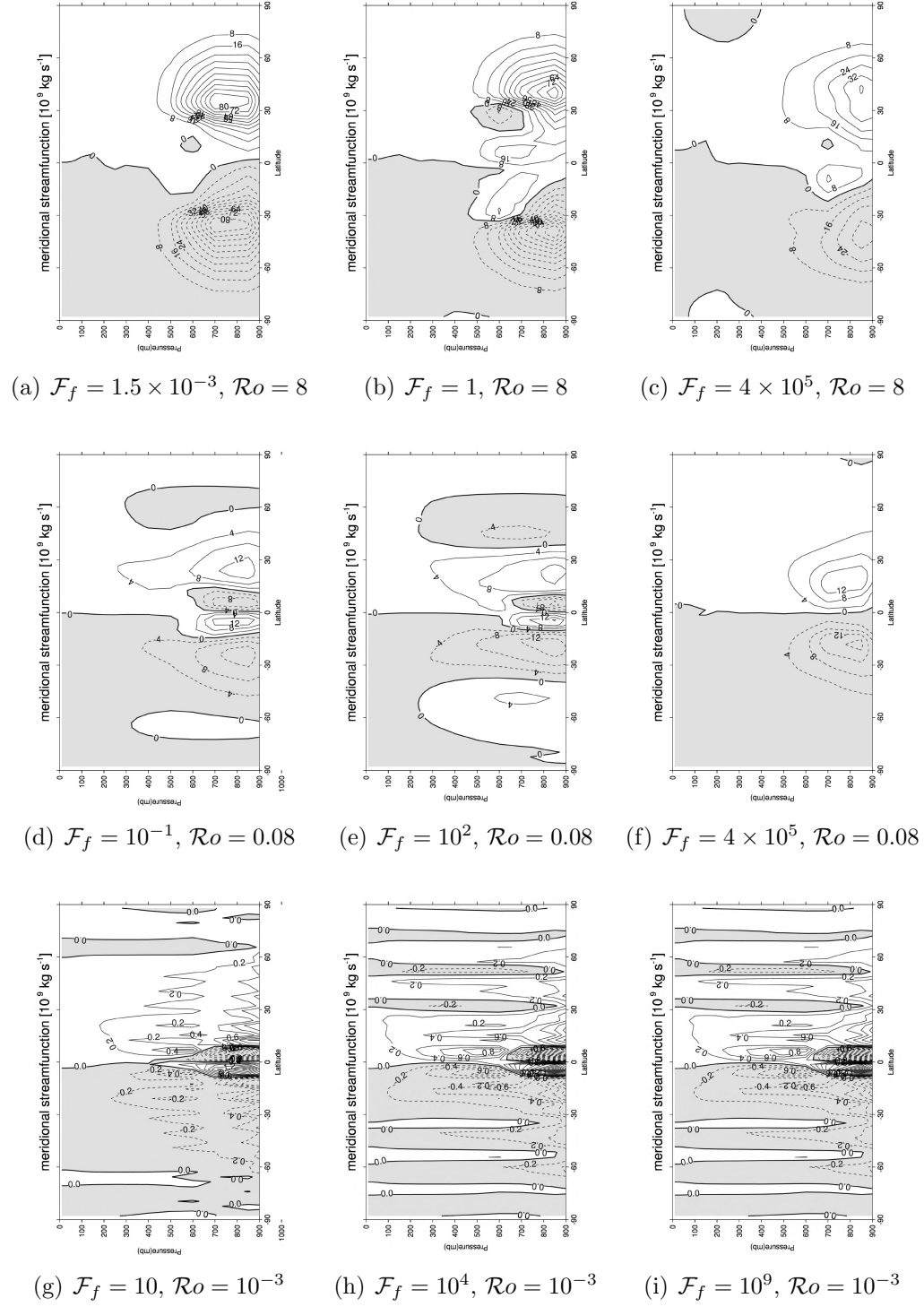


(h) $\mathcal{F}_f = 10^4$, $\mathcal{R}o = 10^{-3}$



(i) $\mathcal{F}_f = 10^9$, $\mathcal{R}o = 10^{-3}$

Figure 2:



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Figure 3:

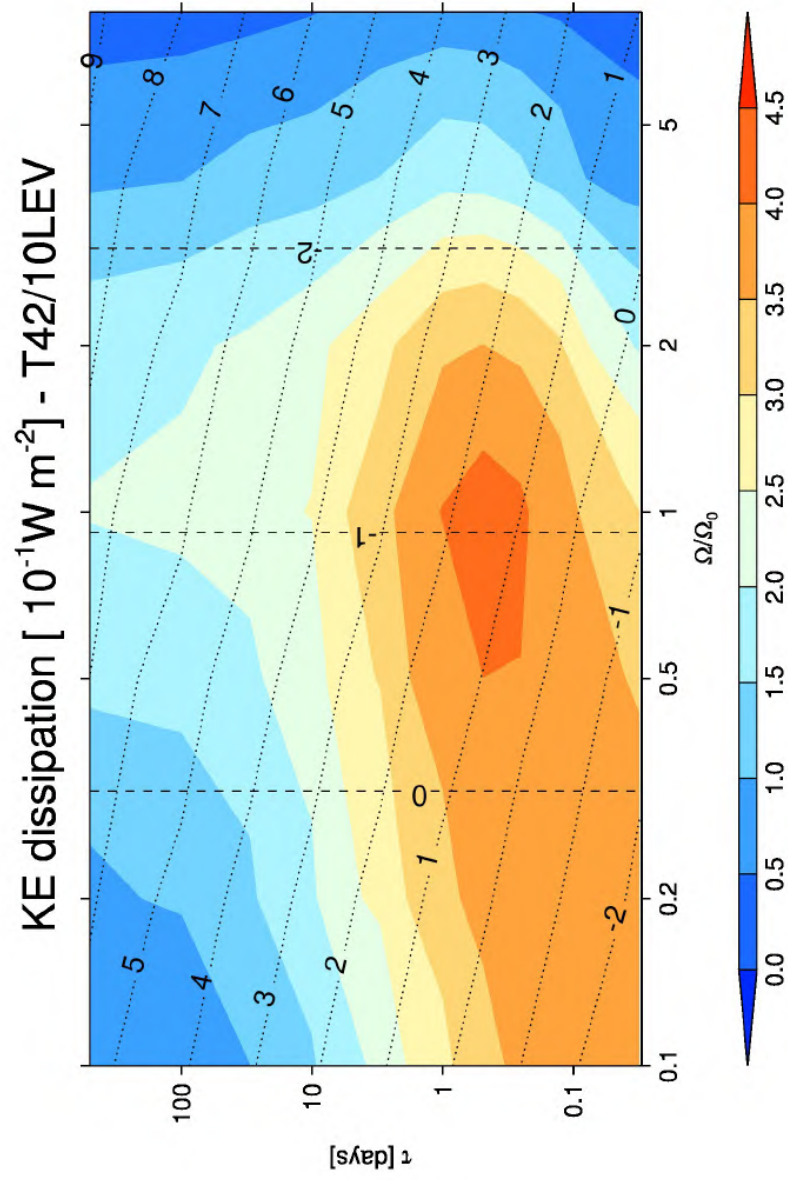


Figure 4:

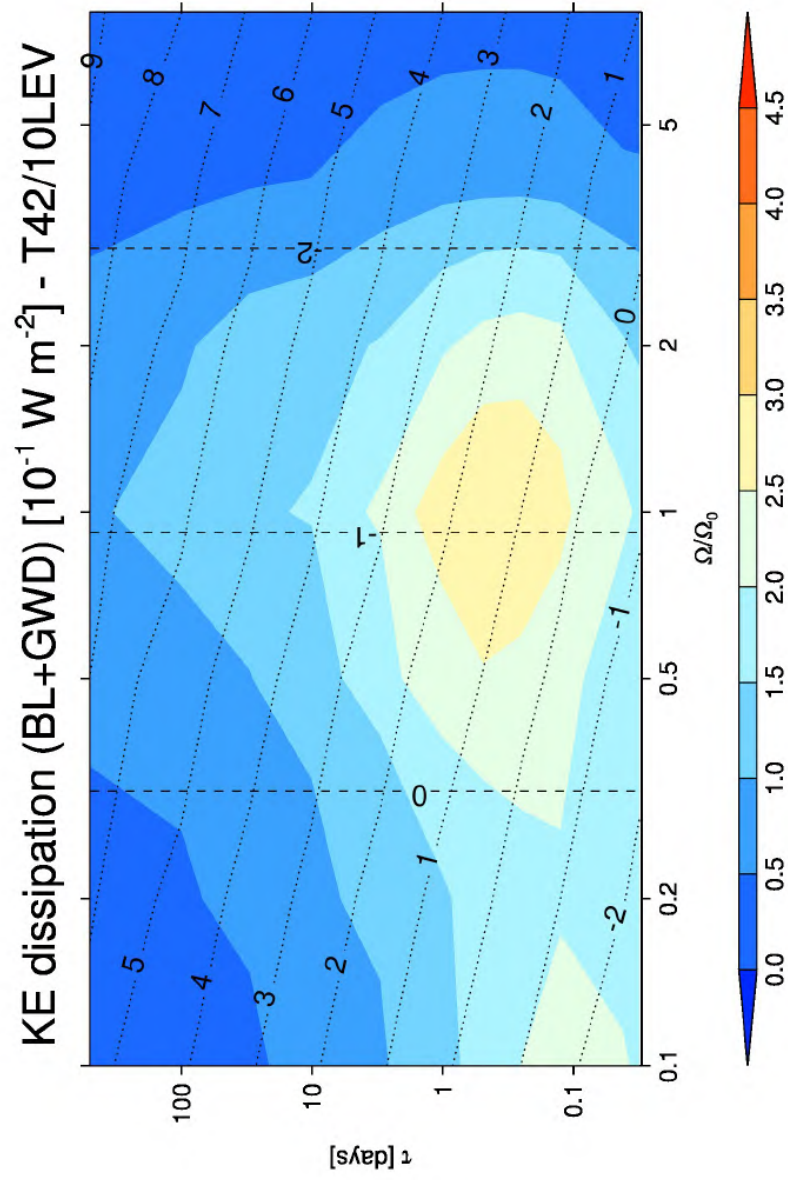


Figure 5:

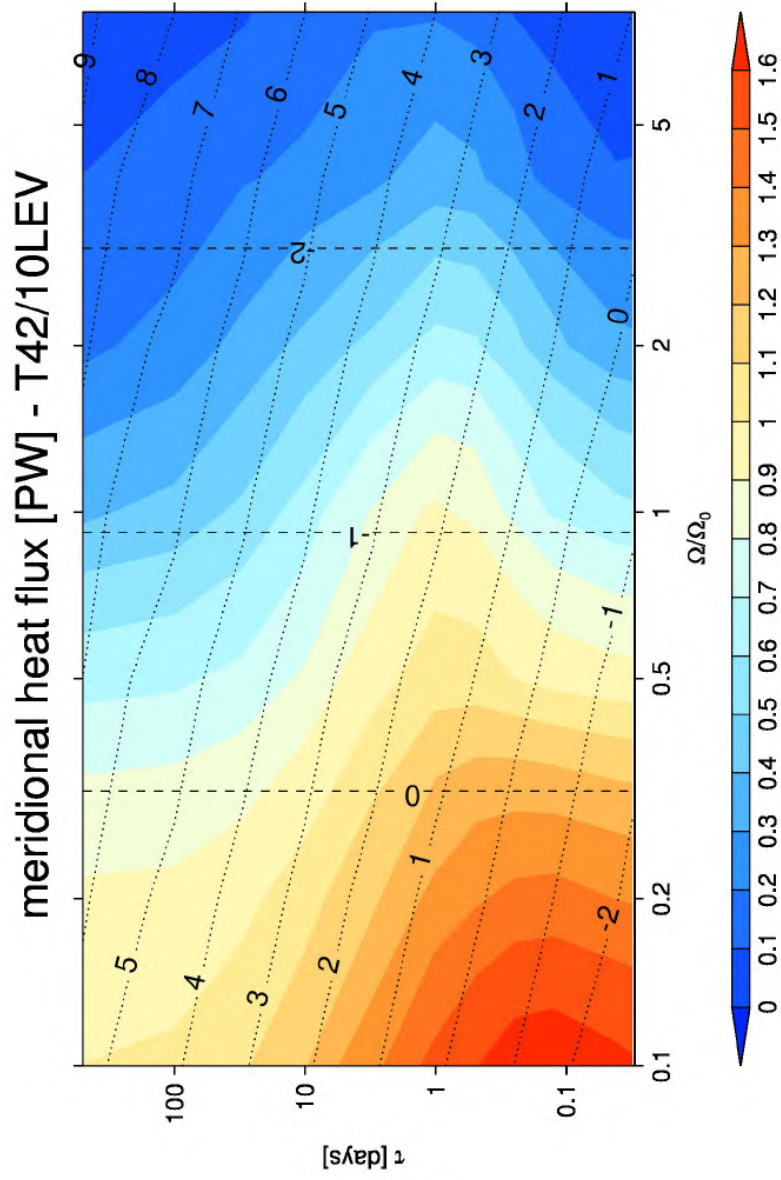


Figure 6:

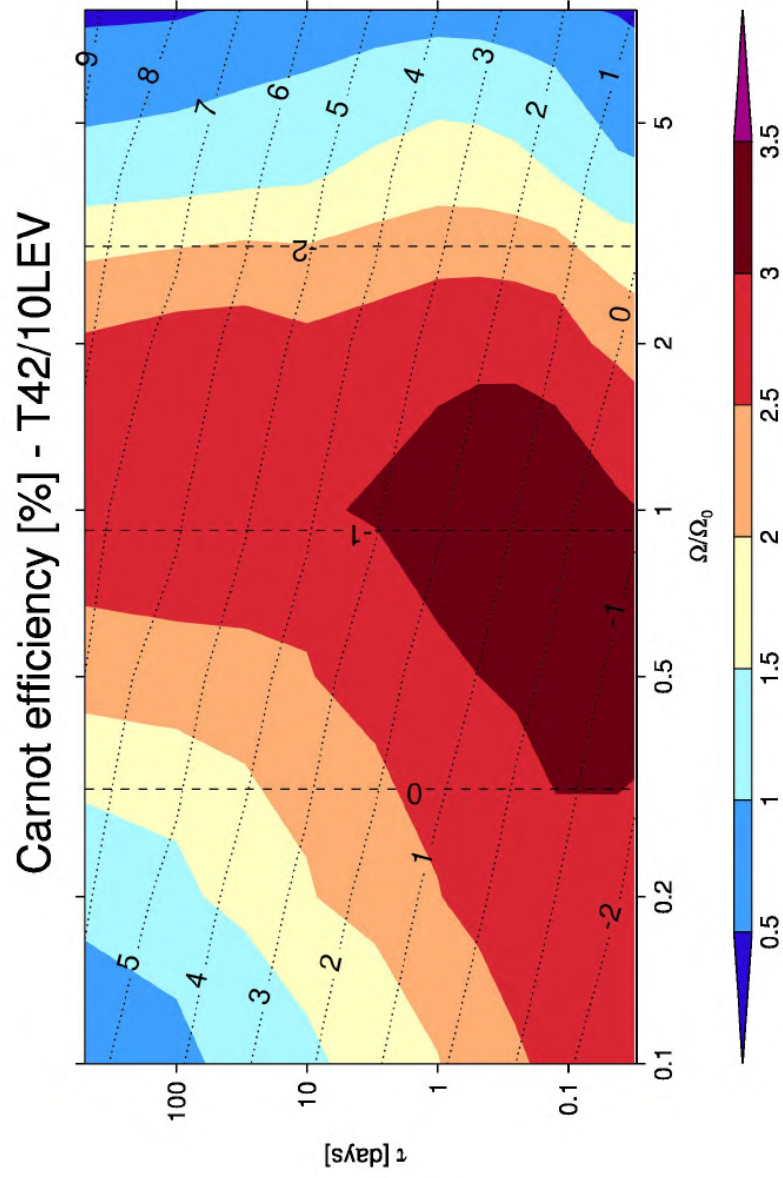


Figure 7:

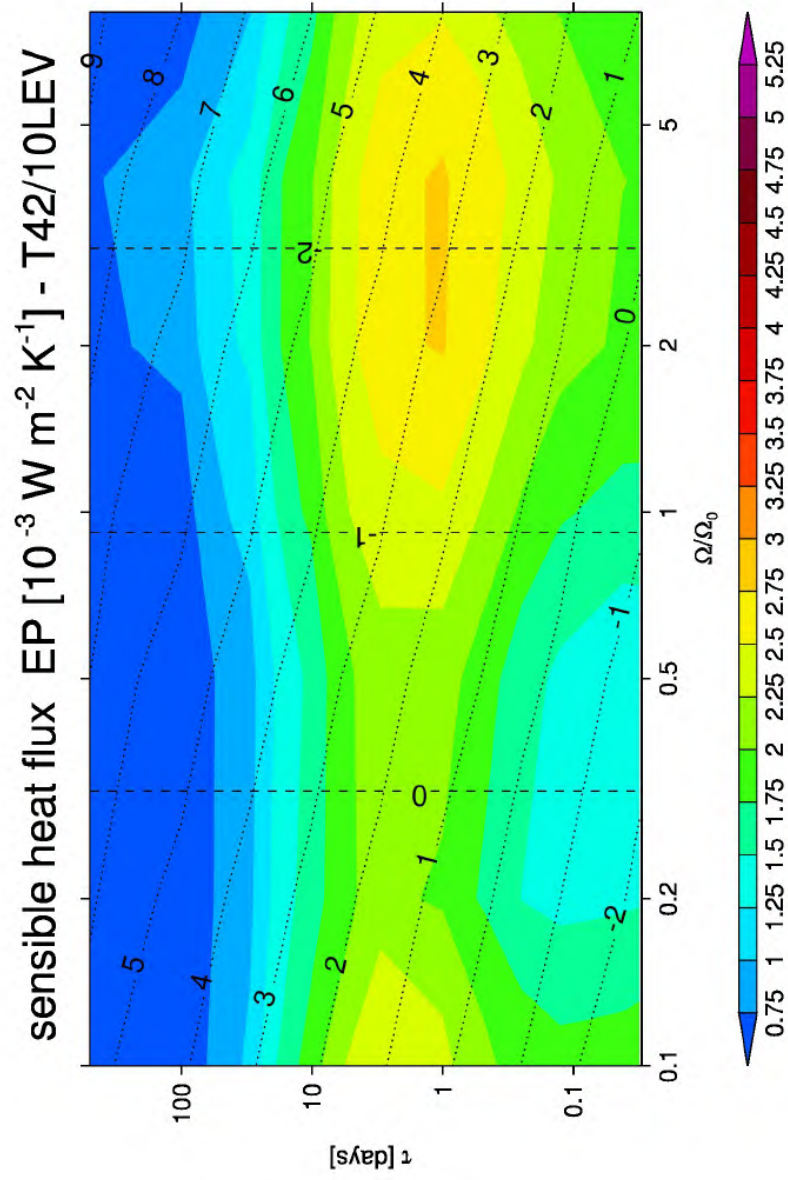


Figure 8:

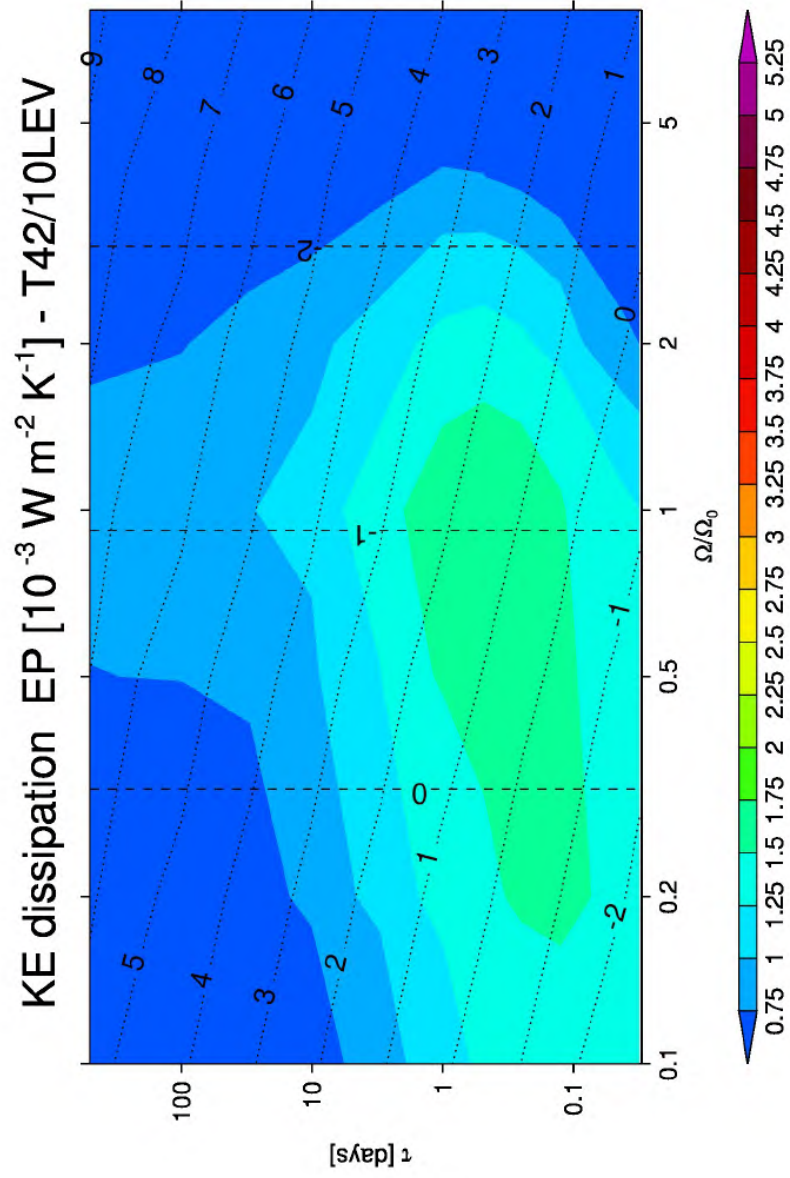


Figure 9:

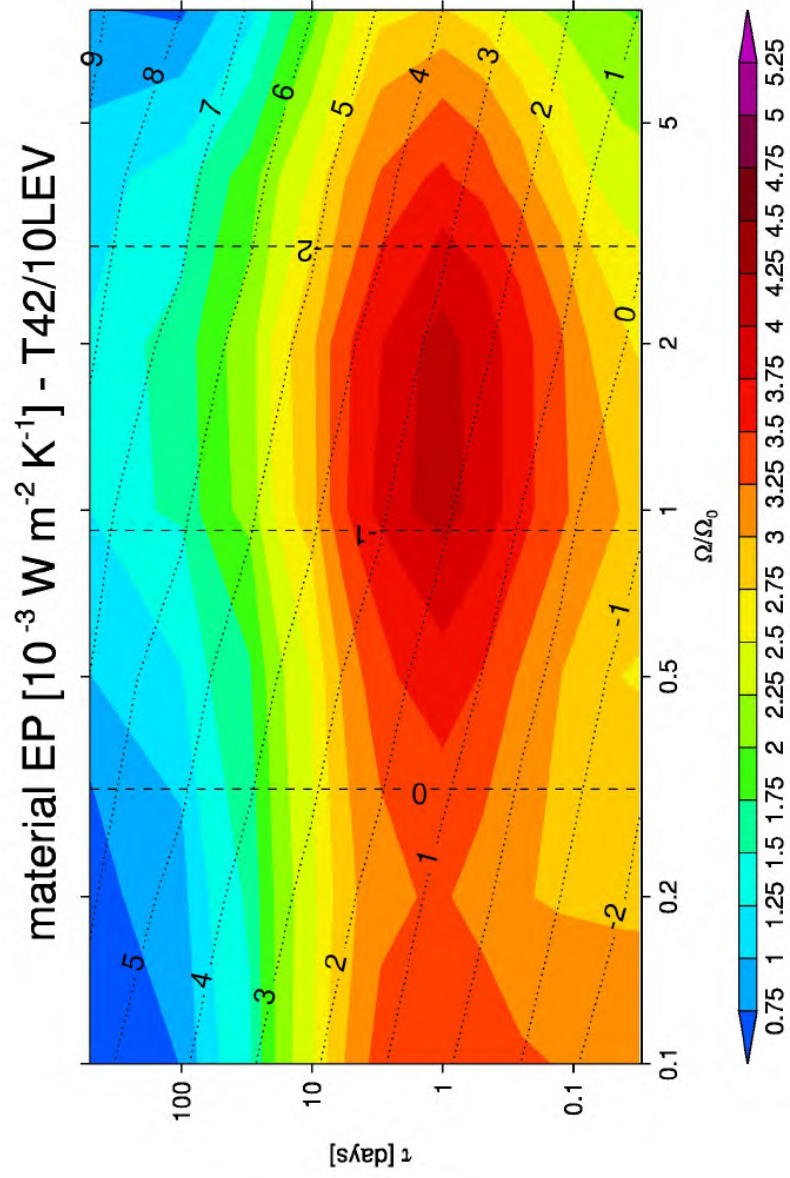


Figure 10:

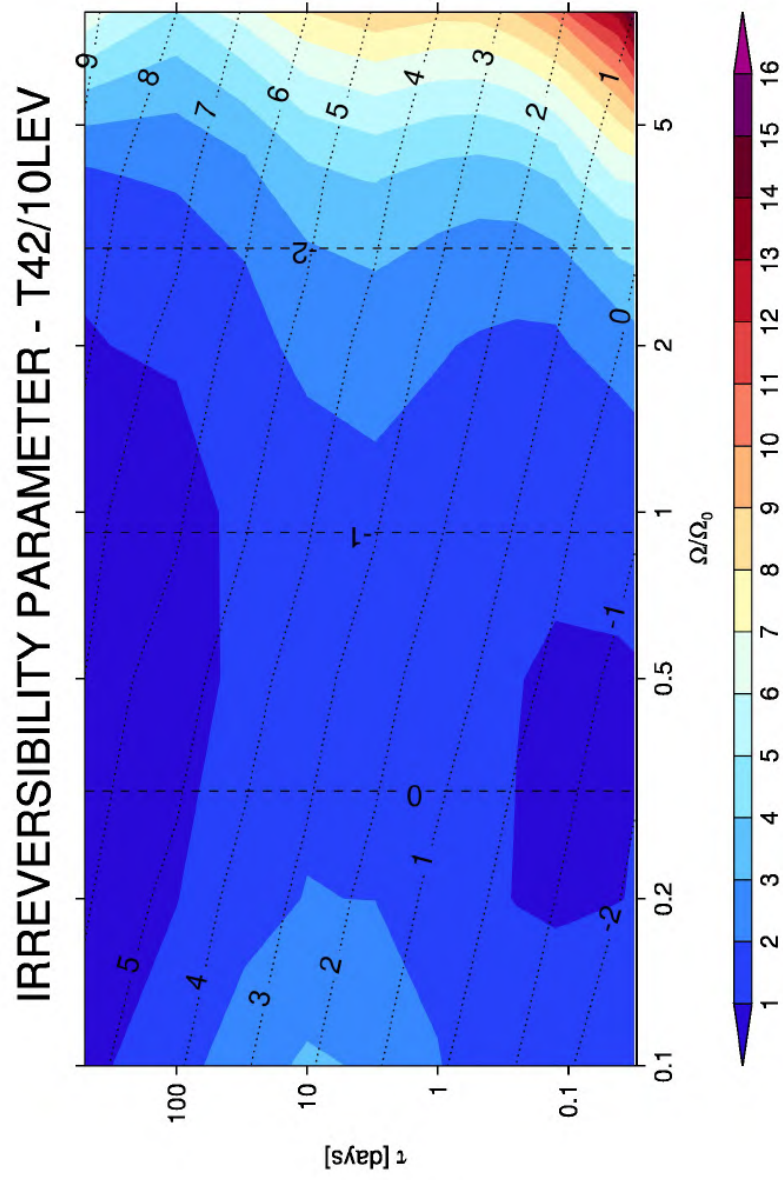
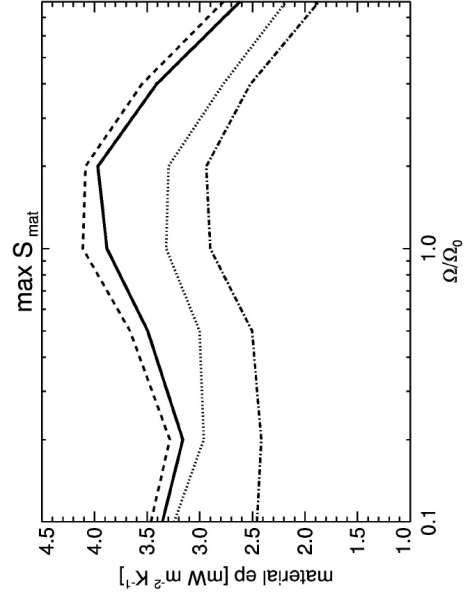
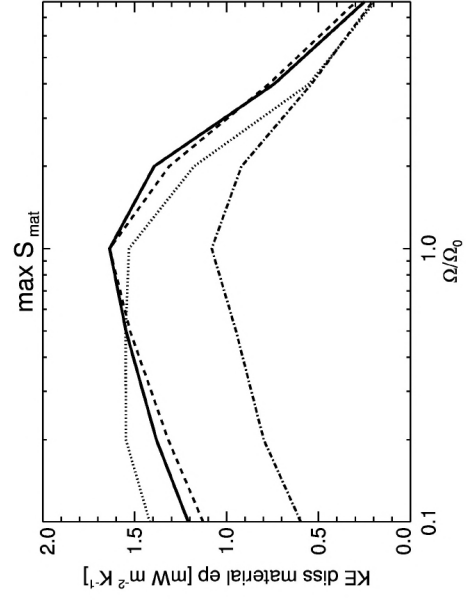


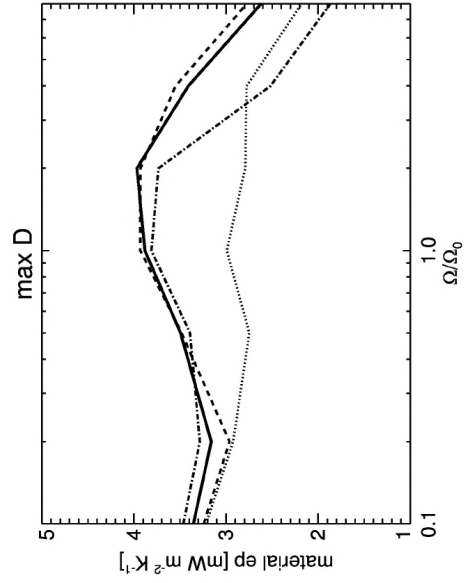
Figure 11:



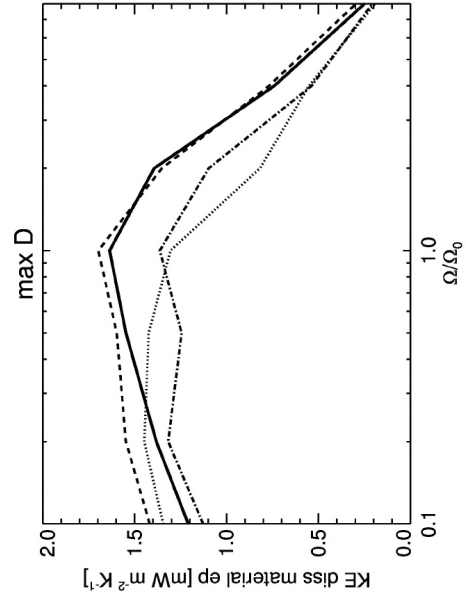
(a)



(b)

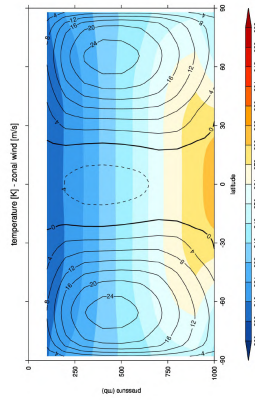


(c)

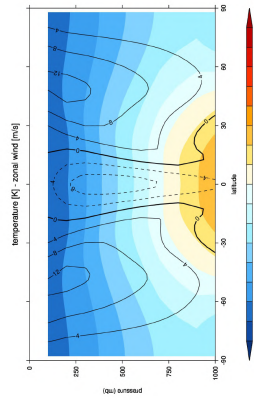


(d)

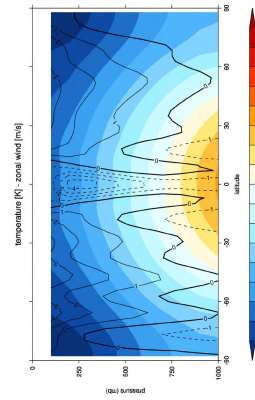
Figure 12:



(a)

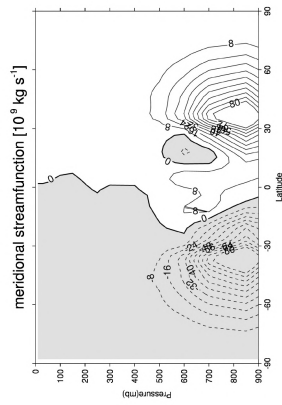


(b)

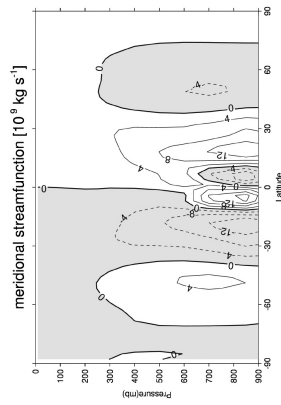


(c)

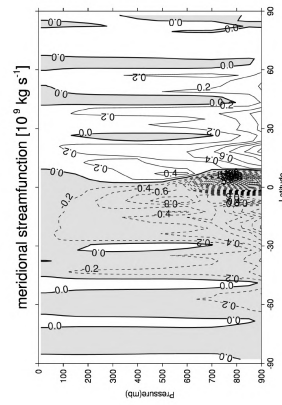
Figure 13:



(a)



(b)



(c)

Figure 14: