Deep convection–driven vortex formation on Jupiter and Saturn

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The surfaces of Jupiter and Saturn have magnificent vortical storms that help shape the dynamic nature of their atmospheres. Land- and space-based observational campaigns have established several properties of these vortices, with some being similar between the two planets, while others are different. Shallow-water hydrodynamics, where the vortices are treated as shallow weather-layer phenomenon, is commonly evoked for explaining their formation and properties. Here, we report novel formation mechanisms for vortices where the primary driving mechanism is the deep planetary convection occurring in these planets. Using three-dimensional simulations of turbulent convection in rotating spherical shells, we propose two ideas: (i) Rotating turbulent convection generates deep axially aligned cyclones and anticyclones; (ii) a deep planetary dynamo acts to promote additional anticyclones, some as large as Jupiter’s Great Red Spot, in an overlying atmospheric layer. We use these ideas to interpret several observational properties of vortices on Jupiter and Saturn.

INTRODUCTION

By modeling the dominant dynamical features present on the surface of Jupiter and Saturn, namely, zonal jet streams and storms/vortices, we can learn about what drives them and their connection to the planetary deep interior. Apart from the well-known and persistent hexagonal storm on Saturn and the Great Red Spot (GRS) on Jupiter, there are numerous other compact vortical storms with various sizes and lifetimes found at various latitudes on these planets. By analyzing the images taken by the Cassini spacecraft through its lifetime, Trammel et al. (1, 2) provide a detailed outlook on the properties of compact vortices on Jupiter and Saturn. In our study, we focus only on the fluid dynamical properties of the vortices where the following general features stand out. First, with the exception of the equatorial regions containing a broad eastward zonal jet, vortices are generally found at all latitudes but tend to occur preferentially in regions of westward zonal flow (1, 3). There is a stark contrast in the number of vortices between the two planets: about 200 with 1000 km or larger diameter on Jupiter, while only 10 to 50 on Saturn (1). Note that while both planets also have smaller vortices, the disparity in the numbers still persists (3, 4). Vortices are also present at or very close to the rotational poles: Saturn has a cyclonic polar vortex at each pole (5), while Jupiter has a cluster of cyclones at each pole (6).

The rotation sense of vortices is also an important property but much harder to infer through direct observations. Li et al. (4) analyzed 500 vortices with a diameter of at least 700 km on Jupiter using Cassini images. They inferred the vorticity of the largest 100 vortices, all of them being anticyclonic. An earlier study (7) using Voyager images reached a similar conclusion. Because of the more hazy atmosphere of Saturn, the contrast of the different fluid dynamic features is much lower compared with Jupiter. Paired with the lack of continuous images of Saturn’s surface, determining the vorticity is not possible for most of the storms. Instead, the vorticity of a compact vortex is assumed to be determined by the vorticity of the local zonal wind shear (1–3). On the basis of this assumption, Saturn’s atmosphere does not appear to have a strong bias toward anticyclones. We must, however, keep in mind that the number of vortices on Saturn is small and under-

goess large variations over time. For example, during the 7-year period from 2008 to 2015, the number of vortices of at least 1000-km diameter changed from about 5 to 20 in the northern hemisphere of Saturn (2). Therefore, a robust statistical trend cannot be inferred for Saturnian vortices.

Much like in the study of zonal jets on gas giant planets [see (8–10) for shallow jets and (11–15) for deep jets], vortices can also be understood as either existing only in the outermost weather layers or extending much deeper into the planetary interior. In the case of zonal jets, exciting progress was recently made by modeling the gravity data from the Juno mission and the Cassini Grand Finale. It indicates that the zonal jet streams on both planets are likely thousands of kilometers deep, about 9000 km in Saturn (16) and about 3000 km in Jupiter (17). These results lend support to the deep zonal jet scenario where planetary convection is the driving force. Given these results, it is certainly worth investigating whether some of the vortices visible on the surface extend deep into the interior.

So far, most vortex formation models consider only the thin outer weather layer, with several using the shallow-water equations with a prescribed forcing in the form of moist convection or solar heating (18, 19). These studies have successfully modeled several properties of the vortices, for example, the bias toward anticyclones (20, 21) and polar vortex formation (22, 23). However, it is also well known that coherent vortices can spontaneously form in three-dimensional (3D) fluid turbulence if rotational effects are substantial [e.g., see (24)]. Could deep planetary convection also generate coherent vortices along with zonal jets? There have been several attempts to study this. Coherent, primarily anticyclonic, vortices have previously been shown to coexist with zonal jets using an elastic, spherical deep convection model with a stably stratified layer (25). Others have investigated vortex formation in Cartesian geometry (26–28), where the curvature effects due to the spherical shell are not present. Here, we report several physically motivated simulations of planetary convection in deep spherical shells to better understand how vortices form and behave under such conditions.

RESULTS

Thin shell case

Keeping the general features of the giant planet interiors in mind (Fig. 1), we explore two cases that will help us to better understand
vortex formation in rotating spherical shells. We first consider the dynamics in a thin rotating spherical shell—a generic representation of the outer convective layers in giant planets that couple only weakly with the interior magnetic field (14, 29)—which spans a region from 0.97$r_o$ to $r_o$, where $r_o$ is the outermost radius of the shell. Such a shell thickness would be equivalent to about 2000 and 1000 km if scaled to Jupiter and Saturn, respectively. Since contemporary 3D simulations cannot attain the highly turbulent geostrophic regime present in the interiors of giant planets, we need to set up the simulation such that it can promote relevant physics. The choice of assuming a relatively thin layer is motivated by the fact that 2D and quasi-2D rotating flows are well known for vortex formation in a broad parameter regime (30). The thin shell nature of the simulation will allow it to excite vortices but still retain the 3D nature of the convective flows. We ignore the fluid compressibility in this simulation for the sake of simplicity and to reduce the computational costs associated with the model. Furthermore, a fourfold symmetry is imposed in the azimuthal direction to limit the computational requirements. The crucial control parameters governing the used nondimensional Boussinesq equations are as follows: the Ekman number defining the ratio of viscous and Coriolis forces is $10^{-5}$, the Rayleigh number defining the convective driving is $3 \times 10^5$, and the thermal Prandtl number is 0.1. Additional model details can be found in Materials and Methods.

This thin layer model self-consistently generates about seven alternating zonal jets in each hemisphere with an eastward jet in the equatorial region (fig. S1), as well as a number of compact vortices. Alternating zonal jet streams have already been reported in earlier simulations of spherical shells representing the deep convection in the outer atmospheres of giant planets [e.g., see (12–15)]. Furthermore, anticyclonic vortices, along with zonal flows, were also reported earlier (25) in a similar but thicker shell setup with an overlying stably stratified layer.

The simulation generates about 40 well-defined, plus additional smaller and shorter-lived, vortices in each hemisphere (Fig. 2). The vortices are formed in the high to mid-latitude regions (fig. S2), and

The partial circles with arrows represent overturning convection. The overturns with thick lines highlight energetic convection plumes that impinge on the adja-
cent layer. Jupiter may be crudely approximated by a two-layer structure with the deep interior consisting of metallic hydrogen with larger conductivity, and a lower
conductivity molecular hydrogen outer layer (38). Saturn, on the other hand, likely has a three-layer structure: an innermost metallic hydrogen layer sustaining the dy-
namo, a molecular hydrogen-rich low-conductivity outermost layer, and a stably
stratified layer due to helium rainout (57) between the other two layers. Although
the exact location and nature of the intermediate helium-enriched layer are uncer-
tain, its existence is highly favored (57, 58) due to the very high axisymmetry of
Saturn’s magnetic field (59). Jupiter may also have a similar stable layer, but it is
likely much thinner (60); furthermore, Jupiter’s magnetic field is not highly axisym-
metric (61, 62).

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properties of both 2D and 3D setups, is a good compromise to simulate relevant physics.

**Thick shell case**
Next, we investigate what happens when a hydrodynamic layer directly interacts with another layer where the electrical conductivity is large enough to allow generation of magnetic fields (Fig. 1). We consider a rotating spherical shell that spans a region from 0.1\(r_o\) to \(r_o\), within which the fluid density decreases by a factor of about 150 along the radius in the shell. Similar to Dietrich and Jones (33), the electrical conductivity of the fluid changes with radius as a hyperbolic tangent function: It is constant until \(\approx 0.84r_o\), and then sharply decreases by seven orders of magnitude within \(\approx 0.07r_o\), and stays constant afterward (see "Profile 1" in fig. S4). Such an electrical conductivity profile defines a deep dynamo layer and an overlaying atmospheric layer where Lorentz forces will be negligible. Therefore, we label the radial level at 0.85\(r_o\) as the “transition” radius separating the dynamo and the atmospheric region. Note that our chosen electrical conductivity profile differs from several earlier studies: The hyperbolic tangent function introduces sharper conductivity gradients in the transition radius as compared with a combination of constant for some radius and exponentially decreasing afterward (34–37). The electrical conductivity in Jupiter starts to decay with a superexponential rate after about 0.9\(r_o\) (38), which leads to decoupling of the atmosphere from the dynamo in a relatively small radius range. To capture such a decoupling, it is crucial to include a strong electrical conductivity decay offered by a hyperbolic tangent function with a large step change. We then solve the convective magnetohydrodynamics in this setup. The crucial parameters governing the anelastic system of equations are as follows: the Ekman number is \(10^{-6}\), the Rayleigh number is \(4 \times 10^9\), the thermal Prandtl number is 0.1, and the magnetic Prandtl number is 1.5. We refer the reader to Materials and Methods for more details about the model.

In this two-layer setup, the dynamics and the emergent features are very different from the earlier thin shell case. The simulation generates a largely dipole dominant magnetic field in the interior
The deeper dynamo layer enforces a nearly rigid body rotation in the fluid below \( \approx 0.85 r_e \). Above this radial level, a hypothetical magnetic Lorentz force is likely playing a role in promoting the formation of anticyclones. This hints at magnetic Lorentz forces being a key ingredient for triggering them. Another simulation with a more conventional profile of constant conductivity until \( 0.85 r_e \) and an exponential drop afterward (“Profile 2” in fig. S4) did not produce similar anticyclones. This is likely due to the presence of a gradual transition of the fluid dynamics from the dynamo layer to the hydrodynamic layer, which makes it one system rather than two layers interacting with each other. Since the anticyclones are formed in the outer low-conductivity layer, they do not have a corresponding magnetic signal associated with them (fig. S7). However, the simulation does have a tendency to generate more anticyclones in the regions containing higher magnetic energy. This hints that magnetic Lorentz forces are likely playing a role in promoting the formation of the anticyclones.
DISCUSSION
We have presented two simulations of highly turbulent convection in rotating spherical shells. The first simulation demonstrates the spontaneous formation of zonal jets, cyclones, and anticyclones in a thin spherical shell. The second one shows a novel vortex formation mechanism where plumes originating in a deep dynamo layer exclusively excite anticyclones in an overlying low-conductivity atmospheric layer.

The results from these simulations can help us better understand the atmospheric dynamics visible on the surface of Jupiter and Saturn. Although vortex formation (through an inverse cascade of energy) is well known in rotating Cartesian boxes with either forced or convective turbulence (26–28), our simulation is the first—barring Heimpel et al. (25), who report anticyclones in a stably stratified layer above a convective layer—to produce both cyclones and anticyclones, as well as zonal jet streams, in a global spherical geometry with convective driving. Compared with local Cartesian box simulations, our thin shell simulation can be considered a step further toward the goal of modeling the global atmospheric dynamics of gas giant planets. The global geometry allows us to simultaneously capture the vortex-free near-equatorial regions, alternating zonal jets, formation of storms with vorticity governed by local shear flows, and nonlinear interactions between storms (including mergers). The simulation shows that rotating convection can generate deep axially aligned storms that mimic many properties of the storms on Jupiter and Saturn.

The second case shows how a dynamo layer might interact with an overlying atmospheric layer. The dynamo region is conventionally viewed as an inert layer proving a drag force for the atmospheric zonal jets (12–15). Our simulation, instead, shows that it could play an important role dynamically and may provide a source for seeding and sustaining large anticyclones in the atmospheric layer. Given that, there is a possible connection between our results and Jupiter’s GRS. Jupiter’s deep dynamo is expected to be in the magnetostrophic regime where small-scale and planetary-scale giant convective cells are expected (43–46). If Jupiter’s dynamo layer does seed anticyclones in the atmospheric layer, as suggested by our simulation, then a GRS-type anticyclone with a quiet center can be readily formed. Our simulation forms several GRS-scale anticyclones. Furthermore, long-lived vortices in this simulation also showed a GRS-like westward drift: about 0.5°/day in the simulation, while about 0.3°/day for GRS in recent times (47). In the simulation, the deep dynamo sets up a slow westward drift in the low to mid-latitudes. The convection plumes in the dynamo region accordingly drift westward, dragging the correspondingly large vortex in the atmospheric layer. The stability of the GRS, however, could not be captured in the simulation. The most stable storms in our simulation existed for about 20 days, while the GRS has a lifetime of hundreds of years (47). Possibly, Jupiter is able to sustain much longer-lived convective plumes, spawning stable GRS-like vortices, than our simulation. A broader systematic parameter study of this simulation will be useful to investigate the possibility of longer-lived large anticyclones.

On the basis of these two case studies, we can begin piecing together a global picture of the deep convection–driven fluid dynamics occurring in Jupiter and Saturn. An atmospheric layer with dynamics governed by rapidly rotating deep convection will tend to promote either cyclonic or anticyclonic vortices/storms depending on the direction of the vorticity set by the local zonal wind shear. This phenomenon is likely active in the atmospheric layers of both Saturn and Jupiter. On top of that, if the interior structure of a planet allows the dynamo layer to directly interact with the outer hydrodynamic layer, then energetic plumes from the deep dynamo layer will exclusively excite anticyclonic vortices in the atmospheric layer, tilting the balance toward more anticyclones than cyclones. This scenario is likely active in Jupiter, explaining why it has many more anticyclones than cyclones (4, 19). On Saturn, however, the dynamo layer is largely decoupled from the outer hydrodynamic layer due to the likely presence of a thick stably stratified layer. Even if its dynamo layer generates energetic plumes, they get trapped by the overlying stably stratified layer, explaining the absence of a clear and consistent preference for anticyclones on Saturn (2, 3, 19).

The vortex formation mechanisms discussed above, however, do not exclude the presence of other mechanisms that have been proposed. For example, moist convection–driven shallow storms might be forming in the outermost layers of both Saturn and Jupiter. Furthermore, the idea of anticyclone formation through convective pumping of a thin stably stratified layer from below (25) might also be active on both planets. The atmospheres of Jupiter and Saturn exhibit extremely rich dynamics, suggesting the presence and complex interaction of various storm formation mechanisms. Here, we offer two possibilities that appear promising. The Juno spacecraft currently in orbit around Jupiter will likely shed more light on this issue. The on-board microwave radiometer (48) observes thermal radiation from levels as deep as 1000 bars. If large atmospheric vortices on Jupiter maintain some correlation with the temperature at depth, then we might be able to ascertain if some of the storms on the surface go deeper into the atmosphere.

MATERIALS AND METHODS
Thin shell case
The first simulation discussed here consists of a thin spherical shell bounded by a lower boundary at 0.977 r_o and an outer boundary at r_o. The spherical shell rotates about a fixed axis, whose direction vector is represented by \( \hat{\varepsilon} \), with an angular velocity \( \Omega \). The fluid dynamics in this shell is governed by thermal convection set up by a temperature gradient of \( \Delta T \) between the inner and the outer boundary. The fluid is assumed to be incompressible, i.e., \( \nabla \cdot \vec{u} = 0 \), where \( \vec{u} \) is the velocity. We use the classical Boussinesq approximation to model the fluid dynamics. Furthermore, we work with nondimensional system of equations where the distances are scaled by the shell thickness \( d = r_o - r_s \), time is scaled by \( \Omega^{-1} \), and temperature is scaled by \( \Delta T \). With these assumptions, the equations governing the velocity and temperature \( T \) are

\[
\frac{\partial \vec{u}}{\partial t} + \vec{u} \cdot \nabla \vec{u} + 2\Omega \times \vec{u} + \nabla P = \frac{Ra}{Pr} g(r) T \vec{\hat{r}} + E \nabla^2 \vec{u} \tag{1}
\]

\[
\frac{\partial T}{\partial t} + \vec{u} \cdot \nabla T = \frac{E}{Pr} \nabla^2 T \tag{2}
\]

where \( P \) is pressure. The nondimensional gravity \( g(r) \) varies as \( (r_o/r)^2 \), assuming that most of the planetary mass is concentrated below the inner boundary at 0.977 r_o. The relevant nondimensional parameters are as follows: Rayleigh number \( Ra = ag_o \Delta T d^3/\nu k \), Ekman number \( E = \nu/(\Omega D^2) \), and Prandtl number \( Pr = \nu/\kappa \), where \( a \) is the thermal expansion coefficient, \( g_o \) is gravity at the outer boundary, \( \nu \) is viscosity, and \( \kappa \) is thermal diffusivity. These control parameters are set as \( E = 10^{-3} \), \( Ra = 3 \times 10^{14} \), and \( Pr = 0.1 \). The velocity matches a stress-free condition on both boundaries. The temperature is assumed to be constant on each boundary.
The above system of equations is solved using the open source MagIC code [https://magic-sph.github.io; (49)], which uses the pseudospectral approach. The latitudinal and longitudinal quantities are expanded using the Legendre polynomials, and the radial ones are expanded using the Chebyshev polynomials. The code also uses the toroidal-poloidal decomposition to maintain strict divergenceless nature of the relevant quantities. The code uses the fast spherical harmonic library SHTns [https://nscheaff.bitbucket.io/shhtns/; (50)].

The equations are time advanced using an explicit second-order Adams-Bashforth scheme for the Coriolis and the nonlinear terms, and an implicit Crank-Nicolson scheme for other terms (51).

We simulate the system for about 1500 rotations, which was enough to show a statistically stationary behavior. To limit the computational requirements, we used fourfold symmetry in the azimuthal direction (14, 52), which effectively makes it a quarter-spherical-wedge simulation with periodic boundary conditions on the meridional planes of the wedge. The first 930 rotations were carried out on a simulation grid of [1056, 2112, 160], where the numbers represent resolution in the azimuthal, latitudinal, and radial directions, respectively. The corresponding maximum spherical harmonic degree is 1408. The remaining time period was simulated on a larger grid of [1152, 2304, 200]. Apart from sharpening the fluid dynamic features, the higher resolution did not change the simulation behavior. Despite the large grid we used, a hyperdiffusion scheme, which suppresses energies at small length scales, had to be applied to run the simulation. Following several earlier studies (14, 15, 25), the viscosity in our setup becomes a function of spherical harmonic degree after a certain cutoff. MagIC multiplies the following function to the main viscous diffusion operator

\[
d(\ell) = 1 + D \left[ \frac{\ell}{\ell_{\text{max}} + 1 - \ell_{\text{hd}}} \right]^\beta
\]

where \( \ell \) defines the amplitude of the function, \( \ell_{\text{max}} \) is the maximum spherical harmonic degree used in the simulation, \( \ell_{\text{hd}} \) is the degree after which the hyperdiffusion starts, and \( \beta \) defines the rise of the function for degrees higher than \( \ell_{\text{hd}} \). For our simulation, we use \( D = 10, \beta = 6, \) and \( \ell_{\text{hd}} = 450 \).

### Thick shell case

In this simulation, we use the anelastic system of equations (53, 54) to model compressible, subsonic flows present in the interior of giant planets. In this approximation, the thermodynamic quantities are assumed to be a combination of a static background (tilde) and small fluctuations (prime), i.e., \( x = \bar{x} + x' \). The equations used in our simulation are in the entropy variable form with entropy contrast between the inner and the outer boundaries driving the convection (29). The magnetic field is scaled by \( \sqrt{\rho_o} \mu_o \lambda_i \Omega_i \), where \( \rho_o \) is the density on the outer boundary, \( \mu_o \) is the magnetic permeability, and \( \lambda_i \) is the magnetic diffusivity on the inner boundary. The scales for distance and time are the same as above. With these assumptions and rescaling, the evolution for velocity is governed by

\[
\nabla \cdot (\hat{\rho} \vec{u}') = 0
\]

\[
\left( \frac{\partial \vec{u}'}{\partial t} + \vec{u} \cdot \nabla \vec{u}' \right) = -\nabla \frac{p'}{\hat{\rho}} - \frac{2\varepsilon'}{E} \times \vec{u} + \frac{\hat{R} \rho \hat{\varepsilon} \vec{s}'}{Pr} + \frac{1}{\hat{P} m_i E \hat{\rho}} (\nabla \times \vec{B}) \times \vec{B} + \frac{1}{\hat{\rho}} \nabla \cdot \vec{S}
\]

where \( p' \) is the pressure perturbation, \( \hat{g}(r) \) is the gravity (assumed to vary linearly), \( s' \) is the entropy perturbations, \( \vec{B} \) is the magnetic field, and \( Pr_m = v/\lambda_v \). The Rayleigh number \( Ra \) in this formulation is given by \( \alpha_o g_o \hat{R} \tau_d^2 \Delta c_p \approx 1 \). The traceless rate-of-strain tensor \( \vec{S} \) is defined by

\[
S_{ij} = 2\hat{\rho} \left( e_{ij} - \frac{1}{3} \delta_{ij} \nabla \cdot \vec{u} \right) \text{with } e_{ij} = \frac{1}{2} \left( \frac{\partial u_i}{\partial x_j} + \frac{\partial u_j}{\partial x_i} \right)
\]

where \( \delta_{ij} \) is the identity matrix. The energy conservation gives time evolution equation for entropy perturbation \( s' \)

\[
\rho \overline{\nabla} \left( \frac{\partial s'}{\partial t} + \vec{u} \cdot \nabla s' \right) = \frac{1}{Pr} \nabla \cdot (\rho \overline{\nabla} s') + \frac{Pr Di \lambda_{\text{norm}}}{P m_i E Ra} (\nabla \times \vec{B})^2
\]

where \( \overline{\nabla} \) is the background temperature and the viscous heating contribution is given by

\[
\Phi_v = 2\rho \left[ e_{ij} e_{ij} - \frac{1}{3} (\nabla \cdot \vec{u})^2 \right]
\]

The dissipation number \( Di = \alpha_o g_o d_l/c_p \), where \( \alpha_o, g_o \) are the thermal expansivity and gravity on the outer boundary and \( c_p \) is the specific heat at constant pressure. The evolution of the magnetic field is governed by

\[
\frac{\partial \vec{B}}{\partial t} = \nabla \times (\vec{u} \times \vec{B}) - \frac{1}{\hat{P} m_i} \nabla \times (\lambda_{\text{norm}} \nabla \times \vec{B})
\]

where \( \lambda_{\text{norm}} \) is the local magnetic diffusivity normalized by its value at the inner boundary \( r_s \). The magnetic field also follows the divergenceless condition \( \nabla \cdot \vec{B} = 0 \). We use the anelastic version (55) of the MagIC code to simulate this system of equations.

The inner boundary is at 0.1\( r_o \), and the outer is at \( r_o \). The control parameters for this simulation are as follows: \( E = 10^{-6}, Ra = 4 \times 10^9, Pr = 0.1, \) and \( Pm_i = 1.5 \). The density of the fluid changes by five density scale heights (a factor of about 150) along the radius.

The electrical conductivity follows a step change using a hyperbolic tangent function (33); it remains constant until about 0.85\( r_s \), decreases by factors of 10\(^6\) in a small radius, and stays constant afterward (see Profile 1 in fig. S4). The velocity matches a stress-free condition, the entropy is held constant, and the magnetic field matches a potential field on the shell boundaries. Instead of starting the simulation from small flow and magnetic field perturbations, we used a saturated dynamo simulation (without any conductivity change in its interior) from our earlier study (56) and lowered its Ekman number from 10\(^{-5}\) to 10\(^{-6}\) in incremental stages. Once the 10\(^6\) simulation showed statistical saturation in its kinetic and magnetic energy time series, we then introduced a conductivity jump in stages, reaching a decrease of 10\(^7\) in the outer layer. The simulation ran for about 2000 rotations in the final form. About 50 rotations were simulated on a grid of size [2112, 1054, 400], and the rest were on a grid with [1152, 576, 400]. Here, too, increasing the resolution did not change the simulation nature substantially. This simulation also required hyperdiffusion with parameters \( D = 20, \beta = 3, \) and \( \ell_{\text{hd}} = 450 \) (see earlier section) for the higher-resolution grid and \( D = 5, \beta = 2, \) and \( \ell_{\text{hd}} = 120 \) for the lower-resolution grid.

**SUPPLEMENTARY MATERIALS**

Supplementary material for this article is available at http://advances.sciencemag.org/cgi/content/full/6/46/eabb9298/DC1

REFERENCES AND NOTES


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