5.1 INTRODUCTION

In this chapter, we review the recent progress in understanding the jet structures on exoplanets and for the Sun. For the former, the primary focus is on hot-Jupiters, given that many more observations are available for them presently (see, e.g., Cho, 2008; Showman et al., 2011; Heng and Showman, 2015). Here we make no attempt to be comprehensive, emphasizing only the more robust aspects of observation and numerical modeling—and, in particular, those that relate more directly to jets, the topic of this book. Even so, we apologize at the outset for not including all of the voluminous studies that may be relevant to jets of the exoplanets and the Sun. In addition, the views expressed primarily reflect the authors’ own understanding and biases, and may not be entirely in line or agree with those of the community at large presently. Hence, the readers are encouraged to use the present chapter as a starting point for their own exploration and analyses.

Before embarking on our discussion, the first thing we note is that, because we are still in the early stages of observation and theoretical modeling of exoplanets, not much can actually be said about the presence and strength (let alone the morphology) of the jets on exoplanets and the Sun. In addition, the views expressed primarily reflect the authors’ own understanding and biases, and may not be entirely in line or agree with those of the community at large presently. Hence, the readers are encouraged to use the present chapter as a starting point for their own exploration and analyses.

The situation with HD189733b, one of the best observed exoplanets so far, presents a good illustration of the disagreements. Grillmair et al. (2008) report discrepancies in spectral emission features between two epochs. Charbonneau et al. (2008) find spectral features in secondary eclipse that are inconsistent with measurements obtained earlier by Watkins and Cho (2010) Cho et al. (2015). In many of these studies, the jet is supersonic—even hypersonic, and it exhibits very little zonal (longitudinal) and temporal variabilities. Several theoretical studies, however, have emphasized the possibility of wave-induced variability (e.g., Cho et al., 2003; Watkins and Cho, 2010; Cho et al., 2015) as well as vortex-induced variability (e.g., Thrastarson and Cho, 2010)—stressing the coupling between the variability and the jet, rather than the variability simply being dictated by the jet. Only recently one case of observed variability has not been disputed, thus far. Currently, there is no direct observational evidence of jets on an exoplanet.

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5.2 EXOPLANETS

5.2.1 Current Observations

At present, observations of “surface” (i.e., atmospheric) features on exoplanets are extremely challenging. This is because the planetary disk is not resolved, unlike the planets in the Solar System. Accordingly, interpretation of acquired observations are highly controversial. For example, there are numerous cases of observations of the same planet at different epochs which disagree significantly (sometimes by the same group!). While this may be indicating an intrinsic variability of the planet, the variations are primarily attributed by the observers to variability in the instrument and/or data handling.

The significance of variability for exoplanets is in its putative connection with a high-speed prograde (eastward) equatorial jet, often produced in numerical simulation studies (e.g., Cho et al., 2008; Showman et al., 2009; Rauscher and Menou, 2010; Heng et al., 2011; Mayne et al., 2014; Tsai et al., 2014; Cho et al., 2015). In many of these studies, the jet is supersonic—even hypersonic, and it exhibits very little zonal (longitudinal) and temporal variabilities. Several theoretical studies, however, have emphasized the possibility of wave-induced variability (e.g., Cho et al., 2003; Watkins and Cho, 2010; Cho et al., 2015) as well as vortex-induced variability (e.g., Cho et al., 2003; Thrastarson and Cho, 2010)—stressing the coupling between the variability and the jet, rather than the variability simply being dictated by the jet. Only recently one case of observed variability has not been disputed, thus far. Currently, there is no direct observational evidence of jets on an exoplanet.

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Grillmair et al. (2007). At the same time, Charbonneau et al. (2008) find a different eclipse depth from that reported by Knutson et al. (2007). On the other hand, Sing et al. (2009) obtain measurements at two wavelengths during a transit and report that they do not observe the water signature previously reported by Swain et al. (2008). In Agol et al. (2010), six pairs of transits and eclipses of HD189733b are studied, concluding with “1-σ (68% confident) constraints on the variability consistent with a low variability on the dayside and a high variability on the nightside. Further, observational results of Knutson et al. (2012) differ from those of Charbonneau et al. (2008), with the former work expressing that the values of the latter work require revision.

The same state of affairs exists for HD209458b, another of the best observed exoplanets. Crossfield et al. (2012) report a mean transit depth consistent with previous measurements, showing no evidence of variability in transit depth at the 3% level, but they report a mean eclipse depth somewhat higher than that previously reported for this planet. In a more recent observation by Zellem et al. (2014), they report inconsistency with an earlier study by Knutson et al. (2008), but suggest that the inconsistency is due to variation in the instrument usage between the two studies. In addition, Zellem et al. (2014) revise their previous 4.5 μm measurement of HD209458b’s secondary eclipse (i.e., when the planet is behind its host star from the point of view of the observer) emission downward by approximately 35%, a very large amount. They also report that the dayside brightness temperature is 1499 ± 15 K and the nightside emission temperature is 972 ± 44 K for this planet, a difference of nearly 500 K that is broadly consistent with the setup used in many simulations.

In contrast, Crossfield et al. (2010) observe a non-transiting giant exoplanet, ups And b, and find the phase curve amplitudes at two different epochs to be consistent to 1.7σ certainty. This is an example of non-variability (in time). An example of variability in time that has not experienced dispute so far has recently been put forward for the hot-Jupiter HAT-P-7b by Armstrong et al. (2016). For this planet, Armstrong et al. (2016) report variations in the phase curve peak offset with the peak brightness repeatedly shifting from east side of the planet’s substellar point to the west side. The existence of such a behaviour has been previously shown by Cho et al. (2008) in their simulation study (see Figure 13 therein).

Note that HAT-P-7b is an extremely hot planet with radius 1.4 Rₖ, where Rₖ is the radius of Jupiter (7.1 × 10⁷ m), and with a dayside brightness temperature of ~2,860 K and an equilibrium temperature of ~2,200 K. It transits its host star with a period of ~2.2 days and has been continuously observed for four years by the Kepler telescope at optical wavelengths. It has also been intensively observed at infrared wavelengths with the Spitzer telescope. Hence, in principle, it is one of the best objects to focus on for modeling work. However, observational studies of this planet have suffered from the usual disagreements with the measured amplitudes, as well as with the albedos and temperatures derived from them. The disagreements, here again, have been attributed to the differing wavelengths, datasets, and analysis methods used. Nevertheless, observations appear to agree that the variability occurs on a timescale of tens to hundreds of days for this planet.

Interestingly, several full-phase observations of the exoplanets, HD189733b (Knutson et al. 2007), HD209458b (Zellem et al. 2014), and WASP-43b (Stevenson et al. 2014), report that a “hot spot”—thermally brightest region on the disk—to be eastward of the substellar point. Some circulation model studies (e.g., Showman et al. 2009, Kataria et al. 2013) result in a fast equatorial eastward jet and the hottest region eastward of the substellar point, which is used to claim a rough “agreement” (in the degree of offset). In turn, this has been interpreted in some observational studies as “observational confirmation” that hot-Jupiters have fast (i.e., at least several kilometers per second) zonal winds and a strong equatorial jet—even though there is limited or no information about the latitudinal position of the brightest region and even though other dynamic mechanisms could potentially result in the brightest region being eastward of the substellar point.

For example, Zellem et al. (2014) report that the hot spot is shifted eastward of the substellar point by 40.9 ± 6 degrees, asserting that this is in “rough agreement” [sic] with circulation models predicting equatorial superrotation; this is despite the fact that many of them have shifted hot spots concomitant with a supersonic jet resulting from solving the primitive equations with free-slip boundary conditions (e.g., Showman et al. 2009), which is not physically valid (e.g., Cho et al. 2015). Both the simple distribution and location of the putative hot spot, as well as the robustness of obtained numerical simulation results, have been questioned in several studies (see, e.g., Cho et al. 2008, Thrastarson and Cho 2010, Cho et al. 2015). More recently, Kepler observations show that the light-curves of some hot-Jupiters are “asymmetric”—i.e., for the hottest planets, the light-curve peaks before the secondary eclipse, whereas for planets cooler than ~1900 K, it peaks after the secondary eclipse.

For objects other than hot-Jupiters, the evidence of surface variability is thought to be somewhat more clear and numerous. For example, Demory et al. (2016) report a “4σ-detection” of variability in the dayside thermal emission from the transiting super-Earth 55 Cancri e. Here the signal varies by a factor of 3.7—a very large variation. There is also ample evidence of variability in brown dwarf atmosphere, based on observations in the current literature; see, for example, Artigau et al. (2009), Radigan et al. (2012), and Apai et al. (submitted). It is common, according to the survey of Buenzli et al. (2014). However, no direct observational evidence of jets on super-Earths or brown dwarfs have been obtained so far as well.

As already noted, although the disagreements have been attributed to variability in the instrument and data handling, McCullough et al. (2014) argue that the slope in the short wave spectral region observed could be due to starspots. Note that the unocculted spot coverage is unknown. However, starspots do not account for the lack of the alkali metal wings in the data or for the fact that the slope appears at a higher effective altitude than the 1.4 μm water feature, bringing into question the role of the spots.
5.2.2 Review of Recent Simulations

In general, observations currently rely heavily on simulations for interpreting the obtained data. Unfortunately, as in observations, current simulation studies are not free from discrepancies and disputes—although arguably less than in observations. Thus far, simulation results roughly fall into two general categories: i) those which exhibit supersonic equatorial jets, with hot spots shifted eastward of substellar point by some finite angular distance, and no variability (e.g., Showman et al. 2009); and, ii) those which exhibit subsonic equatorial jets in either direction with complex distribution of, not-necessarily steady, hot (and cold) regions, and often with pronounced variability and noticeable dependence on initial condition (e.g., Cho et al. 2008; Thrastarson and Cho 2010). Note that the two categories are not “mutually exclusive”. For example, an eastward-shifted hot spot can also be present with a subsonic equatorial jet.

The apparent dichotomy is partly due to the use of different initial, boundary, and forcing/drag conditions—as demonstrated by Cho et al. (2015). An example is given in Figure 5.1 which shows time-averaged, zonal-mean zonal velocity $\bar{u}$ from simulations initialized with a 1000 m s$^{-1}$ east or west equatorial jet. In all the simulations, strong thermal damping is applied in the upper region. In the simulations of Figure 5.1a, strong momentum drag along with weak thermal damping are applied in the lower region; and, in the simulations of Figure 5.1b, momentum drag and thermal damping are not applied in the lower region. Note that $\bar{u}$ is essentially the same regardless of the initial jet direction, in the former case; however, $\bar{u}$ is strongly dependent on the initial jet direction, in the latter case. Note also that the behavior in Figure 5.1a is in good agreement with many previous studies whereas the behavior in Figure 5.1b is not, clearly demonstrating that a source of disagreement in simulations is the difference in the dissipation condition employed (Thrastarson and Cho 2011; Cho et al. 2015). Another main source is numerical in origin, resulting in variations even under exactly matching physical setup (Polichtchouk et al. 2014; Cho et al. 2015).

In the simulations of Figure 5.1 a prograde equatorial jet appears to be generic—although its morphology and peak amplitude can be markedly different between different simulations. Showman and Polvani (2011) address the aforementioned category i) behavior, focusing on the generation of the equatorial jet. They argue—following, e.g., Wu et al. (2001)—that the jets result from the interaction of the mean flow with standing Rossby waves induced by differential (day–night in the exoplanet case) thermal forcing. In the proposed scenario, Rossby waves develop phase tilts that pump eastward momentum from high latitudes to the equator, inducing equatorial superrotation. An analytic theory is presented to demonstrate this mechanism as well as a two-dimensional (shallow-water) model “support” for the 3D numerical simulation results.

However, several caveats should be noted. First, the shallow-water model does not cover the entire range of the parameter regime (e.g., supersonic and layer inhomogeneity) of the category, as have been pointed out explicitly in Cho et al. (2008): the shallow-water model is valid, for example, only under incompressible (i.e., subsonic) condition and cannot be directly compared with 3D simulations exhibiting flows at finite Mach numbers. Second, a potentially significant term is included in the momentum equation of the shallow-water model in Showman and Polvani (2011) to represent the effects of momentum transport between layers, a 3D effect; but, the representation, which is used for terrestrial condition, is not valid. Third, the dependence on the background flow, which could be crucial, but is unconstrained by theory or observations, is also not taken into account in the modeling. Finally, the analytical model shows a weak symmetry breaking, leading to superrotation only in the (longitudinal) average, rather than uniformly over the entire equator (as in the 3D simulations): there is a considerable leap from the shallow-water scenario to the fast multiple kilometer per second equatorial jets seen in the 3D simulations.

Note that bottom drag employed in simulations of Figure 5.1a is physically unrealistic. When such a drag is not
Figure 5.2 Instantaneous longitude–latitude distributions of the flow ($\text{m s}^{-1}$) and temperature (K) fields at log($p$) = \{2.0, 4.0, 5.6, 7.0\}, where $p$ is the pressure in Pa, from the simulations in Figure 5.1b. The reference flow vectors are shown at the bottom right in each panel, and the temperature ranges for each row are shown at the right; the projection is cylindrical-equidistant and centered on the substellar point, (longitude, latitude) = (0, 0). Both the flow and temperature distributions are complex, with multiple irregular hot (cold) regions often situated away from the substellar (antistellar) point. Comparing the two simulations, the instantaneous fields are significantly different, except at low $p$ level—as expected, given the very short radiative cooling time there: the cooling time is proportional to $p / T^4$, when the temperature perturbation of the region associated with the cooling time is small compared to $T$. Note also that the vertical temperature gradient fields in the two simulations are different, which would lead to different emergent heat flux distributions over the globe as well as a modification of the cooling time (e.g., Cho et al. 2008).
In their study, Komacek and Showman (2016) focus on explaining the trend of day-night temperature contrast on hot-Jupiters with scaling arguments and MITgcm simulations in the cubed-sphere grid configuration. They report that the full-phase infrared light curves of low-eccentricity hot Jupiters show a trend of increasing dayside–nightside brightness temperature difference with increasing equilibrium temperature. In the simulations, they use the Newtonian cooling scheme as well as a linear frictional drag with two components—a “basal” [sic] drag and a crude parameterization of Lorentz force drag, following Perna et al. (2010); here the basal drag is strongest at bottom of the computational domain and linearly decreases upward and the “Lorentz drag”, applied at the bottom, is spatially constant but varies between different simulations. The employed simple drag representation for the species Lorentz drag has been criticized by Koskinen et al. (2014) and Cho et al. (2015) because of its lack of physical realism for both the upper and lower regions of the modeled domain. Komacek and Showman (2016) also argue that the longitudinal propagation of waves mediates dayside–nightside temperature differences in hot-Jupiter atmospheres. Indeed, wave heating and cooling mechanisms have been discussed explicitly by Watkins and Cho (2010) for gravity waves and Cho et al. (2015) for planetary waves; the latter type of waves have been discussed earlier by Cho et al. (2003). Note that both types of waves can be damped in hot-Jupiter atmospheres by radiative cooling, saturation, and encounters with critical layers—effecting communication over long distances (“teleconnection”), in both the lateral and vertical directions. An example of the interaction and propagation of planetary waves can be seen in Figure 5.3 in which upwardly propagating planetary waves strongly modify the background jet (and associated temperature) structure. In the figure, propagating waves power variability, if 1) vertical resolution near the bottom is increased [cf. Figure 5.3] and Figure 5.3), 2) linear Rayleigh drag is not applied near the bottom of the domain [cf. Figure 5.3a and Figure 5.3]), or 3) sensitivity to initial condition is present [cf. Figure 5.3a and Figure 5.3b].

Parmentier et al. (2016) use the thermal structure from MITgcm simulations to determine the expected cloud distribution and Kepler light-curves of hot-Jupiters. Post-processing the simulation data with plane-parallel radiative transfer and “equilibrium cloud” models, they report that the change from an optical light-curve dominated by thermal emission to one dominated by scattering (reflection) naturally explains the observed trend from positive to negative offset of the hot spot location. They speculate that, for the cool planets, the presence of an asymmetry in the Kepler light curve is indicative of the cloud composition, because each cloud species can produce an offset only over a narrow range of effective temperatures. They also add that the cloud composition of hot-Jupiters likely varies with equilibrium temperature, suggesting that a transition occurs between silicate and manganese sulfide clouds at a temperature near 1600 K, analogous to the L/T transition on brown dwarfs: the cold trapping of cloud species below the photosphere naturally produces such a transition and predicts similar transitions for other condensates, including TiO. Of course, all of these results—as well as those below—depend on the accuracy of the jet and thermal structures obtained in the GCM simulations.

Kataria et al. (2016) present results from a MITgcm simulation study of nine hot-Jupiters that make up a large transmission spectral survey using the Hubble and Spitzer Space Telescopes. These observations exhibit a range of spectral behavior over optical and infrared wavelengths, suggesting perhaps diverse cloud and haze distributions. By utilizing the specific system parameters for each planet, they explore the parameter-space spanned by planet radius, surface gravity, orbital period, and equilibrium temperature. In this study, they show that their model “grid” recovers trends shown in traditional parametric studies of hot-Jupiters, particularly equatorial superrotation and increased day-night temperature contrast with increasing equilibrium temperature. They report that spatial temperature variations—particularly between the dayside and nightside and west and east terminators—can vary by hundreds of degrees, which may imply large variations in Na, K, CO, and CH4 abundances in those regions. They also compare theoretical emission spectra generated from their models to available Spitzer eclipse depths for each planet and find that the outputs from their cloud-free, solar-metallicity models generally provide a good match to many of the datasets.

Mayne et al. (2014, 2017) have adapted the Unified Model (UM) and use it to study HD209458b. They compare their results with test cases, Showman et al. (2009), and available observations. They also focus on the robustness and evolution of a superrotating equatorial jet and its interaction with the deep atmosphere. The UM solves the full 3D Navier-Stokes equations with a height-varying gravity, appropriate for deep atmospheres; these equations are valid for non-hydrostatic atmospheres. In their study, the occurrence of a superrotating, supersonic equatorial jet is robust to changes in various parameters they consider, and over long timescales—even in the absence of strong inner or bottom boundary drag, similar to what has been reported in Cho et al. (2015). Note that Cho et al. (2015) have argued that such a jet is not physical in their simulations, since supersonic jets are produced under hydrostatic conditions with rigid (free-slip) top and bottom boundary conditions. It is also worth noting that shocks are not resolved in simulations in any of the studies. As an aside, Mayne et al. (2017) also find that the jet amplitude is diminished when the equator-to-pole temperature gradient in the deep atmosphere is forced over long timescales, showing dependence on the details of the setup.

Anumaden et al. (2016) extend the model of Mayne et al. (2014) by incorporating a radiation scheme based on the two-stream approximation and correlated-k method with opacities from the ExoMol database (Tennyson and Yurchenko, 2012) and compare their results with observations and Showman et al. (2009). In their study, Anumaden et al. (2016) find a reasonable agreement between observations and both their day-side emission and hot spot offset. But, they report that their night side emission is too large. In addition, although their results are qualitatively similar to...
Figure 5.3 Hovmöller plots of instantaneous zonal velocity $u(t,p)|_{(\phi, \theta)=(0, 30)}$ for $t = [0, 300] \tau$ and $\log(p) = [4.0, 7.3]$; here $\phi$ is the longitude, $\theta$ is the latitude, $\tau$ is the rotation period, and $p$ is the pressure in Pa. The color bars indicate the velocity $(\text{m s}^{-1})$ for each row. Initial jet amplitudes and directions (latter indicated above each plot) are as in Figure 5.1. In a) and b), the resolution is T21L1000 (i.e., max degree and order of 21 each and 1000 vertical layers) and the flow at the bottom is strongly Rayleigh dragged; the vertical levels are equally spaced in $\log(p)$ and $p$ for simulations in a) and b), respectively. In c) and d), the resolution is T85L40 and the flow is not dragged near the bottom; the vertical levels are equally spaced in $\log(p)$. In b), c), and d), vertically propagating planetary waves power variability, if the vertical resolution near the bottom is increased [cf. a) and b)], Rayleigh drag is not applied near the bottom of the domain [cf. a) and d)], or sensitivity to initial condition is present [cf. c) and d)].

those of Showman et al. (2009), they report several quantitative differences: their simulations show significant variation in the position of the hottest part of the atmosphere with pressure, as shown in Cho et al. (2015). This is in contrast to the “significant vertical coherency” reported by Showman et al. (2009). In addition, they also find significant quantitative differences in calculated synthetic observations. In Amundsen et al. (2016) the temperature-pressure profile in their deep atmosphere continues to evolve, without reaching equilibration.

Mendonça et al. (2016) introduce and show benchmark tests for their new GCM, THOR, that solves the 3D non-hydrostatic “Euler equations”. Their model uses an icosahedral grid to address the pole problem (as with the cubed-sphere grid for MITgcm) and a tunable explicit dissipation “spring dynamics” scheme, chosen to ensure stable model integration. Hence, the strength of dissipation does not change with latitude or longitude. Note that, although inviscid equations are solved, substantial dissipation of numerical origin is still present. In addition, split-explicit method is used for the time stepping, together with a horizontally explicit and vertically implicit integration. The model is designed to run on graphical processing units (GPUs) and is a part of the open-source Exoclimes Simulation Platform. The developers have validated their code with the Held-Suarez test case for Earth-like conditions and qualitatively reproduced the results of Menou and Rauscher (2009) and Heng et al. (2011) for hot-Jupiter-like conditions.

Fromang et al. (2016) use RAMSES, a finite-volume shock-capturing code, to also solve the compressible “Euler equations” in the $\beta$-plane; the $\beta$-plane is a tangent plane approximation, retaining the Coriolis term contribu-
ibility also creates velocity fluctuations that propagate upward and steepen into weak shocks at pressure levels of a few mbars (n.b., 1 bar = $10^5$ Pa). They conclude that hot-Jupiter equatorial jars are potentially unstable to both a barotropic Kelvin-Helmholtz instability and a vertical shear instability.

Prior to this study, Polichtchouk and Cho (2012) solve the hydrostatic 3D equations on the full sphere and demonstrate the possibility of baroclinic instability in hot-Jupiter-like conditions; the instability is sensitive to the morphology of the jet and the temperature distribution. In their study, when it occurs, the instability evolves concomitantly with a barotropic instability. The instability naturally transports kinetic energy downward, as part of the equilibration process. Additionally, they show that, for a high-speed subsonic equatorial jet, instability occurs at the flanks (but not at its core), when adequate resolution is present—as confirmed by Fromang et al. (2016), solving the non-hydrostatic 3D equations on the beta-plane.

GCMs solving the full Navier-Stokes equations have also been directly coupled with detailed cloud models. Helling et al. (2016) extend the model used by Dobbs-Dixon et al. (2010) and Dobbs-Dixon and Agol (2013) and focus on aspects of cloud distribution on HD189733b and HD209458b. Helling et al. (2016) report that, in their study, both planets are covered in mineral clouds throughout the entire model domain—independently of differences in hydrodynamic models. Further, the clouds are chemically complex, composed of mineral particles that have a height-dependent material composition and size: therefore, single values to characterize metallicity and C/O ratio are not valid. They also argue that their results concerning the presence and location of water in relation to the clouds explain some of the observed difference between the two planets and that obscuring clouds exist high in the atmosphere of HD209458b, but much deeper in HD189733b.

However, the above interpretation does not uniquely or wholly fit the currently available transit spectra. For example, in the case of HD209458b, a clear atmosphere provides a satisfactory fit to the transit data; see, e.g., Figure 7 of Lavvas et al. (2014). In particular, the Na D line wing and the 1.4 micron water band show up roughly at the same level as the short wavelength slope, indicating that these features probe a similar pressure level. The revised data from Sing et al. (2016) also present a K line wing. On HD189733b, fitting the HST transit spectrum with a clear atmosphere is difficult; see, e.g., Figure 7 of Lavvas and Koskinen (2017). There are no Na D or K line wings and the short wavelength slope lies well above the 1.4 micron water feature. This configuration requires the presence of a high altitude absorber with a small particle size. Indeed, high altitude haze composed of soots can form on HD189733b with abundances sufficient to explain the alleged transit spectrum.

Lee et al. (2016) also use an extended model of Dobbs-Dixon et al. (2010) and Dobbs-Dixon and Agol (2013), focusing on clouds formation under 3D dynamics for HD189733b. The simulation includes the feedback effects of cloud advection and settling, gas phase element advection and depletion, and the radiative effects of cloud opacity. The cloud particles are modelled as a mix of mineral materials, which change in size and composition as they travel through atmospheric thermo-chemical environments. All local cloud properties such as number density, grain size, and material composition are time-dependent. Gas phase element depletion as a result of cloud formation is included in the model. In situ effective medium theory and Mie theory are applied to calculate the wavelength-dependent opacity of the cloud component.

In their study, Lee et al. (2016) find that the mean cloud particle sizes are typically sub-micron (0.01–0.5 μm) at pressures less than 1 bar, with hotter equatorial regions containing the smallest grains. Denser cloud structures occur near the terminators and in the deeper regions (>1 bar). Silicate materials, such as MgSiO3(s), are found to be abundant at mid–high latitudes, while TiO2(s) and SiO2(s) dominate the equatorial regions. Elements involved in the cloud formation can be depleted by several orders of magnitude. The interplay between radiative-hydrodynamics and cloud kinetics leads to an inhomogeneous, wavelength dependent opacity cloud structure with properties differing in longitude, latitude, and depth. This suggests that transit spectroscopy would sample a variety of cloud particles properties (e.g., sizes, composition, and densities). Presumably, such modeling could in principle help to elucidate or constrain information about jets.

Several studies have appeared in the literature very recently and we briefly mention them here. Tremblin et al. (2017), based on a two-dimensional steady-state atmospheric circulation model, model the advection of the potential temperature $\Theta$ due to mass and longitudinal momentum conservation; here $\Theta = T(p_r/p)^\kappa$, where $p_r$ is a reference pressure (often set to 1 bar) and $\kappa$ is the adiabatic index. They argue that longitudinal–vertical mixing implies a larger radius for the planet, reproducing the observed radius of HD209458b. Rogers and McElwaine (2017) show that a dynamo can be maintained in the atmosphere of a hot-Jupiter by conductivity variations arising from strong asymmetric heating from the planets’ host star, independently of the deep-seated dynamo in the planet. They argue that the presence of a dynamo significantly increases the magnetic field strength on the surface. Zhang and Showman (2017) investigate the effects of atmospheric bulk compositions on temperature and wind distributions for tidally locked sub-Jupiter-sized planets, using the MITgcm. Penn and Vallis (2017), using a shallow-water model with time-dependent forcing, confirm the inhomogeneous shallow-water equations study by Cho et al. (2008) that the peak of an exoplanet thermal phase curve is, in general, offset from the secondary eclipse when the planet is rotating. They also consider the inverse problem of constraining planetary rotation rate, a critical parameter for dynamics, from an observed phase curve.

5.2.3 Some Critical Issues

As mentioned, current simulations are in agreement about the small number (usually three) of broad jets expected on synchronized giant exoplanets. However, it is not clear that
the agreement is entirely significant, given that a number of modeling issues remain to be addressed and resolved. The situation should certainly improve as more repeated observations of the same planets are made, which could provide better constraints. Crucial scales and parameters—such as the Rossby deformation scale $L_R$, Rhines scale $L_\beta$, and plasma beta parameter $\beta_p$ (see, e.g., Cho [2008])—are still poorly known, and pose a great challenge for simulation as well as overall understanding; here $L_R$ is a measure of the distance traversed by a propagating gravity wave in one rotation period of the planet $\Omega_p$; $L_\beta$ is the scale at which nonlinear processes is balanced by the gradient of the Coriolis acceleration and can serve as a crude measure of jet width; and, $\beta_p$ is a measure of the speed of the gravity wave relative to the Alfvén wave. Paradoxically (perhaps), the jet strength is not proportional to the separation distance of the planet from its host star (Cho et al. [2008]), at least in the Solar System: there is some evidence that the peak strength of the jets may be inversely proportional to the separation distance and the depth in the atmosphere where the heat deposition occurs.

Studies already clearly show that improved numerical algorithms and models with greater resolution and physical complexity are crucially needed for a better understanding of jets on exoplanets (see, e.g., Polichtchouk and Cho [2012], Polichtchouk et al. [2014], Cho et al. [2015]). One reason for this is because many of the planets are in highly geostrophic, “unbalanced” regimes. Higher resolution will be needed to resolve the narrower jets that result from faster rotation rates expected for synchronized planets closer-in than approximately 0.02 AU, for example. Also, small-scale gravity waves would be even more dominant than in present situations, and would be even more crucial for accuracy. Additionally, ideal and non-ideal magnetohydrodynamic (MHD) simulations will also need to play a greater role, as ionization becomes stronger—even at depths much greater than previously thought (Koskinen et al. [2013], Cho et al. [2015]).

To this point, some simulations have employed a linear, Rayleigh drag as a “crude representation of magnetic drag effects” that putatively stem from thermal ionization. As pointed out in Cho et al. [2015], there are two major concerns with this. First, thermal ionization is insignificant in the modeled region, as temperature is too low and density is too high. This is so even taking into account the low-ionization potential of alkali metals (e.g., K, Na, Ca) because these are species are not present in abundant amounts, assuming solar abundances—i.e., $n_{H^+}/n_e \lesssim 2 \times 10^{-16}$ and $n_{K^+}/n_{H^+} \approx 3 \times 10^6$, where $n_x$ is the x-specie number density and “n” subscript refers to the neutral component. Secondly, even if the ion-induced drag were significant via a non-thermal mechanism, for example, it cannot be represented as a simple isotropic drag-to-rest on the momentum field: ion velocities and the intrinsic field orientation need to be modelled self-consistently for accurate representation (e.g., Koskinen et al. [2010]).

While modeling work to date has covered parts of the presumably relevant parameter-space with different types of physical and numerical models, it is important to note that many non-trivial modeling choices made are often rather similar between different studies. For example, when the models are forced with temperature relaxation, the assumptions about the structure of equilibrium and initial temperature distributions do not reflect the large uncertainty and range of conditions that are plausible. It is not clear if more complex forcing (e.g., incorporating sophisticated radiative transfer), alleviates this uncertainty or even worsen the situation. Likewise, the uncertainty and effects of deeper layers of the atmospheres has only been addressed to limited extent. Usually, in the models, the deeper layers contain a large mass of slowly evolving air that is poorly resolved and almost always assumed to start from rest. But in reality, this large mass of air could have a variety of dynamic patterns that could presumably affect the layers above.

Verifications of newly developed or adapted codes for highly nonlinear atmospheric dynamics problems is not a trivial task. Rigorous testing with analytical solutions is not possible and equatable comparison with other published solutions is difficult because they usually show chaotic time-dependence and because specific assumptions (e.g., on boundary conditions or neglected terms) are not implemented in different codes or often not even adequately detailed. This generates unnecessary confusion.

A community effort for concerted validation and comparison of the numerous codes currently in use would be extremely valuable at this time, as shown by Polichtchouk et al. [2014] and Cho et al. [2015]. These studies clearly show that one can arrive at erroneous conclusions if simulations from different codes are not at least qualitatively reproduced with the same setup. Even then, erroneous conclusions can still be drawn, as codes which perform well ostensibly in one region of the physical parameter-space do not perform well in another (or, more precisely, extended) region. For example, when pushed to a highly geostrophic region, numerical accuracy of the code can become seriously degraded by small-scale, fast oscillations generated in that part of the parameter-space (e.g., Thrastarson and Cho [2011]). In addition, the atmosphere can exhibit multiple equilibrium states or be driven unwittingly to an unphysical state (Cho et al. [2015], Cho and Thrastarson (in prep.)). In general, the critical question of “realistic and accurate” forcing and initialization—and their effects—is still unsettled, and the question deserves much more attention and scrutiny than has generally received thus far.

5.3 THE SUN

We now turn our attention to the Sun. The Sun is a Main-sequence star. Such stars share a number of common dynamical elements, which include turbulent convection, differential rotation, and magnetic activity. They “burn” hydrogen in their interiors and exhibit a wide range of fluid dynamical behaviour, under various arrangements of radiative and convective zones. For example, in the Sun, $p-p$ chain fusion reaction leads to a radiative core (RC) but the much more rapid C-N-O nuclear burning cycle outside the core gives rise
to a convective envelope outer layer, called the convection zone (CZ). The CZ occupies the 29% of the Sun’s radius ($R_\odot = 6.96 \times 10^8$ m), from its surface down to $\sim 2 \times 10^8$ m in depth. On the surface, vigorous convection, differential rotation (manifested by broad zonal flows) and magnetic cycles are readily visible (e.g., Gough and Toomre, 1991; Christensen-Dalsgaard, 1991). 

### 5.3.1 Observations

Unlike for the stars that host the exoplanets, the surface activity on the Sun can be observed in exquisite detail. Remarkably, even its interior flow can be probed. This is because of the Sun’s proximity. NASA’s Solar Dynamics Observatory (SDO) mission has already provided a large amount of data on solar dynamics and magnetic activities, and this data is complemented by the data from the Solar and Heliospheric Observatory (SOHO) mission as well as by ground-based observatories, which include the Global Oscillation Network Group (GONG) and the New Solar Telescope (NST). Helioseismic soundings using resonant p-modes (acoustic waves) have produced the angular velocity profile, $\Omega_\odot = \Omega_\odot(r, \theta)$, where $r$ is the radial distance from the center and $\theta$ is the colatitude, for most of the Sun’s interior (e.g., Gough and Toomre, 1991; Christensen-Dalsgaard, 2002). Note that $\Omega_\odot$ is not a constant here, in contrast to $\Omega_p$. The $\Omega_\odot$ profile is shown in Figure 5.3.

The Sun’s visible surface has long been known to rotate differentially, with a rotation period of $\sim 25$ days near the equator and $\sim 33$ days at high-latitude regions. This leads to a roughly “three zonal jet” structure, pole-to-pole. These jets are referenced against the quasi-uniform rotation of the RC. The $\Omega_\odot$ profile obtained from the soundings clearly shows the deep structure and the remarkable feature that the rotation contours are roughly constant on conic surfaces (rather than on cylindrical surfaces). Note also that the CZ is bounded by a shear layer near the surface and another, very strong shear layer—called the tachocline (Spiegel and Zahn, 1992)—that span across the bottom of the CZ and the top of the RC.

Doppler measurements of the photosphere and local helioseismic inversions show a poleward meridional flow in the surface layers (Hathaway, 1996; Haber et al., 2002; Zhao and Kosovichev, 2004). Such an azimuthal flow, which is nevertheless closely tied with the mean zonal flow, may be driven by the so-called “gyroscopic pumping” mechanism (McIntyre, 1998). In this mechanism, a negative radial shear (i.e., outward reduction) of $\Omega_\odot$ near the surface suggests a local retrograde (westward) mean flow, which is driven by the divergence of Reynolds stresses associated with supergranulation. In this way, poleward circulation may be considered as a “boundary effect” that occurs near the surface. Although helioseismic inversions so far suggest that the shear layer penetrates too deeply into the CZ for this mechanism to operate effectively (Giles et al., 1997; Braun and Fan, 1998), the inversions are currently not sensitive enough to definitively rule it out.

At the lower boundary of the CZ, if the tachocline is in thermal wind balance (see, e.g., Tassoul, 2000), then the strong radial rotational shear from helioseismic inversions entail the presence of latitudinal thermal gradients. This is expressed by the relation,

$$\Omega_\odot \cdot \nabla v_\varphi = \frac{g}{2c_p r} \frac{\partial \pi}{\partial \theta},$$

(5.1)

where $\Omega_\odot = \Omega_\odot \hat{e}_z$ is the rotation vector, $v_\varphi$ is the longitudinal (eastward) velocity component and the overbar denotes an average over the longitude $\varphi$ and time $t$; $g$ is the gravitational acceleration, $c_p$ is the specific heat per unit mass at constant pressure, and $s$ is the specific entropy per unit mass (e.g., Brun and Toomre, 2002; Durney, 1999; Elliott et al., 2000; Kitchatinov and Rüdiger, 1995; Miesch, 2005). Note that Equation 5.1 is valid for an ideal gas in hydrostatic balance with nearly adiabatic background and with $R_o \ll 1$.

The existence of the tachocline, and its associated overshoot region, may have a profound effect on the dynamo action and mean flow in the Sun—for example, in promoting the generation of strong toroidal fields, cyclic variability, and zonal jets. Turbulent motions in the convection zone expel magnetic fields toward the boundaries, and the asymmetry between upward and downward flows gives rise to a systematic pumping of mean and fluctuating fields downward (Tobias et al., 2001; Ziegler and Rüdiger, 2003). The
stable stratification of the overshoot region helps to suppress magnetic buoyancy instabilities; and, the rotational shear of the tachocline amplifies mean toroidal fields, which could enhance the dissipation of small-scale fields through the distortion and fragmentation of closed magnetic loops. The latter is related to the rotational smoothing process, in which opposite field polarities are brought together by vertical shear and are dissipated on a timescale of $\sim [\lambda^2/(\eta \delta^2)]^{1/3}$ (Spruit 1999); here $\lambda$ is the longitudinal wavelength, $\eta$ is the magnetic diffusivity, and $\delta = dU/dr$ is the radial shear with $U$ the characteristic speed.

Although high-precision oscillation measurements have greatly advanced our understanding of the dynamical structure of the Sun’s surface and interior, several important issues remain to be resolved. For example, there is a discrepancy between the spectroscopically-determined abundance of heavy elements on the surface and the abundance deduced from global helioseismology data (in conjunction with numerical modeling), indicating that understanding of the basic physics of the interior is still incomplete. This includes fundamental issues related to the properties of non-ideal plasma, which affect the accuracy of abundance estimates in helioseismology. In addition, it is still not known whether the inner, energy-generating part of the RC rotates faster or slower than the outer part. This is of fundamental importance for understanding the overall formation and evolution of the Sun—including the jets.

5.3.2 Jets in Simulations

Numerical simulations of 3D convection in a rotating spherical annulus have become quite advanced. In these simulations, the large-scale flow is in a state close to a “geostrophic balance” (see, e.g., Tassoul 2000, Christensen 2002), if $Ro \ll 1$. In this situation, Equation 5.1 reduces to the Taylor-Proudman theorem, in which zonal velocity (specifically, $\Omega_\theta$) contours are cylindrical and aligned with the rotation direction $\Omega_\phi$. In a less rapidly-rotating annulus (which is more appropriate for the Sun), a non-cylindrical rotation component is exhibited and maintained by the baroclinic forcing term on the right-hand side of Equation 5.1. The baroclinicity is such that the rotation profile is more conical, with $\Omega_\phi$ decreasing toward the poles. This would entail a poleward entropy gradient (that is, $\partial S/\partial \theta < 0$ in the northern hemisphere). However, it is important to note that Equation 5.1 is not valid for a slowly rotating (i.e., $R_o \gtrsim 1$) object.

On the other hand, although the Sun is often thought of as a slowly rotating star, the rotation is rapid enough that the meridional acceleration of zonal shear flows by the Coriolis force can exceed the convective momentum transport by the Reynolds stresses in the meridional plane. This is because locally $R_o$ can be much less than unity—for example, if the speed and length scales based on the Reynolds stress (and $\Omega_\phi$ of the RC) are used. The smallness of $R_o$ for such structures appear to be well-supported by 3D global convection simulations (e.g., Brun and Toomre 2002, Elliott et al. 2000). This may also be of importance for gaseous exoplanets. Note that, for the Earth, $R_o$ for convective scales tends to be large, in contrast to what is seen here.

Global solar convection simulations that incorporate a tachocline of rotational shear do indeed exhibit pumping, organization, and amplification of toroidal fields (Browning et al. 2006). This can be seen in Figure 5.5. In the figure, the CZ is dominated by fluctuating, non-axisymmetric fields that account for approximately 95% of the total magnetic energy (Figure 5.5a). In contrast, the field in the tachocline is dominated by oppositely directed, axisymmetric toroidal bands in the northern and southern hemispheres (Figure 5.5b) that account for more than 60% of the magnetic energy and that which have persisted for a simulated duration of 17 (Earth) years, apart from a brief 105-day interval of symmetric parity (Browning et al. 2007). The mean poloidal field is similarly more ordered, with a stronger dipole moment and less frequent reversals.

In general, rotational shear can be established through the advection of angular momentum by meridional circulation, as well as through convective Reynolds stresses. In a statistically steady state, these must be balanced so that

$$\nabla \cdot \left( \rho_0 \left( \mathbf{v}_m \left( L - \mathbf{v}_p \mathbf{v}_m \right) r \sin \theta \right) \right) = 0, \quad (5.2)$$

where $\mathbf{v}_m$ is the meridional velocity component, $\rho_0 = \rho_0(r)$ is the density averaged over horizontal surfaces and time, $L = r \sin \theta (\Omega r \sin \theta + \mathbf{v}_p)$ is the specific angular momentum (Tassoul 2000), and primes indicate small perturbations about the mean. More broadly, Equation 5.2 generalizes to

$$\rho_0 \mathbf{v}_m \cdot \nabla L = -\nabla \cdot \mathbf{F}, \quad (5.3)$$

where $\mathbf{F}$ is an angular momentum flux that includes contributions from Reynolds and Maxwell stresses as well as viscous diffusion. Note that the latter is negligible in stars, but not in numerical simulations.

According to Equation 5.3, an acceleration of the jet $\mathbf{v}_p$, powered by the divergence of $\mathbf{F}$, induces a meridional flow across constant surfaces of $L$. This is effected through the Coriolis acceleration of $\mathbf{v}_p$ and may also involve a contribution from the thermal wind, which satisfies Equation 5.1. Because $\nabla L$ is directed cylindrically outward, away from the rotation axis, retrograde and prograde forcing induce circulations toward and away from the rotation axis, respectively: this is the gyroscopic pumping mechanism discussed above. But, note that cylindrically inward gradients satisfy the necessary condition for instability according to the Rayleigh criterion (see, e.g., Tassoul 2000 and references therein). In the gyroscopic pumping, circulations are set up in response to the convective Reynolds stress and the rotational shear, which satisfy the dynamical balance of Equations 5.1 and 5.3—a delicate balance between large forces. Hence, the meridional circulation in numerical simulations typically exhibits very large fluctuations (\sim 300%) about its temporal mean (Miesch and Toomre 2009). This is probably not realistic.

Baroclinicity, which could aid in establishing a non-cylindrical rotation component, does arise in global convection simulations of the Sun. As noted, this is due to the in-
fluence of rotation on the convective heat flux, which tends to establish a poleward entropy gradient and non-cylindrical $\Omega_{\odot}$ profiles [Brun and Toomre, 2002; Elliott et al., 2000; Miesch et al., 2000]. However, Miesch et al. (2006) suggest that thermal coupling to the tachocline may also play a significant role in establishing a non-cylindrical profile. Here the sub-adiabatic stratification of the lower tachocline is essential for ensuring that a Sun-like rotational shear can generate a poleward entropy gradient via meridional circulation. Ultimately, the profile obtained depends on a complex interaction between convection, differential rotation, meridional circulation, and thermal stratification. This is illustrated in Figure 5.6.

Reynolds stresses maintain a significant differential rotation but angular velocity contours are cylindrical in accordance with the Taylor-Proudman theorem, when a uniform entropy is imposed at the lower boundary (Figure 5.6a). On the other hand, a conical rotation profile is obtained (Figure 5.6b), if a poleward latitudinal entropy gradient imposed at the lower boundary (Figure 5.6c). Note that the mean rotation profile and specific entropy variation shown in Figure 5.6c,e satisfy Equation 5.4 in the lower convection zone. However, thermal wind balance breaks down in the upper CZ, where $R_e$ can be of order unity locally. Similar results are obtained by interior convection simulations (e.g., Elliott et al., 2000; Robinson and Chan, 2001; Brun and Toomre, 2002; Brun et al., 2004). However, the relative amplitude of the thermal perturbation associated with the non-cylindrical wind component is $\sim 10^{-6}$ to $\sim 10^{-5}$, corresponding to a very small temperature variation of $\sim 10^K$ (Miesch et al., 2006). Such variations are too small to be detected by helioseismology (Gough et al., 1996). Hence, at present it is unknown whether Equation 5.1 is satisfied in the deep solar CZ.

As was the case for exoplanets, numerical resolution also plays a significant role. Low resolution simulations lead to a multi-celled meridional circulation profiles (not shown); but, with higher resolution, simulations exhibit a meridional circulation dominated by a single cell in each hemisphere (Figure 5.6d). In the latter, there is a pole-ward flow in the upper convection zone and an equator-ward flow in the lower CZ. According to Miesch et al. (2008), the thin counter-cells near the lower boundary are likely the effect of the thermal and mechanical boundary conditions. Simulations of solar convection that include an underlying stable zone exhibit equator-ward circulation throughout the overshoot region as a consequence of the turbulent alignment of downflow plumes (Miesch et al., 2000).

5.3.3 Some Important Issues

Throughout the preceding discussion of the Sun, magnetic fields have essentially been left out. This is clearly not justified in the interior. Significantly, global MHD convection simulations indicate that dynamo-generated magnetic fields tend to suppress the rotational shear established by Reynolds stresses and baroclinicity, which is in contrast to the helioseismology data.

In some dynamo theories, toroidal magnetic fields are generated and stored in the tachocline: in these theories, the meridional flow transports the toroidal field in the tachocline towards the Equator (which could explain the “butterfly diagram”) and the emergence of loops on the surface that form the sunspot regions. However, the strength of the return meridional flow is largely unknown and the required eddy diffusivity is high—roughly 10 times that predicted by standard mixing-length theory (MLT). Of course, the accuracy of MLT is also uncertain. Another idea for the interior is that proposed by Balbus (2009), who argues that weak magnetization renders CZ more prone to baroclinic instabilities than without magnetization. Although the problem, in principle, requires the knowledge of the functional relationship between entropy and rotation, even simple models readily produce results in broad agreement with the helioseismology data. The theory, however, does not apply to the tachocline, where a simple thermal wind balance may not be valid.
A possible alternative to the above ideas is a model that incorporates the presence of the near-surface rotational-shear layer, in which the magnetic field is generated in the bulk of the convection zone but the butterfly pattern is generated in the shear layer. Interestingly, local-helioseismology does provide evidence that the meridional circulation may consist of two-tiered radial cells. Synoptic analysis of magnetic patterns on the solar surface, such as rotation of sunspot groups and their inclination relative to the Equator, known as Joy’s law (see, e.g., Thomas and Weiss 2008) may provide additional important information. Unfortunately, at present the interpretation of local-helioseismology inversion is a subject of some debate—although steadily more realistic numerical simulations is alleviating the debate.

It has been argued above that the solar tachocline plays an essential role in the solar dynamo, and ultimately the jets, as the likely region in which mean toroidal flux is generated and stored, eventually emerging from the solar photosphere as bipolar active regions. Transport of helical magnetic flux into the tachocline and the generation of nonhelical fields via rotational shear may also promote large-scale field generation in the convection zone by circumventing dynamical quenching constraints. Furthermore, simulations of convection indicate that the presence of a tachocline can help organize and amplify mean fields and can modify the differential rotation profile throughout the solar envelope via baroclinic forcing. Poleward angular momentum transport in the tachocline owing to instabilities or penetrative convection may also influence the global rotation profile, possibly offsetting equatorward transport by giant cells in the convection zone (Gilman et al. 1989).

In summary, the following crucial questions come to the fore: “How do the upper and lower boundary layers influence the internal dynamics of the CZ?” and “How is the tachocline confined?” Pertaining to the first question, issues include baroclinic forcing, gyroscopic pumping, magnetic helicity flux, tachocline instabilities, and inertial, gravity, and Alfven waves, as well as the ways in which each affects the thermal, mechanical, and magnetic coupling between the RC and CZ. Concerning the second question, the relative roles of fossil magnetic fields, dynamo-generated fields, tachocline instabilities, and internal gravity waves need to be better studied. The tachocline is thought to be maintained against the downward spreading of differential rotation induced by gyroscopic pumping, radiative diffusion, and baroclinic circulations. Tachocline instabilities (Gilman and Fox 1997), fossil magnetic fields (Gough and McIntyre 1998), and internal gravity waves (Talon et al. 2002) may all act to halt this spreading and thereby to maintain uniform rotation in the radiative interior. Understanding this most fundamental question is an essential prerequisite to understanding the tachocline’s broader role in the dynamics of the solar interior. Moreover, both have far reaching consequences for understanding other type of main-sequence stars as well.

**5.4 DISCUSSION**

We conclude this Chapter by summarizing and commenting on some key observations and simulation results of jets on exoplanets and on and in the Sun. Experience from studies of the Earth and other Solar System planets shows that a hierarchy of theoretical models is necessary to build a robust understanding of both exoplanets and the Sun, and ultimately to help interpret observations. Solar System planets can—and should—be used as validation, as well as a general guide. But, we must bear in mind still the many things that
are poorly understood even for the Solar System planets. For example, what is the deep structure of the Jovian jets? Often, the way the atmosphere responds to the various forcing and damping is a subtle affair, requiring detailed knowledge of the species composition and distribution; knowledge of the physical properties of fluids in high pressure environment is critical but still the stuff of forefront research. Uranus and Venus starkly show us that the atmosphere can respond in a manner that is quite unexpected: Uranus has zonal jets, even though it is heated at the pole, while Venus’ cloud top is dominated by an equatorial jet and a stable polar vortex even though the planet rotates extremely slowly.

On synchronized planets, waves and eddies (not just the mean flow) would likely be involved in transporting heat between the dayside and the nightside, as well as between the tropics and the poles. There is likely to be upwelling motion at the substellar point and downwelling motion away from this point and inside polar vortices—if they exist (e.g., Cho et al., 2003). Wave momentum flux and adjustments could reduce shears and temperature gradients, and these mechanisms remain to be modeled accurately. As noted, current simulations are in agreement about the small number of broad jets to be found on synchronized planets, but they do not resolve small-scale processes nor incorporate realistic forcing and boundary conditions. Moreover, ideal and non-ideal MHD simulation studies are still at the beginning stages.

For the Sun, here we have reviewed many of the subtleties that must be confronted. In the past few decades, helioseismology has provided evidence that CZ is bounded above and below by strong shear layers, and the presence of such layers can have a critical role in influencing the overall dynamics. Our understanding of the solar interior and surface has advanced substantially, due to helioseismology and continuous observations of solar oscillations by the Helioseismic and Magnetic Imager (HMI) onboard SDO and by GONG. High-precision measurements of oscillation frequencies have provided the radial sound speed profile and the distribution of the angular velocity through the whole interior, except perhaps the very inner core of the Sun. Observations of local processes steadily becoming available, along with advances in computing, will add to the remarkable advances that have already been made in this research area.

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