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Fast magnetic reconnection in laser-produced plasma bubbles

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Recent experiments have observed magnetic reconnection in high-energy-density, laser-produced plasma bubbles, with reconnection rates observed to be much higher than can be explained by classical theory. Based on fully kinetic particle simulations we find that fast reconnection in these strongly driven systems can be explained by magnetic flux pile-up at the shoulder of the current sheet and subsequent fast reconnection via two-fluid, collisionless mechanisms. In the strong drive regime with two-fluid effects, we find that the ultimate reconnection time is insensitive to the nominal system Alfvén time.

Magnetic reconnection [1, 2], the change of magnetic topology in the presence of plasma, is observed in space, laboratory, and, most recently, laser-produced high-energy-density (HED) plasmas [3–6]. Reconnection plays a key role in energy release by plasma instabilities, as in solar flares or magnetospheric substorms, and the change in topology allows the rapid heat transport associated with sawtooth relaxation in magnetic fusion devices [7]. The observation of reconnection in HED plasmas suggests that it may also play a role in inertial confinement fusion, and indeed recent work has now observed the generation of large-scale magnetic fields during inertial fusion implosions [8].

We investigate recent experimental observations [3–6] of fast magnetic reconnection between the HED plasma bubbles created by focusing terawatt (TW)-class lasers (\sim kJ/ns) down to sub-millimeter-scale spots on a plastic or metal foil. The foil is ionized into hemispherical bubbles that expand supersonically off the surface of the foil. Each bubble is found to self-generate a strong magnetic field of order megagauss (MG), which forms a toroidal ribbon wrapping around the bubble. If multiple bubbles are created at small separation, the bubbles expand into one another, and the opposing magnetic fields are squeezed together and seen to reconnect (Fig. 1). The rates of reconnection are observed to be fast, and unex-

plained by classical Sweet-Parker theory [3, 4]. Reconnection has also been observed between laser-produced plasma bubbles immersed in a background plasma and magnetic field [9], though we focus here on reconnection between HED bubbles driven by kJ-class lasers.

It is of great interest to bring these results in line with what is already known about reconnection. As mentioned above, there are a number of new features in these laser-driven experiments, such as the high energy density in the plasma and magnetic field. Perhaps their most notable feature is the very strong reconnection drive: the opposing magnetic fields are driven together by the expanding bubbles at sonic and super-Alfvénic velocities. Such a strongly-driven regime has not been previously accessible in experiments, but may aid understanding of reconnection at the Earth’s magnetopause (where the driver of reconnection is the super-Alfvénic solar wind) and other astrophysical contexts with high plasma β (accretion disks, stellar interiors), or with colliding supersonic, magnetized flows (supernovae remnants, or at the heliopause). At the magnetopause, for example, the strong inflow can drive “flux-pileup” reconnection. [10, 11].

Here, we report the results of the first fully kinetic, particle-in-cell (PIC) simulations with parameters and geometry relevant to these experiments. Our main aim has been to understand the apparent Alfvénic reconnection rates observed in the experiments. We demonstrate that it can be explained as a combination of two effects: (1) a strong pile-up of magnetic flux [12] at the shoulder of the current sheet, caused by the sonic inflow and corresponding ram pressure that drives substantial compression of the plasma and magnetic field between the two bubbles, and (2) the intervention of collisionless effects (Hall current and electron pressure tensor) when the current sheet width falls below the ion skin depth ($d_i = c/\omega_{pi}$, where c is the speed of light, and ω_{pi} the ion plasma frequency). Previous discussion of these experiments has proposed anomalous resistivity [5], or the Hall effect [6], but have not discussed pileup, which we find is an essential ingredient. To our knowledge, these laser-driven experiments are the first in which flux pile-up as well as collisionless effects play important roles, produc-

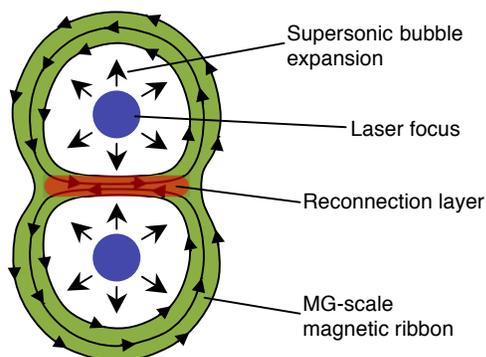


FIG. 1: Magnetic reconnection driven between expanding plasma bubbles, view from top.

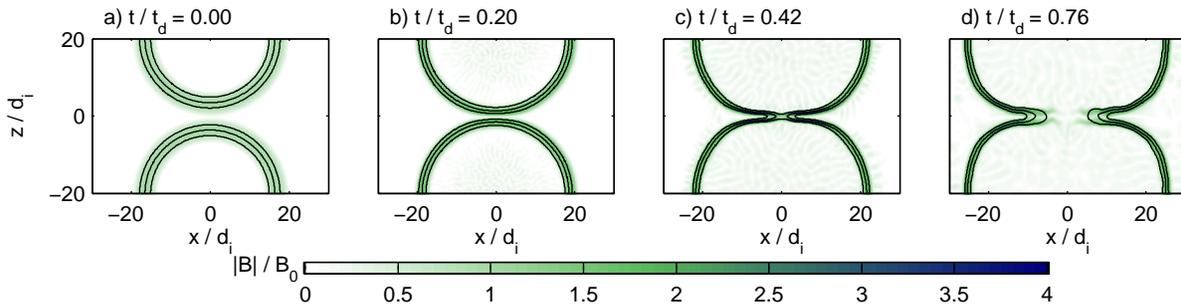


FIG. 2: Evolution of $|B|$ and contours of magnetic flux during the simulation.

ing highly time-dependent, impulsive reconnection dynamics, and Alfvénic peak reconnection rates.

In the laser-driven experiments, the reported reconnection rates were very fast compared to standard, classical predictions from Sweet-Parker theory, and instead were consistent with a simple hydrodynamic time based on the observed bubble expansion rate. The reconnection inflows (1.3 and 4×10^5 m/s for Rutherford [5] and Omega [4], respectively), ranged from about 0.7 to 3 times the associated Alfvén speed. ($V_A = B/\sqrt{\mu_0\rho}$, calculated based on B fields (100 T, 50 T), and mass density ρ derived from ion species (Al, CH) and electron density (both $\sim 5 \times 10^{25}$ m $^{-3}$), evaluated at the bubble edge, before the onset of the bubble collision.) Despite the difficulty making these novel measurements, the overall picture is that the reconnection is very strongly driven and occurs at Alfvénic to super-Alfvénic rates. Naturally then, the inflows are also much larger than allowed by classical Sweet-Parker theory, which allows inflows only a fraction of the Alfvén speed, $V_{SP} \sim V_A/S^{1/2}$, where the Lundquist number $S = LV_A\mu_0/\eta$, based on system size L and classical resistivity η , roughly 200 and 1000 for Rutherford and Omega. Meanwhile, the ion skin depths d_i (40 μm , 40 μm) are equal or larger than Sweet-Parker current sheet widths $\delta_{SP} = L/S^{1/2}$ (15 , 30 μm), suggesting that the experiments are in the regime for two-fluid, collisionless reconnection. The two-fluid reconnection mechanisms have been extensively observed in simulations [13, 14] and have begun to be verified experimentally [15]. Given this, however, the reported Alfvénic reconnection rates remain extraordinary, because even these fast, collisionless theories still typically find reconnection inflows “only” near 0.1 – 0.2 V_A .

We now present 2-d particle-in-cell simulations of bubble reconnection. The simulations track a pair of expanding bubbles through their interaction and reconnection, which is driven by the substantial energy stored in the plasma flow. Our simulations adopt an initial condition corresponding to a time about halfway through the experiments, after the bubbles are created, have expanded, and have generated their magnetic ribbons, but before the pair of bubbles has begun to interact. (It thus avoids the physics of the magnetic field generation pro-

cess, which is a 3-dimensional, $\nabla n \times \nabla T_e$ two-fluid effect [3, 4].) We then model the subsequent expansion and interaction of the plasma bubbles.

We define the following initial conditions. The system here runs, in (x, z) coordinates, over the domain $[-L_x, L_x] \times [-L_z, L_z]$. The system is periodic in both x and z for computational expediency. It therefore contains two half-bubbles, one centered at $(0, -L_z)$ and the other at $(0, +L_z)$. Define the radius vectors from the center of each bubble, $\mathbf{r}^{(1)} = (x, z + L_z)$ and $\mathbf{r}^{(2)} = (x, z - L_z)$. Then the initial density is $n_b + n^{(1)} + n^{(2)}$, where n_b is a background density, and the densities $n^{(i)}$ are

$$n^{(i)}(x, z) = \begin{cases} (n_0 - n_b) \cos^2\left(\frac{\pi r^{(i)}}{2L_n}\right) & \text{if } r^{(i)} < L_n, \\ 0 & \text{otherwise.} \end{cases} \quad (1)$$

Here L_n is the initial size of the bubbles. Next, the velocity field is initialized as the sum of the fields

$$\mathbf{V}^{(i)} = \begin{cases} V_0 \sin\left(\frac{\pi r^{(i)}}{L_n}\right) \hat{\mathbf{r}}^{(i)} & \text{if } r^{(i)} < L_n, \\ 0 & \text{otherwise.} \end{cases} \quad (2)$$

These simulations adopt uniform initial electron and ion temperatures (T_{e0} and T_{i0}) for simplicity. The magnetic field is initialized (Fig. 2(a)) as the sum of two toroidal ribbons, with

$$\mathbf{B}^{(i)} = \begin{cases} B_0 \sin\left(\frac{\pi(L_n - r^{(i)})}{2L_B}\right) \hat{\mathbf{r}}^{(i)} \times \hat{\mathbf{y}} & \text{if } r^{(i)} \in [L_n - 2L_B, L_n], \\ 0 & \text{otherwise.} \end{cases} \quad (3)$$

Here B_0 is the initial strength of the magnetic field, and L_B the half-width of the initial ribbons.

Most parameters can be matched to the experiment, except for a heavy electron mass and relatively low speed of light, as is typically the case for PIC simulations. Lengths are scaled to units of the ion-skin depth to correctly retain the two-fluid effects compared to the experiments. L_n ranges from 20 – 100 d_i , for the Rutherford and Omega experiments. L_B was found from experimental measurements and associated simulations [4] to be $\sim 0.1L_n$ in Omega. Note that the current sheet width

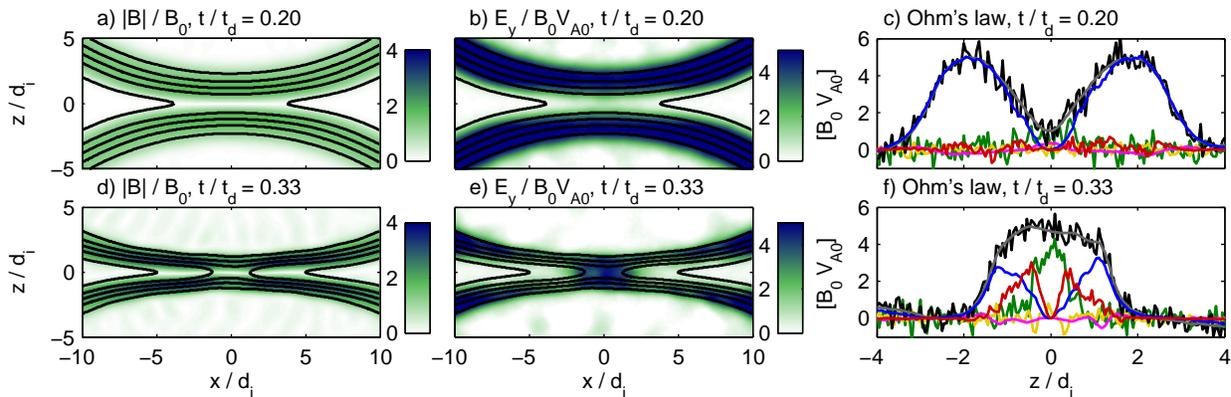


FIG. 3: (a,d) Magnetic fields $|B|/B_0$. (b,e) Reconnection electric fields $E_y/V_{A0}B_0$. (c,f) Generalized Ohm's law $E_y = -(\mathbf{V}_i \times \mathbf{B})_y + (1/ne)(\mathbf{j} \times \mathbf{B})_y - (1/ne)\partial P_{e,y}/\partial t - (1/ne)\nabla_\perp T_{e,y\perp}$. Black: sum of Ohm's law RHS; gray: E_y , blue: $-V_{i,z}B_x$ (with V_i the ion flow); red: $(1/ne)(\mathbf{j} \times \mathbf{B})$; green: $-(1/ne) \times \partial_x T_{e,yx}$; yellow: $-(1/ne) \times \partial_z T_{e,yz}$; pink: $-(1/ne)\partial P_{e,y}/\partial t$, where $P_{e,y}$ is the momentum density in the electron fluid. The stress tensor $T_{e,ij}$ includes the inertia terms $m_e n v_{e,i} v_{e,j}$ due to electron flow along the i and j directions.

is therefore of order the ion skin depth, which is essential for the two-fluid reconnection mechanisms discussed here. B_0 is scaled to the plasma pressure, to achieve the high-beta regime of the experiments.

Particle simulations were conducted with the relativistic, electromagnetic, explicit particle-in-cell code PSC[16]. We choose the PIC parameters $M_i/m_e = 100$, $T_{i0} = T_{e0} = 0.02m_e c^2$, and an initial 100 particles per cell. Total energy is conserved to better than 1% in the simulations reported. The results here are for collisionless simulations. While this is largely consistent with the above estimate that reconnection in these experiments is dominated by collisionless mechanisms, collisions are clearly also non-zero, and this will be an important part of future modeling as greater experimental fidelity is sought.

Figure 2 presents the evolution of the magnetic field in one simulation with parameters close to those of the Rutherford experiments [3], $L_n = 20d_i$, $n_b/n_0 = 0.1$, $L_B/L_n = 1/6$, $B_0^2/n_0 T_{e0} = 0.25$, $V_0 = 3(T_{e0}/m_i)^{1/2}$. The image backgrounds show $|B|$, over which are plotted contours of Ψ , which is defined such that $\mathbf{B}(x, z) = \hat{\mathbf{y}} \times \nabla \Psi + B_y(x, z)\hat{\mathbf{y}}$. Panel (a) gives the initial condition of Eq. 3. The next three panels show the evolution at later times, with time scaled in units of the dynamical time $t_d = L_n/V_0$. As the bubbles interact, a reconnection x-line forms at the leading point of tangency between the bubbles. Near $t/t_d = 0.42$, the reconnection is about half complete, and the reconnection rate is near its maximum. Notice the strong pileup of flux upstream of the reconnection region. In contrast to reconnection considered in other contexts, here the strong inflows have sufficient force to compress the flux in the current sheet by a ratio of about 4. This will be shown momentarily to have important consequences for the scaling of the peak reconnection rate. Finally, by $t/t_d = 0.76$, in (d), reconnection is complete, and the two bubbles have almost

completely merged. Notably, complete reconnection has occurred within about 1 dynamical time.

Figure 3 shows, at two characteristic times (a,b,c: $t/t_d = 0.2$, (d,e,f): $t/t_d = 0.33$), $|B|$, the reconnection rate $E_y = \partial \Psi / \partial t$, and a detailed cut of how the generalized Ohm's law is fulfilled along a cut across the x-line. The time $t/t_d = 0.2$ is characteristic of the compression or pileup phase: the reconnection rate is small, and the electric field is simply supported by the $\mathbf{E} \times \mathbf{B}$ plasma flow of the ribbons (c). However, by $t/t_d = 0.33$ the reconnection is now proceeding at its maximum rate. The electric field, of similar magnitude to that earlier in the flow, is now sustained (1) in the current sheet by the Hall effect ($\mathbf{j} \times \mathbf{B}/ne$), since the current sheet is now of the ion skin depth scale, and (2) at the x-line by the the electron stress tensor. Similar to typical results from 2-d PIC reconnection studies, within the stress tensor, the off-diagonal pressure component $(\partial_x T_{yx})/ne$ is dominant, accounting for 80% of the electric field.

We find that the reconnection rate $\partial \Psi / \partial t$ sustained at this time is super-Alfvénic compared to the *nominal* $B_0 V_{A0}$, where B_0 is given in the initial condition, and V_{A0} is evaluated with n_0 and B_0 . Furthermore, $E_y/B_0 V_{A0}|_{max} \simeq 4.7$ is about a factor of 50 higher than the typical $0.1BV_A$ found from reconnection simulations including the Hall effect. Here we have normalized to the nominal Alfvén speeds to correspond to the experiments: the magnetic fields measured (and simulated), e.g. in Fig. 4(c1) of Ref. [4], are in fact an *initial* condition before the bubbles have begun to interact.

This point is important because two-fluid reconnection mechanisms (the Hall effect and, in these PIC simulations, the electron pressure tensor), are typically found to allow for reconnection inflows of $0.1-0.2 V_A$. On its face, this is not sufficient to explain the apparent Alfvénic reconnection observed in the experiments. However, in

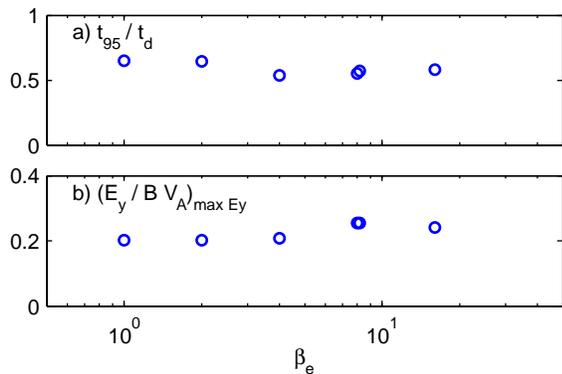


FIG. 4: (a) Time for complete reconnection, and (b) peak reconnection rate normalized to the instantaneous Alfvén speed, vs. initial plasma $\beta_e = 2n_0 T_{e0} / B_0^2$.

this strongly driven regime, the Alfvén speed is also time-dependent, due to the compression of the magnetic flux. Therefore we also track the *instantaneous* Alfvén speed using the maximum upstream magnetic field $|B|$ and the density in the current sheet. Indeed, once one accounts for the factor of 4 magnetic flux compression (a factor a 16 in BV_A), and the slightly reduced density at the current sheet ($n/n_0 \simeq 0.4$) one finds that $E_y/BV_A \simeq 0.25$, much more in line with typical two-fluid physics results. Therefore, we find that previously-established two-fluid physics results can also apply to this geometry, once one accounts for the flux pile-up.

Finally, we confirm the importance of flux-pileup in this strongly driven regime. Figure 4 shows a sequence of simulations in which we have varied the strength of the initial magnetic field B_0 , keeping other parameters fixed. The initial inflows become increasingly super-Alfvénic compared to the nominal Alfvén speed as the initial magnetic field is decreased. We then measure the magnetic reconnection rates in two ways. First, in Fig. 4(a) we measure the time for total reconnection, based on the time in the simulation to reconnect 95% of the initial flux [$\Psi_0 = (4/\pi)L_B B_0$]. Interestingly, we find that the total reconnection time only depends quite weakly on B_0 ; therefore in the strongly driven regime the reconnection time is insensitive to the nominal Alfvén speed.

At the same time, we can also study the system at the time of peak reconnection. Figure 4(b) shows the peak reconnection rates, normalized to the instantaneous Alfvén speeds, calculated as above. We find that with these corrections, over a wide range of initial magnetic fields and therefore flux pileup ratios, the ultimate reconnection rates are consistently about $0.2 BV_A$, roughly in line with previous two-fluid investigations.

In past research, flux pileup has been found to operate when reconnection is driven faster than allowed by the reconnection model (resistive MHD [12], or Hall MHD [17]). If the inflow rates are larger than the re-

connection rates, then flux must pile up. This is apparent in Fig. 3, where the reconnection rate at the x-line is less than the flux advection rate from the plasma flow. Meanwhile, the two-fluid physics result $E_y \propto BV_A \propto B^2$ finds that reconnection rates increase sharply with pileup, so that eventually these two effects can come into balance. However, pileup can only proceed if no other effects can stop the compression process, and an important consideration here would be the accumulation of plasma and magnetic pressure in the current sheet. However, with the view of the laser-driven experiments, where the initial current sheet densities are much smaller than the bubble core densities, we do not believe this would be sufficient to prevent reconnection.

In fact, our simulations find substantial pressure increase in the current sheet, mostly due to strong heating of the current sheet by compression. The observation of current sheet heating is qualitatively consistent with the Rutherford experimental results, which attributed the heating to magnetic energy dissipation [3]. We suggest instead that the converging plasma flows far dominate the magnetic fields as an energy source, and that the heating is from compression.

To conclude, recent experiments on magnetic reconnection in laser-produced plasma bubbles are found to be in a strongly driven reconnection regime unexplored in previous laboratory experiments, where a combination of flux-pileup and two-fluid effects account for extremely fast reconnection. The pile-up is found to be an enormous effect, boosting the relevant Alfvén speeds to match the specified plasma inflow. The strong inflow and flux pileup lead to a reconnection time independent of the nominal Alfvén time.

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