respectively. The reflected probe beam is collected using f/20 lens. The amplitude and polarization of the specularly reflected probe is simultaneously monitored by splitting it into two arms, one measuring the reflectivity with a photodiode (PD2) and the other measuring the polarization state using an analyzer (extinction ratio $10^{-5}$) in front of another identical photodiode (PD3), as shown in Fig. 1. The photodiode P1 measures the shot to shot laser fluctuations. All three PD signals are simultaneously recorded for each laser shot along with the delay stage position. The ellipticity is determined at various fixed delay positions by plotting PD3 signal with respect to analyzer angle. Higher temporal resolution ellipticity data is obtained by continuously varying the delay in steps of 1 pm at different fixed analyzer positions and computing the ratio PD3/PD2 for each delay. The ratio (PD3/PD2) is used above so as to account for plasma reflectivity variation as function of time delay. The temporal behavior of the magnetic field from -2 to +10 ps time delay is monitored. The constant background noise due to the second harmonic radiation from the plasma generation by the pump pulse is negligible (< 0.1%) as compared to the reflected probe.

Figure 2 presents the temporal evolution of the magnetic field. The inset shows reflectivity and induced ellipticity of the probe for a p-polarized pump. The sharp reflectivity dip is used to independently establish the start of the magnetic pulse. The magnetic field is derived from induced ellipticity ($\beta$) for our experimental conditions using $[12]$ $\beta(t) = 3.32 \times 10^{-26} \int n_e(t,t) B^2(t,t) dt$, where $n_e$ is the electron density in cm$^{-3}$, $B^2$ is magnetic field in MG, and $t'$ is the path length in $\mu$m. The magnetic field is deduced assuming a spatially uniform $B$ over a linear density gradient $n_e(x) = 0.4n_e^{403}(x/L)$, where $L$ is the plasma slab length, as shown in the inset of Fig. 1. The factor $0.4n_e^{403}$ corresponds to the turning point density for the 403 nm probe. We integrate over the trajectory in the plasma, $dl = \sqrt{1 + y'^2}dx$, where $y' = \sin \theta_0 / \sqrt{\varepsilon(x) - \sin^2 \theta_0}$, $\varepsilon(x)$ being the dielectric function. The plasma expansion velocity is estimated to be $5 \times 10^7$ cm/s from Doppler shift measurements of the reflected probe. With $L(t) = 1 / \varepsilon_{xy} t$, the magnetic field as a function of time delay is obtained as $B(t) = 80 \sqrt{\beta(t) / (1 + 0.5t)}$ MG. From the results shown in Fig. 2, the magnetic field pulse generated by a p-polarized pump has a peak value of 27 MG and duration (FWHM) of 6 ps. In comparison, s-polarized pump results in a peak value of 14 MG. This confirms the importance of RA, induced by p-polarization, in magnetic field generation. The magnetic field in case of s-polarized pump can be explained by being arising from parametric instabilities near critical density, which develop at a slower rate as seen in figure 2. Further, any realistic laser focusing algorithm involves some RA contribution even for s-polarization, due to critical surface rippling and geometrical effects [8, 15].

The mechanism of quasi-static magnetic field generation in ultrashort laser-plasma experiments can be understood by invoking the following magnetic field evolution equation.
\[
\frac{\partial F}{\partial t} = -\nabla \times \left( \frac{7}{4\pi} \cdot \frac{e}{n_e} \right) + \frac{e}{n_e} \left( \nabla T_e \times \nabla n_e \right) + \frac{e}{\sigma} \left( \nabla \times \nabla \nabla \nabla \nabla F \right) + \frac{e^2}{4\pi\sigma} \nabla^2 F.
\]

Equation (1) is derived by taking the curl of the equation of motion of background plasma electrons that carry the plasma shielding current \( J \). \( J_0 = \nabla \times \vec{B} - \vec{j}_{\text{hot}} \). The first term in equation (1) is the standard electron magnetohydrodynamic (EMHD) [16] source due to the Hall effect, the second term is the thermoelectric source and the third term is the source due to hot electrons (generated near the pump critical surface by the RA of the laser pulse); the last term gives the magnetic field decay due to resistive damping of the plasma shielding currents (\( \sigma^{-1} \), the background plasma resistivity). Electron inertia and magnetic field convection effects are assumed to be negligible in equation (1).

We estimate the first term (important during the first 200 fs when the pump is on) by calculating the cycle-averaged product of high frequency current and magnetic fields [17], \( \nabla \times \left( J \hat{B} \times \vec{B}_h / n_e \right) \), inside the plasma. We obtain the magnitude of this term by carrying out a particle-in-cell simulation using the laser-plasma interaction code LPIC++ [18]. In the simulation, a p-polarized light pulse with \( \sin^2 \) envelope and 100 fs FWHM laser pulse is incident at an angle of 55° on a linear density ramp. Fig. 3(a) shows numerically obtained spatial and temporal profiles of the quasi-steady-state magnetic field. This confirms the earlier analytical results of spatial profiles of laser generated magnetic fields [18]. At the same time, the values of \( B \) and \( \partial B / \partial t \) obtained from simulation compare well with our experimental results (Fig. 3b). The numerically obtained maximum value is \( \partial B / \partial t \) is 22 MG/ps at 120 fs, which is close to the experimentally deduced \( \partial B / \partial t \). The former value reached in 150 fs is 1.2 MG, beyond which the simulation results saturate. This is expected, because to model this regime we have considered only one term, which is valid strictly for the duration of the pulse.

The second term in equation (1) is assumed using a temperature gradient of 100 eV over a transverse scale of 15 μm (pump spot radius) and density gradient of 10^{21} cm^{-3} over 1 μm in the normal direction. This gives \( \partial B / \partial t \) ~ 0.05 MG/ps, which is much smaller than the experimental results. Hence, the thermoelectric source term is neglected in subsequent analysis.

The remaining terms (magnetic diffusion and hot electron source [8, 19]) in the magnetic field evolution equation govern the evolution of \( B \) after the pump pulse is removed. We model this regime by a phenomenological 0-dimensional evolution equation

\[
\frac{\partial B}{\partial t} = S(t) - \frac{B}{\tau}.
\]  

where \( S(t) \) is the source term. \( S(t) \) is mainly due to the fast electron currents and \( B / \tau \) is a 0-d representation of the magnetic diffusion term. Taking \( S(t) = S_0 \exp \left( -t / t_0 \right) \), we get

\[
B(t) = \frac{S_0}{1 / \tau - 1 / t_0} \left[ \exp \left( -t / t_0 \right) - \exp \left( -t / \tau \right) \right].
\] 

The first term in equation (3) denotes the natural decay of the hot electron source produced by RA mechanism. The second term describes the resistive decay of the fields generated by the plasma return currents. As shown in Fig. 2, this expression gives an excellent fit to our experimental data for p-polarized pump with \( S_0 = 53.7 \) MG/ps, \( t_0 = 0.7 \) ps and \( \tau = 5.6 \) ps.

To get an insight into these numerical values, we estimate conductivity \( \sigma \) from the magnetic diffusion term. Using \( \tau \sim (4\pi\sigma/e^2)(\Delta x)^2 \), for the best fit value \( \tau = 5.6 \) ps and \( \Delta x = 15 \mu m \), we get \( \sigma = 1.8 \times 10^{14} \) sec^{-1}. However the \( \sigma_{\text{solid}} \) or \( \sigma_{\text{classical}} \) observed [20] at \( T_e = 100 \) ev has value of \( \approx 4.5 \times 10^{15} \) sec^{-1}, which is an order of magnitude greater than that obtained from magnetic field decay! This clearly brings out the importance of turbulence induced anomalous resistivity effects in the damping of shielding plasma currents. An upper estimate of anomalous resistivity due to turbulence of electrostatic waves, obtained by taking the effective collision frequency \( \nu_{\text{eff}} \sim f_{\text{bulk}} \) (where \( f \) is a fraction of order unity), is in reasonable agreement with our estimate of \( \sigma \). We note that the hot electron source term is effective for about 0.7 ps, which is longer than the laser pulse (FWHM=100fs). Thus we argue that the electrostatic plasma waves generated by the RA mechanism (typically
with $E^2/4\pi n_r T_e \gg 1$) during the laser pulse will continue to slowly damp (or convect away) and accelerate electrons for a few hundred femtoseconds after the laser pulse. To get an estimate of $S_0$, which is the magnitude of the source term $[e/\sigma(\nabla \times j_{hot})]$, we evaluate $j_{hot} \approx f_0 eT/T_{hot}$ by taking a conversion fraction ($f_0$) of incident energy into hot electrons as 0.3 and $T_{hot} \approx 20$ keV (estimated by using the well known scaling laws for resonance absorption [21]). This yields hot electron current density $j_{hot} \sim 4.5 \times 10^{20}$ statampere/cm$^2$. Using anomalous conductivity obtained above we get $S_0 \sim 61$ MG/ps, which is again in close agreement with our phenomenological fit.

In conclusion, we have measured and characterized picosecond megagauss magnetic pulses generated by the interaction of ultrashort laser pulse with a solid. Our measurements extend to overdense region of the target and hence are of relevance to electron transport and fusion related issues. The experimentally observed rise times and magnitude of magnetic fields closely follow theoretical estimates and simulations. We also observe for the first time, anomalously rapid damping of return plasma shielding currents produced in response to the hot electron currents penetrating the bulk plasma; this is a topic of great significance to the fast ignition scheme. Such ultrashort, localized magnetic fields are useful for investigating magnetic precession and reversal dynamics, which is vital for developing next generation ultrafast switching and storage devices [22]. Further, the generation and characterization of these giant magnetic fields offers a unique opportunity for accessing extreme stellar conditions and testing astrophysical theories in the laboratory [23]. The diagnostics in chemical and biological sciences like Magnetic Circular Dichroism and Magnetic Resonance may also benefit in situations where high fields are essential [24]. The intensities used in our experiments are easily realizable with modern kilohertz repetition rate femtosecond lasers, and we foresee exciting applications for these magnetic pulses.

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[14] The effective phase difference is given as: \[ \delta \phi_{\phi} = \frac{2}{\pi} \int \left[ k(x) - \sin^2 \theta \right] dx - \frac{2}{\pi} \int \left[ \frac{k(x) - \sin^2 \theta}{\eta^2 + \left( \frac{\eta}{\rho} \right)^2 - \frac{3}{2} \left( \frac{\rho}{\rho_0} \right)^2} \right] dx. \]

For our parameters, we get maximum phase difference of 0.12 radian. In the worst-case scenario, where the plane of polarization is at 45° to the plasma gradient, we get ellipticity that is below the baseline (i.e., observed when probe alone is present and, hence, is subtracted from the data) value of 0.07.