

Absorbing states govern the charging of vibrated grains

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The electrostatic charging of insulators is notoriously irreproducible. Using experiments and simulations we study the charging of a single bouncing ball on a vibrating surface. We show that a ball's final charge results from a coupling between its chaotic dynamics and its evolving charge distribution. This results in a permanent charge dipole which traps the particle in an “orientational absorbing state” preventing further charging. The variability of each particle's final charge is a direct consequence of this dynamical coupling.

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Electrostatic charging is ubiquitous in the flow and agitation of powders and grains. The charge generated through repeated interparticle collisions modifies cohesion [1,2], friction [3,4], and dissipation [5]. These interactions govern collective phenomena such as shifting sand dunes [6] and the dynamics of planetary regolith [7]. Industrially, charging promotes processes such as drug segregation in pharmaceutical tablets [8] and hinders materials separation [9,10]. Therefore, predicting the expected charge on a particle is an important challenge.

While the charge transferred during a collision depends upon the material and surface properties [11–14], the total charge on a surface is fundamentally limited. Specifically, the electric field arising from a distribution of charges may exceed the dielectric strength of the air, causing it to breakdown [13,15–21]. However, quantifying the contributions to this field such as geometry, polarizability, and surface charge is far from trivial [22].

Insulating particles add an additional degree of complexity since, unlike in conductors, the surface charge remains static and localized. After a few initial collisions an asymmetric charge distribution develops [23]. But with sufficient agitation it is often assumed that particles will eventually develop a quasihomogeneous state, which can be modeled by a central charge approximation [10,22,24]. Yet evidence exists that permanent charge dipoles *can* arise from prolonged shaking [5,25,26]. These nonuniform charge distributions result in torques and modify interaction strengths [26,27].

Inconsistent tribocharging in granular systems is frequently attributed to inherent uncertainties: imperceptible material impurities, the nonequilibrium nature of the process, or fluctuating environmental conditions [11,14,28]. This raises a fundamental question: is the emergence of permanent dipoles, and the broader phenomenon of charge variability, just a

consequence of these stochastic influences? Here, we demonstrate that the evolving surface charge distribution on an insulating particle can break symmetry spontaneously, even in an idealized system. This process arrests further charging, resulting in a final distribution that is inherently irreproducible due to a fundamental dependence on the particle's unique dynamical history.

Our primary experiment investigates a single polypropylene ball ($R = 5$ mm) bouncing on the concave surface of a 100-mm-diameter watch glass. The curvature of the watch glass restricts the horizontal motion of the ball but allows it to rotate freely. We adopt this setup because it captures the essence of the “caged” dynamics exhibited by a single particle within a vibrated granular pile: rotation and repeated contacts occur, but translation is restricted.

The surface is covered by an inverted Petri dish to protect from stray dust particles. Balls were cleaned in detergent, rinsed in ultrapure water, and dried under a flow of N_2 ($Q < 0.03$ nC). The charges on balls were measured using a Faraday cup. The watch glass is vibrated vertically at a frequency of 40 Hz and a dimensionless acceleration of $\Gamma \sim 2.5$. This induces chaotic bouncing and results in charging of the ball. Individual balls shaken for 5 min exhibit a wide distribution of final charges, ranging from 1.5 to 3.5 nC [Fig. 1(a)].

The mean charge across multiple balls saturates at ~ 2.5 nC within ~ 1 min [Fig. 1(b), $\bullet$], with prolonged shaking (~ 15 hr) yielding no significant increase. It might be assumed this steady state represents the maximum possible charge. However, this is inconsistent with experiments in which we add a low-frequency horizontal oscillation ($f \sim 1$ Hz, $A \sim 2$ –12 mm), resulting in the measured charge almost doubling [Fig. 1(b), $\bullet$]. Therefore we first need to understand what is the theoretical maximum charge of an insulating sphere.

Our primary experiment is complicated by the presence of two insulating materials (glass and polypropylene), both of which can potentially hold a charge. We therefore measure the charge on balls individually vibrated vigorously in a grounded copper pot. This removes any potential influence of the surface charge of the container. To ensure that our results are general and not system specific, we also use balls of different materials [Fig. 2(a), (POM $\blacklozenge$, PTFE $\blacksquare$, PP $\bullet$)] and radii (R). The measured charges, Q , were then used to calculate the

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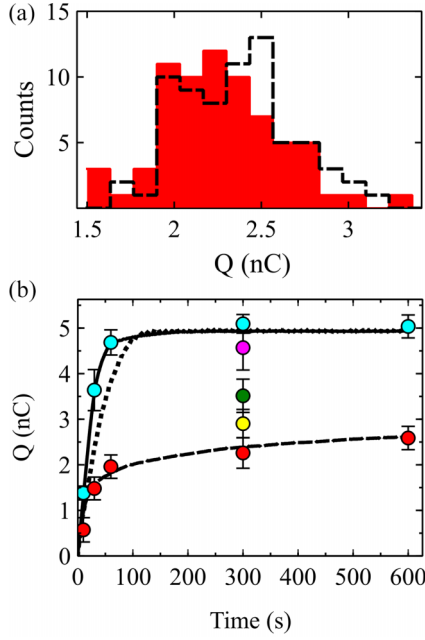


FIG. 1. (a) Charge measurements for balls after 5 min in experiment (solid) and simulation (dashed). (b) Mean charge of 10 balls, with 0 mm (●), 2 mm (●), 4 mm (●), 6 mm (●), and 12 mm (●) horizontal motion. Error bars represent the standard deviation. Lines represent the average of 100 simulations with (solid line) and without horizontal motion (dashed line) and a simulation with no forces acting on the ball due to charges (dotted line).

ball's surface field assuming $E_s = Q/(4\pi\epsilon_0 R^2)$. To calculate the maximum charge on a sphere bouncing on a conducting surface we follow the analytical approach outlined by Lowke and Alessandro [29], making two extensions. First, the field in the gap is the sum of the contributions from the sphere and its image charge. Second, the image charge in turn also induces a polarization charge on the particle, further enhancing the field in the gap. Consequently, as the ball recedes from the surface, the surface field that would result in breakdown passes through a minimum as the gap increases (see Supplemental Material [30] for further details). The surface field (E_s) at this critical gap defines the maximum sustainable charge on the sphere:

$$E_s = \frac{E_0}{2\eta(\epsilon_p)} \left[1 + \left(\frac{2K}{E_0 A R} \right)^{1/2} \right]^2. \quad (1)$$

Here E_0 is the critical electric field, A is a fitting parameter related to the first Townsend coefficient [29] and K is an ionization constant [31]. The factor $\eta(\epsilon_p)$ is a correction factor that accounts for the field due to the polarization of the sphere. For a dielectric constant $\epsilon_p = 3$, $\eta(\epsilon_p) \approx 1.6$ [25,32].

Equation (1) provides an analytical prediction of the particle surface field/maximum charge which agrees well with our data [see Fig. 2(a)]. This also demonstrates that the maximum charge on a sphere is controlled by all contributions to the electric field in the gap [18]. Therefore, replacing the conducting surface with an insulator should modify the maximum charge by changing the total electric field.

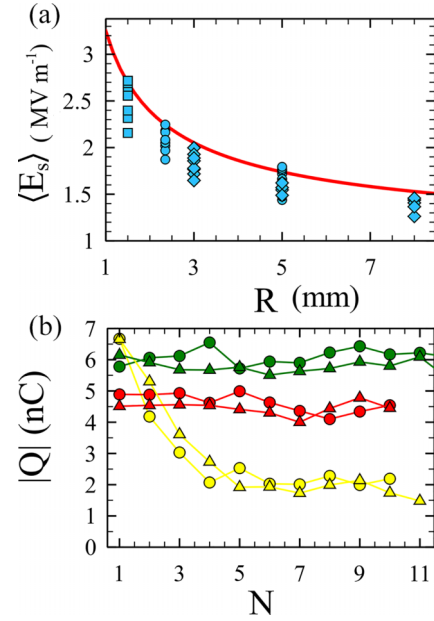


FIG. 2. Maximum surface field for different particle radii and surface materials. (a) The surface field of plastic particles (POM ♦, PTFE ■, PP ●) charged in a copper pot. The line represents the theoretical maximum field between the particle and the surface assuming a particle dielectric constant $\epsilon = 3$. (b) The measured charge of PP particles bounced for 5 min consecutively in pots of copper (●, ▲), PTFE (●, ▲), and glass (●, ▲). All beads were discharged before use, but the pots were not discharged between experiments, allowing the potential accumulation of surface charge.

In a series of 5-min measurements, individual $R = 5$ mm polypropylene (PP) balls were discharged and added consecutively to pots made of copper (grounded), glass ($\epsilon_s \sim 5$ –10), or PTFE ($\epsilon_s \sim 2.1$). Crucially, the pot was not discharged between each measurement, allowing the possibility of surface charge accumulation. Figure 2(b) shows that, in the copper pot, the saturation charge on the ball was consistently $Q_b = 4.7 \pm 0.2$ nC. In contrast, whilst measurements in a PTFE pot initially yielded $Q_b \sim 6.7 \pm 0.2$ nC, the measured charge then decreased significantly with each subsequent measurement.

The total electric field in the gap arises from the sphere's charge Q_b , the accumulated surface charge Q_s , and a dielectric image charge Q_{img} produced by polarization of the surface material. The image charge in the dielectric surface is $Q_{\text{img}} = -\beta Q_b = -Q_b(\epsilon_r - 1)/(\epsilon_r + 1)$ [33]. The image charge in PTFE produces a field ($\beta = \frac{1}{3}$) smaller than that of a metal ($\beta = 1$), resulting in a higher maximum charge Q_b required to reach the breakdown of air. However, in subsequent experiments, as the PTFE surface charge increases, the breakdown threshold of air in the gap is reached at progressively lower values of Q_b [Fig. 2(b)]. PTFE is known to have very low surface conductivity [34], which would mean that surface charge persists between consecutive experiments. After $N \sim 7$ experiments each ball obtains a small but nonzero charge, much smaller than that of the initial experiment. This represents a situation where the charge transferred to the pot per cycle is balanced by the losses.

Metal and PTFE represent limiting cases for understanding the maximum charge: the former is governed purely by the image charge, and the latter is strongly affected by the buildup of surface charge. To understand our primary experiment, we consider how a glass surface relates to these two categories. Figure 2(b) shows that the initial ball charged in a glass pot ($\frac{2}{3} < \beta < \frac{9}{11}$) has a charge which is higher than that of a metal ($\beta = 1$) and lower than that of PTFE ($\beta = \frac{1}{3}$). This is consistent with the role of the image charge in limiting the field. Contrary to the PTFE results, the subsequent balls bounced in a glass pot have a consistent charge, reminiscent of the metal experiments. This implies that significant amounts of surface charge do not accumulate, which is consistent with the higher surface conductivity of glass [34].

In our primary experiment, we observed a near maximal charge when a horizontal oscillation was added to the shaking [Fig. 1(b)]. This motion causes the ball to make contact with an increased area of the glass surface. However, we have shown that glass operates in a “metal-like” regime where there is no significant accumulation of surface charge. This raises a significant question: why should the horizontal motion of the surface have any affect on the particle’s final charge (see Fig. 1)?

We propose that the final charge depends on the degree of coupling between the ball’s chaotic dynamics and its charging history. Since charge on the ball remains localized [26], the initial contacts create a permanent electrostatic dipole. This dipole generates a restoring torque, which makes subsequent contacts with charged regions of the ball increasingly likely. This stochastic feedback mechanism can stabilize specific orientations, spontaneously breaking rotational symmetry and trapping the particle in a lower charge state. This “orientational absorbing state” [35] protects uncharged regions of the ball, preventing the particle from reaching the theoretical maximum charge.

To demonstrate this, we marked clean balls with distinctive colored spots and tracked their orientation using top- and side-mounted cameras (Supplemental Material Video 1 [30]). While the ball initially rotates freely, within ~ 30 s the same colored markers become fixed in the cameras’ fields of view. Despite the ball continuing to bounce vigorously, the particle enters the predicted orientational absorbing state. This prevents further charge growth, even over many hours. To test for rotational bias, the dish was then treated with a dry antistatic spray to suppress charging (< 0.1 nC in 5 min). As a neutral ball bounces on this surface, it rotates continuously. This confirms that any rotational arrest of the particles is an emergent property of the charging process rather than a geometrical constraint.

Electrostatic charging is notoriously sensitive, making it difficult to isolate the effects of minute environmental fluctuations, cleaning protocols, or material differences [11,14,28]. To test the proposed dynamical mechanism without such influences, we now perform simulations (see Supplemental Material [30]). Our simulation replicates the experiment using a sphere discretized into 128 surface patches bouncing on a vibrating curved surface. The patches are distributed uniformly across the ball’s surface and move with the sphere as a solid body. Charge is added to the contacting patch at each collision until the local electric field reaches the breakdown threshold.

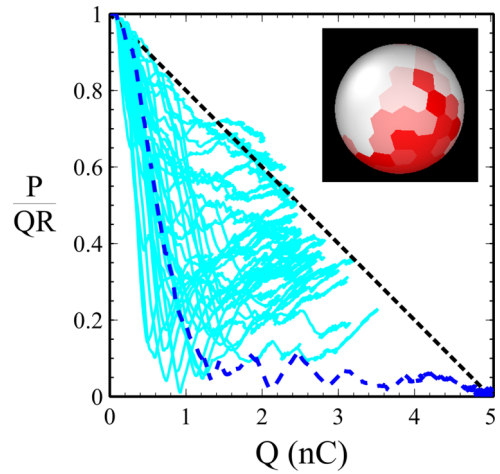


FIG. 3. How charge anisotropy (P/QR) relates to a balls total charge. Trajectories from the full simulation (light blue) and a “no coupling” simulation (dark blue) removing Coulomb forces. All particles initially follow a trajectory similar to that of the no coupling simulation. The final charge and anisotropy of the ball are largely set by the point of deviation. The black dotted line is a guide for the eye representing the approximate maximum charge for different dipole moments.

The total force and torque are calculated by summing the interactions between all patch charges (q_i) and all their dielectric image charges ($\beta = 0.8$). To avoid unphysical singularities at small separations, we model the interaction as that of a point charge with a uniform charge-density disk. This produces Coulombic behavior at large distances while converging to a finite value at contact. Consistent with our glass pot results, we neglect the influence of accumulated surface charge on the dish.

In our simulation the bouncing ball initially rotates freely, occasionally inverting its orientation (Supplemental Material Video 2 [30]). However, the stochastic nature of early collisions inevitably creates an asymmetric charge distribution (inset Fig. 3). The resulting dipole interacts with its own dielectric image, generating a restoring torque that guides the ball toward a preferred orientation. Within ~ 60 s, the chaotic dynamics and electrostatic feedback push the ball into an orientational absorbing state, preventing further charging.

The coupling of the ball’s dynamics with its particular charge configuration results in a different stochastic trajectory for each individual run. Applying a small perturbation at the start of each simulation, we randomize the initial conditions while keeping all physical parameters identical. Despite this, the resulting charge trajectories exhibit significant variation, a direct consequence of the underlying chaotic dynamics. As shown by the dashed histogram in Fig. 1(a), the simulated charges range from 1.5 to 3.5 nC, mirroring the experimental distribution.

Figure 1(b) (dashed line) also shows that the average of 100 such simulations accurately captures the time evolution of the experiment. Whilst we assign a specific charge increment per collision, the charging rate of the ball is largely insensitive to this parameter. Individual patches saturate within a few bounces, showing that the time evolution is primarily limited

by rotational dynamics, specifically the time required for the ball to present uncharged patches to the surface.

We can understand the high variability of the final charge by considering its charging history. Figure 3 shows the evolution of the normalized dipole moment $\frac{P}{QR}$ against the charge of the ball Q . Here $P = |\sum_i q_i \mathbf{r}_i|$ [26]. The dark blue, dashed line in Fig. 3 shows the evolution of the charge when the effect of the electrostatic forces in the simulation is switched off. Initial bounces produce high asymmetry as charge accumulates on a few patches. However, the distribution rapidly homogenizes, leading to a uniformly charged surface. Simulations under the same conditions, but with electrostatic coupling switched on (light blue), initially follow the same curve as the no coupling simulations. However, as the charge increases it influences the dynamics, leading to deviations from the dark blue curve. This signifies the transition to rotational arrest.

Figure 3 shows that random initial perturbations result in very different final charges for each simulation. This highlights how the magnitude of the final charge is governed by chaotic dynamics. Figure 1(a) shows that this variation in charge matches that observed in our experiments. This strongly suggests that the dominant cause of “irreproducibility” in our experiments arises from dynamics and not sensitivity to inherent uncertainties.

Figure 1(b) shows the results from averaged simulations with no coupling (dotted line) or subjected to horizontal oscillations of the surface (solid line). In the first case electrostatic torques are removed and in the second case additional frictional torques overwhelm the dynamics. However, in our experiments, even when the amplitude of displacement of the dish was 6 mm, the measured charge was suppressed relative to the experiments using a 12-mm amplitude. Suppression of the maximum charge would be expected to occur when the existing rotational kinetic energy ($\frac{1}{2}I\omega^2$) is comparable to the work done in rotating the ball. This electrostatic energy scale is set by the Coulomb potential $\sim Q^2/4\pi\epsilon_0 d$, where Q is the typical charge on the bead and d is a length scale comparable to the size of the particle. In our experiments, the horizontal motion of the surface with amplitude A and frequency f has a maximum velocity of $V_s = 2\pi Af$. This causes the

ball to rotate with approximate angular velocity $\omega \sim \frac{V_s}{R}$, here $A = 10$ mm, $f = 1$ Hz, and $d = 5$ mm. Equating the work and rotational kinetic energy terms gives $Q \sim 0.4$ nC. This estimate is consistent with the apparent sensitivity of the final charge to horizontal amplitude of the surface in Fig. 1(b).

The rotational constraints imposed on a caged particle in an agitated granular bed are far more restrictive than those present in our model experiment. Each particle develops local charged contacts which produce torques inhibiting further rotation. Whilst positive and negative charges can develop between grains of the same material [12], the magnitude of charges that develop in mixtures are generally much larger. Such systems should be naturally predisposed to the development of permanent dipoles and exhibit significant charge variability. Charging of a particle will also be particularly pronounced in a vacuum, where breakdown of the air is prevented. In space, the possibility of highly charged particles might therefore result in permanent dipoles arising in less dense systems. Electrostatic charge has been proposed as an important mechanism for resolving the bouncing barrier problem, promoting cohesion of clusters [2,36]. Our work suggests an additional mechanism for the generation of dipole interactions, which are known to influence their formation [1].

Our study demonstrates that the final charge of a vibrated insulating particle is not a fixed material property, but depends on its dynamical history. The coupling of chaotic dynamics and electrostatic torques leads to irreproducible charging. These results are particularly relevant to systems where particles are confined, a situation common to many kinds of granular materials. Finally, our findings suggest that the ideas of statistical physics, such as chaos and dynamical arrest, could be fruitfully extended to the longstanding problem of electrostatic charging.

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Data availability. There are no publicly available research data or software supporting this manuscript. Requests for further information or data should be sent to the authors.

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