Radiation-Reaction-Force-Induced Nonlinear Mixing of Raman Sidebands of an Ultraintense Laser Pulse in a Plasma

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Stimulated Raman scattering of an ultraintense laser pulse in plasmas is studied by perturbatively including the leading order term of the Landau-Lifshitz radiation reaction force in the equation of motion for plasma electrons. In this approximation, the radiation reaction force causes a phase shift in nonlinear current densities that drive the two Raman sidebands (anti-Stokes and Stokes waves), manifesting itself into the nonlinear mixing of two sidebands. This mixing results in a strong enhancement in the growth of the forward Raman scattering instability.

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Parametric instabilities of a laser pulse in a plasma are important due to their applications in the areas of laser-driven fusion and laser wakefield acceleration and have been investigated for decades [1–4]. Stimulated Raman scattering (SRS) is one of the prominent examples of parametric instabilities, carrying significant importance on account of being responsible for the generation of hot electrons in laser-driven fusion [1,2] and plasma wave excitation in laser wakefield acceleration [4]. In SRS, the incident pump laser decays either into two forward moving daughter electromagnetic waves (forward Raman scattering) or into a single backward moving daughter electromagnetic wave (backward Raman scattering) and a plasma wave. The daughter waves have their frequencies upshifted (anti-Stokes waves) and downshifted (Stokes wave) from the pump laser by the magnitude which equals the excited plasma wave frequency.

At high laser intensities $I_l \geq 10^{19}$ W/cm^2, the growth rate of the parametric instabilities becomes smaller due to the relativistic Lorentz factor [3]. However, at ultrahigh laser intensities $I_l \geq 10^{22}$ W/cm^2, the role of the radiation reaction force becomes important, too [5,6]. Such ultraintense laser systems are on the anvil after the commissioning of the Extreme Light Infrastructure (ELI) Project in Europe [7]. Because of radiation reaction force, the laser pulse suffers damping while propagating in a plasma. This damping of the laser pulse makes it vulnerable to plasma instabilities in following ways. First, as the laser loses energy due to the radiation reaction force, it facilitates, apart from the usual parametric decay processes, the availability of an additional source of free energy for perturbations to grow in the plasma. Moreover, its effective intensity decreases, which lowers the relativistic Lorentz factor. Second, the phase shift, caused by the radiation reaction force, in the nonlinear current densities can mediate the mixing of the scattered daughter electromagnetic waves which can now grow faster, utilizing efficiently the additional channel of the laser energy depletion. This necessitates the inclusion of the effect of the radiation reaction force in the theoretical formalism of the parametric instabilities in the plasma.

In this Letter, we include the effect of the radiation reaction force and study the SRS of an ultraintense laser pulse in a plasma, treating the radiation reaction force effects in the classical electrodynamics regime where quantum effects arising due to photon recoil and spin are negligible [5]. For this to be valid, the wavelength and magnitude of the external electromagnetic field in the instantaneous rest frame of the electron must satisfy $\lambda \gg \lambda_C$ and $E \ll E_{\rm cr}$, where $\lambda_C = 3.9 \times 10^{-11}$ cm is the Compton wavelength and $E_{\rm cr} = 1.3 \times 10^{16}$ V/cm is the critical field of quantum electrodynamics [5]. For the laser intensities planned in the ELI Project $I_l \sim 10^{22-23}$ W/cm^2 [7], these two criteria can be fulfilled. In the classical electrodynamics regime, the Landau-Lifshitz radiation reaction force correctly accounts for the radiation emitted by a relativistic charged particle [8]. We incorporate the leading order term of the Landau-Lifshitz radiation reaction force perturbatively in the equation of motion, focusing on the phase slippage caused by it on the quiver momentum of oscillating electrons. This phase shift due to the radiation reaction force tends to enhance the growth rate of the SRS instability. The growth of the forward Raman scattering (FRS) instability gets strongly enhanced while the growth of the backward Raman scattering (BRS) instability does not experience a strong enhancement due to the radiation reaction force.

We consider the propagation of a circularly polarized (CP) pump laser along the $\hat{z}$ direction in an underdense plasma with uniform plasma electron density n_0. Ions are assumed to be at rest as their motion leads to the appearance of an additional ion mode instability which does not couple with the SRS instability [9]. The equation of motion for an electron in the laser field including the leading order term of the Landau-Lifshitz radiation reaction force is

$$\frac{\partial \mathbf{p}}{\partial t} + \mathbf{v}\cdot\nabla\mathbf{p} = -e\left(\mathbf{E} + \frac{1}{c}\mathbf{v}\times\mathbf{B}\right) - \frac{2e^4}{3m_e^2c^5}\gamma^2\mathbf{v} \times\left[\left(\mathbf{E} + \frac{1}{c}\mathbf{v}\times\mathbf{B}\right)^2 - \left(\frac{\mathbf{v}}{c}\cdot\mathbf{E}\right)^2\right], \quad (1)$$

where $\gamma = 1/\sqrt{1 - v^2/c^2}$, e is the electronic charge, m_e is the electron mass, and c is the velocity of the light in vacuum. The other terms of the Landau-Lifshitz radiation reaction force are $1/\gamma$ times smaller than the leading order term [8] and can be ignored. We first solve this equation of motion by ignoring the radiation reaction term and by expressing the electric and magnetic fields in potentials as $\mathbf{E} = -\nabla\phi - \partial\mathbf{A}/\partial ct$, $\mathbf{B} = \nabla \times \mathbf{A}$. In a 1D approximation valid when $r_0 \gg \lambda_0$ (where r_0 is the spot size and λ_0 is the wavelength of the pump laser pulse), this yields the transverse momentum and z component of motion as $\mathbf{p}_\perp = e\mathbf{A}/c$, and $\partial v_z/\partial t = e\nabla_z\phi/(m_e\gamma_0) - e^2\nabla_z|A|^2/(2m_e^2\gamma_0^2c^2)$, where $\gamma_0 = (1 + a_0^2/2)^{1/2}$, $a_0 = eA_0/(m_ec^2)$, $\mathbf{A} = \mathbf{A}_0e^{i\psi_0}/2 + c.c$, $\mathbf{A}_0 = \boldsymbol{\sigma}A_0$, $\boldsymbol{\sigma} = (\hat{\mathbf{x}} + i\hat{\mathbf{y}})/\sqrt{2}$, $\psi_0 = k_0z - \omega_0t$, and ω_0 and k_0 are the carrier frequency and wave vector of the pump laser, respectively [3,10]. A plane monochromatic CP light has $\nabla|A_0|^2 = 0$, so it does not cause any charge separation, leading to no component of velocity in the $\hat{z}$ direction. This is the so-called Akhiezer-Polovin solution for a purely transverse monochromatic CP light in plasmas [3,11,12]. However, the scattering of the laser pulse leads to the total vector potential of the form $\mathbf{A} = [\mathbf{A}_0e^{i\psi_0} + \boldsymbol{\delta}\mathbf{A}_+e^{i\mathbf{k}_\perp\cdot\mathbf{x}_\perp}e^{i\psi_+} + \boldsymbol{\delta}\mathbf{A}_-^*e^{-i\mathbf{k}_\perp\cdot\mathbf{x}_\perp}e^{-i\psi_-^*}]/2 + \text{c.c.}$, where $\boldsymbol{\delta}\mathbf{A}_+ = \boldsymbol{\sigma}\delta A_+$, $\boldsymbol{\delta}\mathbf{A}_-^* = \boldsymbol{\sigma}\delta A_-^*$, $\boldsymbol{\delta}\mathbf{A}_+$, and $\boldsymbol{\delta}\mathbf{A}_-$ represent the anti-Stokes and the Stokes waves, respectively ($|\delta A_\pm| \ll |A_0|$); $\psi_+ = (k_z + k_0)z - (\omega + \omega_0)t$ and $\psi_-^* = (k_z - k_0)z - (\omega^* - \omega_0)t$ [3,10]. The scattered vector potential written above, for all ω and k, represents modulational interaction in plasmas. Beating of the Stokes and the anti-Stokes waves with the pump laser leads to the density perturbation $\delta n/n_0$ (plasma wave excitation). It can be estimated after solving the equation of continuity and the Poisson equation together with the z component of the equation of motion, and reads as $\delta\tilde{n} = (e^2k_z^2/2m_e^2\gamma_0^2c^2D_e)(A_0^*\delta A_+ + A_0\delta A_-)$, where $D_e = \omega^2 - \omega_p'^2$, $\omega_p' = \omega_p/\sqrt{\gamma_0}$, $\omega_p = \sqrt{4\pi n_0e^2/m_e}$, $\delta n/n_0 = \delta\tilde{n}e^{i\psi}e^{i\mathbf{k}_\perp\cdot\mathbf{x}_\perp}/2 + \text{c.c.}$, and $\psi \equiv \psi_+ - \psi_0 \equiv \psi_- + \psi_0 = k_zz - \omega t$ [3,10]. It causes an axial component of velocity and momentum $\beta_z = v_z/c \ll 1$ and $p_z \ll p_\perp$, respectively.

Now, we use the above solutions for transverse and longitudinal components of momenta to simplify the radiation reaction term in Eq. (1) and solve the full equation of motion to include the radiation reaction force perturbatively. Writing the CP laser pulse as $\mathbf{A} = \mathbf{A}_\perp(\boldsymbol{x}_\perp, z, t)e^{i\psi_0}/2 + \text{c.c.}$, with its amplitude varying slowly, i.e., $|\partial\mathbf{A}_\perp/\partial t| \ll |\omega_0\mathbf{A}_\perp|$, $|\partial\mathbf{A}_\perp/\partial z| \ll |k_0\mathbf{A}_\perp|$, $|\phi| \ll |\mathbf{A}|$, $\omega_p^2/\gamma\omega_0^2 \ll 1$, and $\gamma = (1 + e^2|\mathbf{A}|^2/m_e^2c^4)^{1/2}$, we get the transverse component of the quiver momentum as

$$\frac{\partial}{\partial t}\left(\mathbf{p}_\perp - \frac{e}{c}\mathbf{A}\right) = -\frac{e\mu\omega_0}{c}\mathbf{A}\gamma|\mathbf{A}|^2(1 - 2\beta_z), \quad (2)$$

where $\mu = 2e^4\omega_0/3m_e^3c^7$, $\beta_z = (\omega/k_zc)\delta\tilde{n}e^{i\mathbf{k}_\perp\cdot\mathbf{x}_\perp}e^{i\psi}/2 + \text{c.c.}$, and we have assumed $\mu\gamma|\mathbf{A}|^2 \ll 1$, which is valid for laser intensities $I_l \le 10^{23}\ \text{W/cm}^2$, for which the influence of the radiation reaction force has to be taken into account. One may also note that we do not consider the effect of radiation reaction on plasma oscillations. This is justified since $|\phi| \ll |A|$ and the radiation reaction effects associated with the plasma wave are negligible in the case of the collinear movement of plasma electrons and the plasma wave. One can solve Eq. (2) for the equilibrium and the scattered vector potentials by substituting $\mathbf{A}$ and expressing the transverse component of the quiver momentum in an analogous manner as the vector potential $\mathbf{A}$, e.g., $\mathbf{p}_\perp = [\mathbf{p}_0e^{i\psi_0} + \mathbf{p}_+e^{i\mathbf{k}_\perp\cdot\mathbf{x}_\perp}e^{i\psi_+} + \mathbf{p}_-^*e^{-i\mathbf{k}_\perp\cdot\mathbf{x}_\perp}e^{-i\psi_-^*}]/2 + \text{c.c.}$, where $\mathbf{p}_+$ and $\mathbf{p}_-$ have similar polarizations as the anti-Stokes and the Stokes modes. The wave equation for the vector potential after the density perturbation $n = n_0 + \delta n$ by the ponderomotive force becomes

$$\nabla^2\mathbf{A} - \frac{1}{c^2}\frac{\partial^2\mathbf{A}}{\partial t^2} = \frac{\omega_p^2}{\gamma c^2}\left(1 + \frac{\delta n}{n_0}\right)\frac{c}{e}\mathbf{p}_\perp. \quad (3)$$

On collecting the terms containing $e^{i\psi_0}$, Eq. (3) yields the dispersion relation for the equilibrium vector potential as $\omega_0^2 = k_0^2c^2 + \omega_p'^2(1 - i\mu|A_0|^2\gamma_0/2)$. It is evident from the dispersion relation that the radiation reaction term causes damping of the pump laser field. This damping can be incorporated either by defining a frequency or a wave number shift in the pump laser [13]. We define a frequency shift of the form $\omega_0 = \omega_{0r} - i\delta\omega_0$, $\delta\omega_0 \ll \omega_{0r}$ with the frequency shift $\delta\omega_0$ being $\delta\omega_0 = \omega_p'^2\varepsilon\gamma_0a_0^2/2\omega_{0r}$, where $\varepsilon = r_e\omega_{0r}/3c$, $r_e = e^2/m_ec^2$ is the classical radius of the electron, and without the loss of generality we have assumed $a_0 = a_0^*$. This frequency shift should be less than the growth rate; otherwise, the growth of the SRS does not occur and the assumption of the locally constant laser field in deriving the growth rates remains no longer valid. Similarly collecting the terms containing $e^{i\psi_\pm}e^{i\mathbf{k}_\perp\cdot\mathbf{x}_\perp}$ in Eq. (3), we get $D_+\delta A_+ = R_+(\delta A_+ + \delta A_-)$ and $D_-\delta A_- = R_-(\delta A_+ + \delta A_-)$, yielding the dispersion relation

$$\left(\frac{R_+}{D_+} + \frac{R_-}{D_-}\right) = 1, \quad (4)$$

where

$$D_\pm = (\omega \pm \omega_0)^2 - \omega_p'^2\left(1 - \frac{i\varepsilon a_0^2\gamma_0\omega_0}{\omega \pm \omega_0}\right) - k_\perp^2c^2 - (k_z \pm k_0)^2c^2,$$

$$R_\pm = \frac{\omega_p^2a_0^2}{4\gamma_0^3}\left[\frac{k_z^2c^2}{D_e}\left(1 \mp i\varepsilon a_0^2\gamma_0 + \frac{2i\varepsilon a_0^2\gamma_0}{k_zc}\frac{\omega\omega_0}{\omega \pm \omega_0}\right) - \left(1 \mp i\varepsilon a_0^2\gamma_0\frac{\omega}{\omega \pm \omega_0} + 4i\varepsilon\gamma_0^3\frac{\omega_0}{\omega \pm \omega_0}\right)\right]. \quad (5)$$

Because of the presence of the radiation reaction term, coupling between the Stokes and the anti-Stokes modes is modified ($R_+ \neq R_-$), and this form of dispersion relation differs from the dispersion relation derived before [1–3,10]. Without the radiation reaction term $\varepsilon = 0$ and

$R_+ = R_- \equiv R$, the dispersion relation assumes the same form as derived before [1–3,10].

For calculating the growth rate of the FRS in a low-density plasma $\omega'_p \ll \omega_{0r}$, one has to take into account both the Stokes and the anti-Stokes waves, as they are both the resonant modes of the plasma [14]. After substituting for the pump laser frequency shift $\delta\omega_0$ and ignoring the finite $k_\perp$ for the FRS, we get $D_\pm = (\omega \pm \omega_{0r})^2 - \omega_p'^2 - (k_z \pm k_0)^2c^2$. On writing $\omega = \omega'_p + i\Gamma_{\rm FRS}$, where $\Gamma_{\rm FRS}$ is the growth rate of the FRS instability, and assuming that both the sidebands (Stokes and anti-Stokes) are resonant, i.e., $D_\pm \approx (\omega \pm \omega_{0r})^2 - \omega_p'^2 - (k_z \pm k_0)^2c^2 = 0$, we have $D_\pm \approx 2i\Gamma_{\rm FRS}(\omega'_p \pm \omega_{0r})$ and $D_e \approx 2i\omega'_p\Gamma_{\rm FRS}$. Substituting these expressions in the dispersion relation and taking $k_z^2c^2 \approx \omega_p'^2$ and $\omega_p'^2 - \omega_{0r}^2 \approx -\omega_{0r}^2$, we obtain, in the weakly coupled regime $\Gamma_{\rm FRS} \ll \omega'_p$, the growth rate which is well approximated by the following expression

$$\Gamma_{\rm FRS} = -\frac{\omega_p^2 \varepsilon a_0^2}{2\omega_{0r}} \pm \frac{\omega_p^2 a_0}{\sqrt{8}\gamma_0^2\omega_{0r}} \cos(\theta/2) \times \sqrt[4]{(1 + 2\varepsilon^2 a_0^2\gamma_0^4)^2 + \varepsilon^2 a_0^4\gamma_0^2\left(\frac{\omega_{0r}}{\omega'_p}\right)^2},$$
$$\theta = \tan^{-1}\left(\frac{-\varepsilon a_0^2\gamma_0(\omega_{0r}/\omega'_p)}{(1 + 2\varepsilon^2 a_0^2\gamma_0^4)}\right). \quad (6)$$

In the case of no radiation reaction force $\varepsilon = 0$, the relativistic growth rate of the FRS instability is same as derived before [3,10]. Two solutions corresponding to $\pm$ signs represent growing and decaying modes, respectively. The decaying mode is damped faster and induces no experimentally detectable signatures in the laser pulse spectrum. The effective growth rate of the FRS instability is $G_{\rm FRS} = \Gamma_{\rm FRS} - \delta\omega_0$. Figure 1 shows the growth rate of the FRS with (upper panel) and without (lower panel) the radiation reaction force. One can immediately notice that the radiation reaction force significantly enhances the growth rate of the FRS at lower plasma densities $\Omega_p \equiv \omega_p/\omega_{0r} \ll 1$ and higher laser amplitude $a_0 \gg 1$, also apparent from Eq. (6). For $\Omega_p \approx 0.02$ and $a_0 = 300$, there is an order of magnitude enhancement in the growth rate. The radiation reaction term also contributes substantially to the growth enhancement of the FRS at higher plasma densities. In this case, the growth rate is also higher since it is directly proportional to the square of the plasma frequency. The strong growth enhancement of the FRS instability is counterintuitive, as the radiation reaction force is generally considered as a damping force similar to collisions in plasmas. This enhancement occurs due to the mixing between the Stokes and the anti-Stokes modes mediated by the radiation reaction force. In the absence of the radiation reaction force, nonlinear currents that drive the Stokes and the anti-Stokes modes have opposite polarizations. Consequently, the phase shift induced by the radiation reaction force—as seen from the expression of $R_\pm$ in Eq. (5)—is opposite for these modes. This results in the interaction between the nonlinear current terms, culminating into phase shift accumulation in Eq. (4). We term the nonlinear mixing of the two modes due to the radiation reaction force as the manifestation of this accumulation of phase shifts, and it leads to the enhanced growth rate of the FRS instability. One can also intuitively imagine this growth enhancement occurring due to the availability of an additional channel of radiation-reaction-force-induced laser energy decay and its efficient utilization by both the Stokes and the anti-Stokes modes.

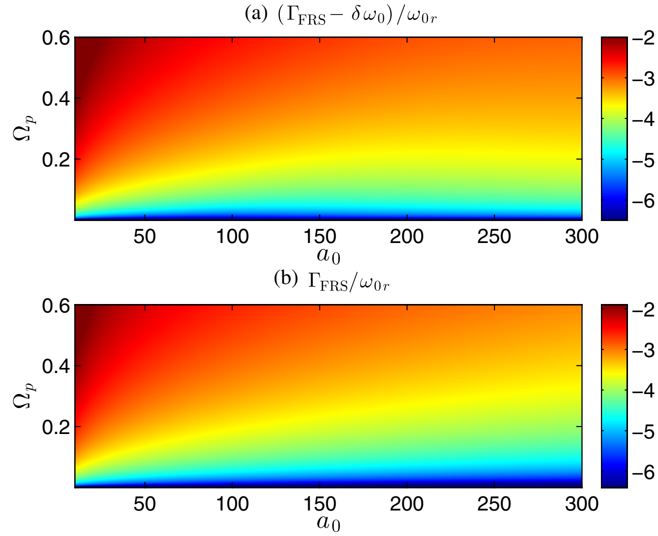

FIG. 1 (color online). Normalized growth rate $(\Gamma_{\rm FRS} - \delta\omega_0)/\omega_{0r}$ of the FRS as a function of the normalized plasma density $\Omega_p \equiv \omega_p/\omega_{0r}$ and normalized pump laser amplitude $a_0 = eA_0/mc^2$ (a) including the radiation reaction force and (b) without the radiation reaction force. The normalized growth rate is plotted on a $\log_{10}$ scale.

Since this growth enhancement depends strongly on the resonant excitation of both the Stokes and the anti-Stokes modes, it is instructive to estimate the conditions under which both modes are excited and also to see if the radiation reaction term enhances the growth of the FRS even when only the Stokes mode is excited in the plasma. Kinematical considerations always allow excitation of the Stokes mode ($D_- = 0$), however, only in a tenuous plasma ($\omega'_p \ll \omega_{0r}$) both the Stokes and the anti-Stokes modes can be simultaneously excited. Assuming that the Stokes mode is excited, one can calculate the frequency mismatch for the anti-Stokes mode which is defined as $\Delta\omega_m = \omega'_p + \omega_{0r} - [\omega_p'^2 + c^2(k'_p + k_0)^2 + D_+]^{1/2}$, and it turns out to be $\Delta\omega_m = -\omega_p'^3/\omega_{0r}^2 + 9\omega_p'^4/4\omega_{0r}^3$. If this frequency mismatch is smaller than the growth rate $\Gamma_{\rm FRS} - \delta\omega_0$ of the FRS instability, then one has to retain both the modes in the dispersion relation while deriving the growth rate of the FRS. Figure 2 depicts the frequency mismatch normalized by the growth rate of the FRS $|\Delta\omega_m/(\Gamma_{\rm FRS} - \delta\omega_0)|$, with a_0 and $\Omega_p \equiv \omega_p/\omega_{0r}$. One

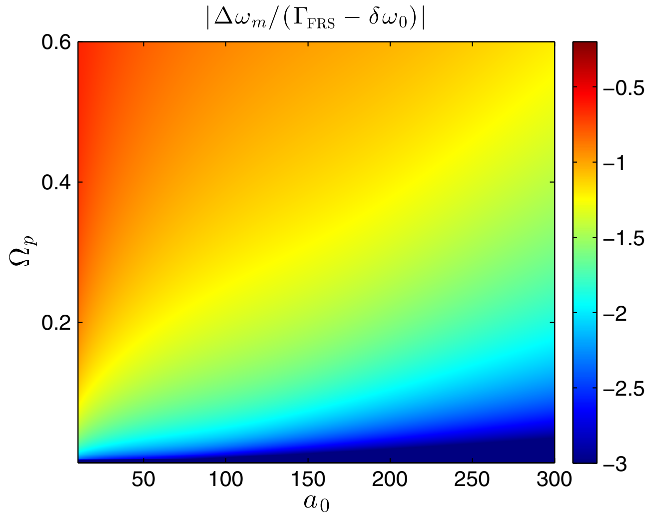

FIG. 2 (color online). Normalized frequency mismatch of the anti-Stokes wave $|\Delta\omega_m/(\Gamma_{\rm FRS}-\delta\omega_0)|$ as a function of normalized plasma frequency $\Omega_p \equiv \omega_p/\omega_{0r}$ and normalized laser amplitude $a_0 = eA_0/mc^2$. The normalized frequency mismatch is again plotted on the $\log_{10}$ scale.

can clearly see that the frequency mismatch for the anti-Stokes mode is usually smaller than the growth rate of the FRS for all values of Ω_p and a_0. This necessitates including both the Stokes and the anti-Stokes modes in the analysis of the FRS. The frequency mismatch is indeed much smaller than the growth rate at lower plasma densities and higher a_0. This is also the parameter regime where strong enhancement to the FRS growth rate occurs. Hence, both the Stokes and the anti-Stokes modes are resonantly excited in the plasma, leading to strong enhancement of the FRS instability due to the radiation reaction force. If one considers only the Stokes mode in the dispersion relation, the growth rate enhancement due to the radiation reaction force is marginal, as the nonlinear mixing of the two Raman sidebands is absent in this case. The growth rate enhancement in this case occurs due to the phase shift caused by the radiation reaction force, which maintains the laser energy transfer to the Stokes mode for a longer time. The BRS is essentially a three-wave decay process, as the anti-Stokes wave is not the resonant mode of the plasma. For the BRS, we have $k_z \simeq 2k_0$ and the instability is always in the strongly coupled regime, i.e., $\Gamma_{\rm BRS} \gg \omega'_p$ (but $\Gamma_{\rm BRS} \ll \omega_{0r}$). One can expand D_e and D_- as $D_e \approx -\Gamma^2_{\rm BRS}$ and $D_- \approx -2i\Gamma_{\rm BRS}\omega_{0r}$, and we get the growth rate of the BRS

$$\Gamma_{\rm BRS} = \frac{\sqrt{3}}{2}\left(\frac{\omega_{0r}}{2\omega_p}\right)^{1/3} \frac{\omega_p a_0^{2/3}}{(1+a_0^2/2)^{1/2}}\left(1+\frac{\varepsilon a_0^2\gamma_0}{3\sqrt{3}}\right). \quad (7)$$

The effective growth rate of the BRS instability is $G_{\rm BRS} = \Gamma_{\rm BRS} - \delta\omega_0$. The radiation reaction term enhances the growth rate of the BRS; however, the enhancement is not strong. Unlike the case of the FRS, no mixing between the anti-Stokes and the Stokes modes is possible in this case due to the absence of the resonant excitation of the former. Again, for $\varepsilon = 0$, one recovers the known growth rate of the BRS [3,10].

To summarize, we have investigated the influence of the leading order term of the Landau-Lifshitz radiation reaction force on the growth of parametric instabilities, namely, the SRS in plasmas. The radiation reaction force strongly enhances the growth of the FRS only when both the Stokes and the anti-Stokes modes are the resonant modes of the plasma. The growths of the FRS—with only the resonant Stokes wave excitation—and the BRS are also enhanced by the inclusion of the radiation reaction force, although the enhancement is a minor one due to the absence of the radiation-reaction-force-induced nonlinear mixing of the anti-Stokes and the Stokes modes. Thus, the radiation reaction force appears to strongly enhance the growth of the SRS involving four-wave decay interaction. These results are important for the ELI Project, as the ultraintense laser pulses are expected to create a dense plasma by strongly ionizing the ambient air and also by producing the electron-positron pairs. The subsequent interaction of this plasma with the laser pulse can lead to the onset of parametric instabilities again—now counterintuitively due to the radiation reaction force—leading to significant change in the frequency spectra and shapes of these extremely intense short laser pulses. Moreover, contrary to the scheme of the nonlinear Compton scattering of a counterpropagating relativistic electron in a strong laser field aiming to discern the signatures of the radiation reaction force on the spectra of high-energy gamma-ray photons [5], enhanced FRS due to the radiation reaction force provides an alternative way to detect the radiation reaction effects on the spectra of low-energy optical photons.

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[13] One can also incorporate the radiation reaction term by appropriately modifying the plasma frequency, which essentially implies change in the laser pump wave vector arising due to its dispersion in the plasma.

[14] We justify retaining both modes in the dispersion relation later by calculating and comparing the frequency mismatch of the anti-Stokes mode with the growth rate of the FRS instability.