responsible for the erosion of the Channeled Scabland (*26*). However, as in the case of the largest Missoula floods (*8*), the peak Chuja flooding seems physically more consistent with a process of cataclysmic dam failure. This dam-break scenario for the peak discharge does not preclude the occurrence of smaller jökulhlaups generated by repeated filling and failure of the ice-dammed lakes. These probably yielded flow peaks on the order of 1×10^6 m^3 s^{-1} (*27*). These relations are remarkably similar to those hypothesized for the late Pleistocene Missoula floods (*8*).

Hydraulic parameters for the Chuja peak flows (Fig. 3) exceed the largest values known (*10*). The superlatives include flow depths of 400 to 500 m, velocities of 20 m s^{-1} (subcritical sections) to 45 m s^{-1} (supercritical sections), bed shear stresses of 5000 N m^{-2} (subcritical) to 20,000 N m^{-2} (supercritical), and stream power per unit area of 10^5 W m^{-2} (subcritical) to 10^6 W m^{-2} (supercritical). These may well have been Earth's greatest floods.

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Why Does the Earth Spin Forward?

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The spins of the terrestrial planets likely arose as the planets formed by the accretion of planetesimals. Depending on the masses of the impactors, the planet's final spin can either be imparted by many small bodies (ordered accretion), in which case the spin is determined by the mean angular momentum of the impactors, or by a few large bodies (stochastic accretion), in which case the spin is a random variable whose distribution is determined by the root-mean-square angular momentum of the impactors. In the case of ordered accretion, the planet's obliquity is expected to be near 0° or 180°, whereas, if accretion is stochastic, there should be a wide range of obliquities. Analytic arguments and extensive orbital integrations are used to calculate the expected distributions of spin rate and obliquity as a function of the planetesimal mass and velocity distributions. The results imply that the spins of the terrestrial planets are determined by stochastic accretion.

The directions of planetary spins are not random. Six of the eight major planets—all except Venus and Uranus—have prograde spin, and those with prograde spin all have obliquity <30°. The predominance of prograde rotation among the planets has often been cited as evidence that they formed by orderly accretion from a thin disk. However, the processes determining planetary spins are different for the terrestrial planets (Mercury, Venus, Earth, and Mars), the giant planets (Jupiter, Saturn, Uranus, and Neptune), and Pluto and Charon. The terrestrial planets consist largely of solids, and the late stages of their formation probably occurred in a gas-free environment by accretion of solid bodies. By contrast, most of the mass and angular momentum of Jupiter and Saturn resides in an envelope of hydrogen and helium, which is believed to have collapsed from the rotating solar nebula onto a core composed of rock and ice. Their spins were probably determined by gravitational torques on or accretion of this envelope (*1*, *2*). Uranus and Neptune are intermediate cases with much smaller gas envelopes than Jupiter and Saturn. The

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obliquities of the four giant planets may have been generated either by large impacts (3–6) or through external torques that rotated the ecliptic plane (7). The formation of Pluto and Charon is poorly understood. In this report we focus on the spins of the terrestrial planets, which form by a process that is simple enough to allow firm theoretical predictions.

The spins of the terrestrial planets were probably determined during their formation, which occurred by the accretion of solid bodies. It is convenient to distinguish two kinds of spin that can arise in this process (3). First, the accreted bodies may have imparted a nonzero mean angular momentum to the forming planet. The angular momentum acquired in this manner is conventionally referred to as the ordered spin component (8). Second, even if the mean angular momentum of the impactors is small, the planet's final spin rate can be substantial if there is imperfect cancellation between the angular momenta of individual impactors, especially if the impactors are massive enough. This process gives rise to the stochastic component of planetary spin (4, 5).

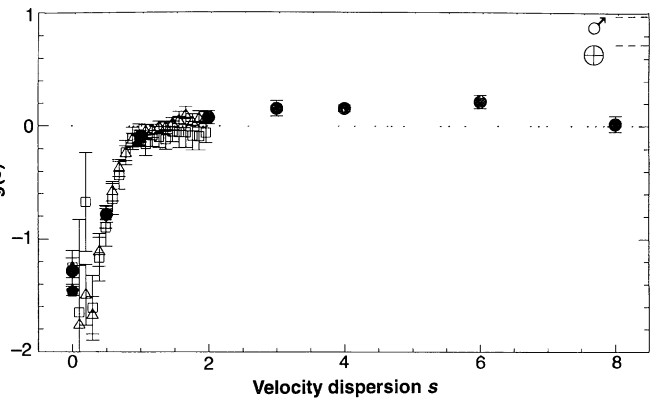

Fig. 1. Values of the function $g(s)$, which measures the dimensionless mean rate of angular momentum accretion (Eq. 5). The horizontal axis is the dimensionless dispersion s, for vertical dispersion $\sigma_z = 0$ (two dimensions) (solid circles), and for $\sigma_z/\sigma_r = 0.5$ as $s \rightarrow 0$ (three dimensions) (solid pentagon). The points are calculated from orbit integrations for planets with dimensionless radius r ranging from 0.001 to 0.01. For comparison, we have also plotted data for $\sigma_z = 0$ from Lissauer and Kary (8, 36) as open triangles ($r = 0.0069$) and open squares ($r = 0.0035$). The function g never attains large positive values, implying that accretion from a uniform disk cannot lead to prograde rotation as rapid as the present rates of Mars (♂) or Earth (⊕), which are marked near the right edge ($g = 0.97$ and 0.72, respectively, from Eqs. 4 and 5). Calculations for $\sigma_z/\sigma_r = 0.5$ and $s > 0$ show no systematic differences from the $\sigma_z = 0$ results, but the statistical errors are larger because of the lower collision rate.

We have examined (9) the specific angular momentum acquired by a planet (assumed to follow a circular, zero-inclination orbit around the sun) that is accreting planetesimals from a differentially rotating disk. Each accreted body imparts a component of specific angular momentum ℓ_z, which is perpendicular to the orbital plane, as well as a component $\mathbf{l}_\perp$ in the orbital plane. Reflection symmetry about the disk midplane guarantees that the expectation value $\langle \mathbf{l}_\perp \rangle = 0$, but $\langle \ell_z \rangle$ can be either positive or negative, implying a (possibly slight) preference for, respectively, prograde or retrograde spin.

The expectation values for the components of the final angular momentum, L_z and $\mathbf{L}_\perp$, of a planet of mass M_p can be written as follows

$$\langle L_z \rangle = M_p \langle \ell_z \rangle, \qquad \langle \mathbf{L}_\perp \rangle = 0 \tag{1a}$$

$$\langle L_z^2 \rangle = \frac{\langle m^2 \rangle M_p}{\langle m \rangle} \langle \ell_z^2 \rangle + M_p^2 \langle \ell_z \rangle^2 \tag{1b}$$

$$\langle L_\perp^2 \rangle = \frac{\langle m^2 \rangle M_p}{\langle m \rangle} \langle \ell_\perp^2 \rangle \tag{1c}$$

where $\langle m \rangle$ and $\langle m^2 \rangle$ are the mean and mean-square masses of the impactors, and we have assumed for simplicity that the impactor mass is much less than the planet mass and that $\langle \ell_z \rangle$, $\langle \ell_z^2 \rangle$, and $\langle \ell_\perp^2 \rangle$ are independent of mass (these assumptions are not correct in general, but the required modifications to these equations are easy to make and do not affect the results below). Stochastic accretion dominates if the first term on the right side of Eq. 1b is larger than the second; in other words, if

$$S_m S_\ell > 1 \tag{2a}$$

with the dimensionless quantities

$$S_m^2 = \frac{\langle m^2 \rangle}{M_p \langle m \rangle}, \qquad S_\ell^2 = \frac{\langle \ell_z^2 \rangle}{\langle \ell_z \rangle^2} \tag{2b}$$

where S_m measures the effective mass of the planetesimals relative to the planet, and S_ℓ measures the degree of cancellation between impactors with positive and negative angular momentum along the z axis. We compute S_m and S_ℓ for cases of interest below. The obliquity ϵ of the planet (the angle between the spin and orbital angular momenta) is given by

$$\cos \epsilon \equiv L_z/(L_z^2 + L_\perp^2)^{1/2} \tag{3}$$

Thus, if $\langle L_\perp^2 \rangle \approx \langle L_z^2 \rangle$, large obliquities are expected.

We now consider the evaluation of S_ℓ. The rate at which a forming planet accretes mass and angular momentum depends primarily on two dimensionless parameters. The first, a dimensionless radius $r \equiv R_p/R_H$, is the ratio of the planet's radius R_p to its Hill radius R_H; we define the Hill, or tidal, radius as $R_H = (GM_p/\Omega^2)^{1/3}$, where Ω is the planet's orbital frequency and G is the gravitational constant (9). During the growth of a planet of constant density, r also remains constant; numerically, $r =$ 0.0022, 0.0030, 0.0042, and 0.0076 for Mars, Earth, Venus, and Mercury, respectively. The second parameter is a dimensionless velocity dispersion $s \equiv \sigma_r/\Omega R_H$, where σ is the radial velocity dispersion of the planetesimal disk. For Earth, $s = 1$ corresponds to a disk of particles with a root-mean-square eccentricity of 0.02 or $\sigma_r = 0.43$ km s^{-1}. In terms of r and s, the usual Safronov number Θ, defined as the square of the ratio of the escape speed from the planet to the velocity dispersion of the planetesimals, is approximately $1/rs^2$.

The value of the dispersion s during the formation of the terrestrial planets is uncertain. In Safronov's models (3, 11), gravitational scattering maintains the disk velocity dispersion near the escape velocity of the forming planet, implying $\Theta \approx 3$ to 5 or, for Earth, $s \approx 10$. More recently, models of "runaway" accretion have been popular. In these models, the largest body in a zone is less efficient in stirring the disk, and $s \approx 0.5$ to 2 at late stages of accretion (12). We will examine a range of assumptions about the value of s below.

A convenient description of the spin rate of a planet is the number of sidereal rotations per revolution around the sun, $\mathfrak{R}$, where positive (negative) values denote prograde (retrograde) spin (8). If r and s remain constant during accretion of a planet, its mean final spin rate will be

$$\mathfrak{R} = \frac{3\langle \ell_z \rangle}{2\Omega R_p^2} \tag{4}$$

We have calculated the distribution of specific angular momentum in accreting bodies for various values of r and s by integrating a large number ($\gtrsim 10^8$) of test particle orbits. We assume that the planet is embedded in a disk of uniform surface density and that any particle that strikes the forming planet is accreted. The conclusions below remain unchanged if the planet has an eccentricity and inclination like those of the terrestrial planets or if the disk surface density varies slowly with distance from the sun. We have performed calculations for both a three-dimensional disk with a ratio of vertical to radial velocity dispersion $\sigma_z/\sigma_r = 0.5$, similar to that found for a wide range of astrophysical disks and implied by theoretical models of protoplanetary disks (13), and for a zero-thickness disk ($\sigma_z = 0$), which has a higher collision rate and, hence, provides better statistics. In most limiting cases we can analytically derive scaling laws for the mean and mean-square angular momentum in terms of r and s; our numerical results are in good agreement with the analytic work.

In the regime of strong gravitational focusing by the planet ($r << 1$, $s << r^{-1/2}$), which is the case most relevant to terrestrial planet formation,

$$\langle \ell_z \rangle = g(s) r \Omega R_H^2 \quad (5)$$

where the function $g(s)$ is shown in Fig. 1. For $s \lesssim 1$, $g(s)$ is negative and decreases as s decreases. For a zero-thickness disk, $g(s)$ approaches a limiting value of -1.28 ± 0.01 as $s \to 0$. For a disk with $\sigma_z/\sigma_r = 0.5$, $g(s)$ is the same as in the zero-thickness case for $s << r^{1/2}$, and $g(s) = -1.46 \pm 0.04$ for $1 >> s >> r^{1/2}$. Thus, in low-dispersion disks, there is a slight excess of impacting bodies with negative angular momentum; so spins are retrograde. For $s >> 1$, scaling arguments imply that $g(s)$ varies approximately as s^{-2}; the proportionality constant is positive and of order unity but is not accurately determined from the data in Fig. 1. We also have, for a three-dimensional disk,

$$\langle \ell_z^2 \rangle = h_z(s) r (\Omega R_H^2)^2 \quad (6a)$$

$$\langle \ell_\perp^2 \rangle = h_\perp(s) r (\Omega R_H^2)^2 \quad (6b)$$

where $r^{1/2} << s << r^{-1/2}$, $h_z(s) = h_\perp(s) = 1/2$ for $r^{1/2} << s << 1$, and $h_z(s) = 1/3$, $h_\perp(s) = 2/3$ for $1 << s << r^{-1/2}$. Note that $\langle \ell_z \rangle^2 << \langle \ell_z^2 \rangle$ because of the near-cancellation of positive and negative angular-momentum contributions.

Figure 1 implies that there is a maximum spin rate that can result from ordered accretion from a uniform disk, whatever the dispersion s may be. The maximum retrograde spin is achieved for $s << 1$, $\langle \ell_z \rangle = -1.46 r \Omega R_H^2$. If the planet accretes at constant r and s, its final spin will then be $\mathfrak{R} = -2.2/r$ (Eq. 4). From Fig. 1, the maximum positive value of $g(s)$ is roughly 0.2, therefore $\mathfrak{R} \lesssim 0.3/r$. Thus, for ordered accretion

$$-2.2 < \mathfrak{R} r < 0.3 \quad (7)$$

corresponding to a range of -750 to 100 days per year for Earth. In terms of the variable S_ℓ (Eq. 2b), this condition can be written approximately as

$$S_\ell^2 \gtrsim \frac{10}{r} \quad \text{for prograde spin} \quad (8a)$$

$$S_\ell^2 \gtrsim \frac{0.2}{r} \quad \text{for retrograde spin} \quad (8b)$$

We now estimate the dimensionless effective mass S_m in Eq. 2b. A simple approach is to assume that a single impactor, of mass m_1, is much more massive than all the rest, in which case

$$S_m = \frac{m_1}{M_p} \quad (9)$$

but the simplicity of this estimate is offset by the artificial assumption that there is a single significant large impactor.

A more realistic assumption is that the number of accreting particles per unit mass follows a power law

$$n(m) \propto m^{-q} \quad \text{for } m > m_{\min} \quad (10)$$

where $m_{\min}$ is much less than a typical planetary mass M_p and $q > 1$. As we show below, we need not assume that the distribution has an upper cutoff. Power laws are produced naturally in coagulation and fragmentation processes, and a variety of arguments—observed size distributions of lunar craters (*14*), asteroids (*15*), and comets (*16*) and theoretical studies of collisionally evolved systems (*17*)—suggest that $1.5 < q < 2$.

A difficulty in estimating S_m for this assumed size distribution arises because $\langle m \rangle$ is infinite for $q < 2$ and $\langle m^2 \rangle$ is infinite for $q < 3$. Of course, a divergence is not physically present because no impacting particle can have a mass exceeding that of the final planet; in fact, the divergence implies only that the mass of the largest impactor to strike the planet, which we again denote by m_1, is not small compared to M_p. A convenient way to eliminate the divergence is to evaluate the means $\langle m \rangle$ and $\langle m^2 \rangle$ at fixed m_1. It is straightforward to show (*18*) that as the number of impacts $N \to \infty$

$$\langle M_p/m_1 \rangle = \left\langle \sum_{i=1}^{N} m_i / \max_{i=1}^{N} m_i \right\rangle = 1/(2-q) \quad (11)$$

for $1 < q < 2$, where $\max_{i=1}^{N} m_i$ is the maximum value of m_i for $i = 1$ to N. This implies that for fixed m_1, we have $N\langle m \rangle = m_1/(2-q)$. Similarly, $N\langle m^2 \rangle = 2m_1^2/(3-q)$ for $1 < q < 3$. Thus, if the averages in the definition of S_m (Eq. 2b) are taken at fixed m_1, we obtain

$$S_m = \left(\frac{2}{3-q}\right)^{1/2} (2-q)$$

$$= \left(\frac{2}{3-q}\right)^{1/2} \frac{m_1}{M_p}, \quad 1 < q < 2 \quad (12)$$

It might seem more natural to take the averages at fixed planetary mass M_p, but that approach is more complicated mathematically and does not yield any worthwhile gain in rigor. For $q \approx 2$ (which we argue below is the case for solar system planetesimals), $S_m = 1.4 m_1/M_p$, nearly the same value that is obtained when a single impactor dominates the mass (Eq. 9). Thus, the relation between S_m and m_1 is not sensitive to the precise choice of the mass distribution of planetesimals.

A by-product of these results is that we can sharpen our estimate of q by examining several lines of evidence regarding the masses of the largest impactors that have struck the terrestrial planets in the course of their formation. (i) Models for the formation of the moon by a giant impact on Earth (*19–21*) imply that the impactor's mass was in the range 0.1 to $0.3M_\oplus$, where $M_\oplus$ is the mass of Earth (crudely, this impactor mass is required so that the angular momentum associated with an impact parameter of $0.5R_\oplus$, where $R_\oplus$ is the radius of Earth, and an impact speed equal to Earth's escape speed is equal to the present angular momentum in the Earth-moon system). Models for the collisional stripping of Mercury's mantle (*22*) require a similar impactor-to-planet mass ratio. (ii) Simulations of the formation of the terrestrial planets by colliding planetesimals (*23*) suggest that planets like Earth have been struck by impactors of mass $m \gtrsim 10^{27}$ g $\simeq 0.2\ M_\oplus$. (iii) Chyba (*14*) has inferred that the largest lunar basin, South Pole–Aitken, was produced by an impactor of mass 1.4×10^{22} g after solidification of the lunar crust. For comparison, geochemical estimates of the total mass of the meteoritic component of the lunar crust yield a total of 4 to 15×10^{22} g. Thus, a fraction 0.1 to 0.3 of the mass accreted in the late, heavy bombardment was delivered in a single large impact; the agreement between this fraction and the fractional mass deduced in arguments (i) and (ii) suggests that the power-law index q is roughly constant over some five orders of magnitude in mass. All of these estimates yield m_1/M_p typically in the range 0.1 to 0.3; using $\langle M_p/m_1 \rangle = 1/(2-q)$, we deduce that q is in the range 1.7 to 1.9. The corresponding range of S_m is 0.15 to 0.4. Of course, there is considerable uncertainty in all these estimates (*24*).

The relatively large values of m_1/M_p that we find provide a consistency check on the power-law model for the impacting masses; much smaller, but still nonzero, values of m_1/M_p would require improbable fine tuning of q to be very near 2. Many investigators have worked with a more general mass model in which there is an upper cutoff to the power-law distribution, but this additional complication is not required by the observations. Indeed, if there were a real cutoff, m_1 should be roughly constant for each planet, whereas the arguments above and other estimates (*5*) suggest that m_1/M_p is much more nearly constant than m_1, which is what we would expect for a mass distribution following Eq. 10 with no cutoff.

For stochastic accretion, the typical rotation rate may be estimated from Eqs. 1, 2b, 6, and 12

$$|\mathfrak{R}| \approx \frac{S_m}{r^{3/2}} \approx \frac{m_1}{M_p} r^{-3/2} \quad (13)$$

This estimate is only a crude one because we have neglected factors of order unity and because of the statistical uncertainty in $\mathfrak{R}$ (the quartiles of the distribution of S_m are about a factor of 2 apart).

We now apply these results to the terrestrial planets. Of the four planets, Mars likely remains closest to its primordial rotation state. Mars has dimensionless radius and rotation rate, $r = 0.0022$ and $\mathfrak{R} = 670$. Two independent arguments then imply that the martian spin arises from stochastic accretion: (i) For Mars, $\mathfrak{R}r = 1.5$, well outside the range given by Eq. 7 for ordered accretion [this difficulty has already been noted by Lissauer and Kary (8)]; (ii) Eqs. 2a and 8a imply that accretion is stochastic if the spin is prograde and

$$S_m \gtrsim 0.3r^{1/2} \qquad (14)$$

which for Mars implies $S_m \gtrsim 0.015$, a condition that is consistent with our earlier conclusion that the observations indicate $m_1/M_p = 0.1$ to 0.3. Thus, the rotation of Mars cannot arise from ordered accretion. If accretion is stochastic, then Eq. 13 implies $m_1/M_p \approx 0.1$, consistent with the range deduced from observations. Models in which accretion is ordered and only the obliquity of Mars arises from stochastic impacts predict significantly smaller values, $m_1/M_p = 0.006$ to 0.02 (5).

The giant impact hypothesis for the origin of the moon already implies that spin accretion for Earth was stochastic, because most of the angular momentum in the Earth-moon system arises in this single impact. Even if we neglect this strong argument for stochastic accretion and consider Earth as an isolated body with its present spin, we find (i) $\mathfrak{R}r = 1.1$, well outside the range allowed for ordered accretion, and (ii) Eq. 14 implies that accretion is stochastic if $S_m \gtrsim 0.015$, which is easily satisfied for distributions with m_1/M_p in the observed range.

Venus has an obliquity near 180° and a sidereal spin period of 243 days. The planet has probably been slowed to its present rate of spin by atmospheric and body tides. It is difficult for these tides to reverse the spin direction, and hence, it appears necessary to conclude that the primordial spin of Venus was retrograde (25, 26). The retrograde spin could result either from stochastic accretion or from ordered accretion in a low-dispersion disk.

Mercury's spin is prograde and trapped in a spin-orbit resonance (the spin period is 2/3 of the orbit period). Trapping at this resonance could have arisen as the spin of Mercury was slowed by solar tides (27–29). It is even possible, though unlikely, for an initially retrograde spin to be accelerated through the synchronous state to this resonance. Thus, Mercury's present spin contains little or no information about its primordial spin. Therefore, the accretion of Earth and Mars is probably stochastic, and the accretion of Venus and Mercury may have been stochastic.

A possible argument against stochastic accretion is that the obliquities ϵ of Mars and Earth are relatively small. For Mars, $\epsilon = 24°$ at present, although the obliquity oscillates over a range of at least 30° owing to a near resonance between the frequencies of the spin axis and node precessions (30). For Earth, $\epsilon = 23°$ at present; studies of the tidal evolution of the Earth-moon system suggest that $\epsilon \approx 10°$ when the moon was near 10 $R_\oplus$ away, as opposed to 60 $R_\oplus$ today (31, 32). To help assess whether low obliquities threaten our conclusions, we determine the expected distribution of obliquities resulting from stochastic accretion.

For large dispersion ($s >> 1$), the final pole direction of the planet will be isotropically distributed, which implies a distribution of obliquities

$$n(\epsilon)d\epsilon = \frac{1}{2}\sin\epsilon\, d\epsilon, \quad 0 \le \epsilon \le \pi \qquad (15)$$

The derivation of the obliquity distribution is more subtle if the dispersion is low ($s << 1$). We define the Kepler sphere of radius R_K around the planet to be the region within which planetesimal orbits are nearly Keplerian; R_K is on the order of, but smaller than, the Hill radius R_H. Since the planet radius R_p is generally much less than R_H, planetesimals impact the planet if and only if their specific angular momentum at the Kepler sphere satisfies $|\mathbf{l}| < (2GM_pR_p)^{1/2} \equiv \ell_p$. Let (θ, ϕ) be spherical polar coordinates on the Kepler sphere. The thickness of the disk is small compared to R_K if $s << 1$, so $\theta \simeq \pi/2$. The components of the specific angular momentum of a particle crossing the Kepler sphere with tangential velocity (v_θ, v_ϕ) are therefore $\ell_z = R_K v_\phi$ and $\ell_\perp = R_K v_\theta$. If the velocity dispersion σ in the disk satisfies $R_K\sigma >> \ell_p$—which is true if $s >> r^{1/2}$—then the impacting particles will have approximately uniform distributions in v_θ and v_ϕ and, hence, in $\ell_\perp$ and ℓ_z, implying a distribution of angular momentum given by

$$n(\ell_\perp, \ell_z)d\ell_\perp d\ell_z = \frac{d\ell_\perp d\ell_z}{2\pi GM_pR_p}$$

where

$$\ell_\perp^2 + \ell_z^2 < 2GM_pR_p \qquad (16)$$

and zero otherwise.

If the angular momentum is dominated by a single impactor, it is straightforward to show from Eq. 16 that the distribution of obliquities is

$$n(\epsilon)d\epsilon = \frac{d\epsilon}{\pi}, \quad 0 \le \epsilon \le \pi \qquad (17)$$

Thus, the probability that a planet has obliquity $\epsilon < 24°$ (or $> 180° - 24°$) is only 0.09 if the pole position is distributed isotropically (Eq. 15) but has the larger value of 0.27 for stochastic accretion from a low-dispersion disk if the spin is determined by a single impact (Eq. 17). With this larger probability, it is not unreasonable to find planets with obliquity as low as 24°.

If planetary spins arise mainly from ordered accretion, the obliquity of Mars is usually explained as the result of one or more off-center impacts, that is, as stochastic "noise" superimposed on the ordered spin. Because the angular momentum contributed by the two processes is generally quite different, this explanation requires fine tuning so that the stochastic component is large enough to create a detectable obliquity but small enough so that the dominant accretion is still ordered.

To summarize, the following arguments imply that the spins of the terrestrial planets were determined by stochastic, not ordered, accretion: (i) the primordial spin of Venus was retrograde; thus, there is no statistically significant preference for prograde spins; (ii) the rotation rates of Mars and Earth are too large to arise from ordered accretion in a uniform disk; (iii) stochastic accretion is expected to dominate if the mass distribution of the impacting planetesimals is the one implied by modern studies of solar system formation.

The role of stochastic accretion in determining planetary spins and obliquities has already been stressed by several investigators, including Safronov (3), Hartmann and Vail (4), Lissauer and Safronov (5), and Kaula (6). This paper extends this work and argues that ordered accretion plays no significant role in the formation of the terrestrial planets. We believe that our arguments are compelling, but several unresolved issues remain: (i) As Lissauer and Kary (8) and Ida (33) have stressed, depletion of the protoplanetary disk around the growing protoplanet may increase the rate of ordered spin accretion. Lissauer and Kary suggest that preferential accretion of planetesimals from the edge of the planet's accretion zone may lead to rapid prograde rotation even if accretion is ordered, thus relaxing the upper bound in Eq. 7. However, rapid rotation only results if the disk is strongly depleted near the planet, and strong depletion has only been demonstrated under restricted conditions, for example, when the velocity dispersion of planetesimals is low and planetesimal scattering is dominated by a single protoplanet. At late times, when planetesimals can be scattered by more than one planet, the orbital distribution of the impactors is likely to be more uniform. Even if depletion is present, the other arguments presented above still imply that accretion is stochastic. (ii) The relatively low obliquities of Earth and Mars, while not inconsistent with stochastic accretion, are unexpected in this process (*32*) and may suggest that their obliquities have evolved since formation. (iii) We have assumed that each impactor is completely accreted when it strikes a planet; in large impacts, particularly at high velocities, fragmentation and loss of ejecta with high angular momentum may be important (*34*). (iv) The spin rates of Jupiter and Saturn, which probably arose during the accretion of their gas envelopes, are not understood theoretically and merit detailed hydrodynamic calculations. In contrast, our results may apply to Uranus and Neptune, ~90% of whose mass consists of solids. Applying Eq. 13 implies $m_1/M_p \approx 0.15$, similar to the values deduced for the terrestrial planets. The large obliquities of Uranus (98°) and Neptune (29°) are also consistent with stochastic accretion, although the obliquities of the giant planets may alternatively have arisen from rotation of the ecliptic plane after planet formation was complete (*7*).

Note added in proof: A new preprint by Laskar and Robutel (*35*) argues that the obliquity of Mars wanders chaotically over the range 0° to 60° and that the obliquities of all the terrestrial planets may have passed through chaotic states in the past. Their results strengthen the case for stochastic accretion by implying that the primordial obliquity of Mars may have been as large as 60°; however, chaotic obliquity evolution provides a possible mechanism for reversing the spin of Venus, so the primordial spin of Venus may have been prograde. Chaotic obliquity evolution does not affect arguments for stochastic accretion based on spin rates because it changes only the orientation, not the magnitude, of the spin vector.

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Mechanisms in the Competitive Success of an Invading Sexual Gecko over an Asexual Native

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The competitive displacement by a sexual gecko species of an asexual resident gecko has been documented over a wide geographic area. To test hypotheses concerning the detailed mechanism of this displacement, an experimental system was developed to follow populations of geckos in a duplicated, controlled environment that closely approximates the natural arena for the competitive interaction. Asymmetric competition occurred only in the presence of light, which attracts a dense concentration of insect food sources. The mechanism of competition was partly due to the behavioral dominance of the larger sexual species over the smaller asexual species in areas near the concentrated food. However, this behavior resulted from an avoidance response of subordinate asexuals rather than overt aggression by the sexual species.

The mechanisms that enable exotic species to thrive at the expense of native species are often unclear. There are many examples of the decline of native species after the arrival of an exotic species (*1, 2*). A competitive mechanism is frequently proposed to explain such phenomena, but rarely has such a mechanism been isolated and tested in an experimental setting, especially in vertebrates.

We have documented the recurrent human-aided arrival and distribution of the house gecko *Hemidactylus frenatus* to islands in the tropical Pacific Ocean and the concomitant numerical decline of species that previously occupied buildings on these islands (*3*). The invader *H. frenatus* is a sexual species, whereas at least two of the species it supplants are asexual parthenogens, *Hemidactylus garnotti* and *Lepidodactylus lugubris*. Around the time of World War II *H. frenatus* reached Oahu, Hawaii, and reached Fiji, Vanuatu, and Samoa probably in the last 20 years (*2*). It was first recorded in Tahiti in 1988 and on the Micronesian islands of Arno Atoll, Ponopei, and Kosrai this year (specimens are in the California Academy of Science).

Documentation of the decline of gecko species that previously occupied the nocturnal, insectivorous house gecko niche comes from two sources: a comparison of historical collection records to current census surveys (*3*) and a 5-year experiment in Suva, Fiji (*4*). On islands like the Cook Islands that have yet to be colonized by *H. frenatus*, *L. lugubris* is extremely common on buildings, whereas *H. garnotti* is patchy in distribution and abundance. These studies present unambiguous evidence of the strong dominance of *H. frenatus* over *L. lugubris*. How-

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