

Negative and positive dust grain effect on the modulation instability of an intense laser propagating in a hot magnetoplasma

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Abstract The modulation instability of intense circularly polarized laser beam in hot magnetized dusty plasma is studied. A nonlinear equation describing the interaction of laser with dusty plasma in the quasi-neutral approximation is derived. The effect of negative and positive dust grains on the laser modulation growth rate is studied. It is shown that the existence of positive dust grains instead of ions can substantially improve the modulation growth rate.

Keywords Dusty plasma · Laser · Modulation · Nonlinear interaction · Growth rate · Magnetoactive

Introduction

Dusty plasmas are frequently found in different places of the cosmic environment. They exist in planetary rings, comet comae and tails, and interplanetary and interstellar molecular clouds [1–7]. They can be revealed in the vicinity of aircrafts [6, 7] and in the controlled plasma fusion [8–10]. In some industrial applications of plasma such as plasma processing of materials, the formation of dusty plasma has also been observed [11–14]. They can also be created during laser ablation experiments [15, 16]. In addition, complex dusty plasmas form in the flame of a humble candle, in the zodiacal light, cloud-to-ground lightings, and volcanic eruptions. Recently [17], it has been suggested that the ball lightning is the dusty plasma medium and it is created during oxidation of nanoparticle networks in the normal

lightning strike on soil. Furthermore, dusty plasmas can be produced and investigated in laboratories. Dust grains not only can be intentionally added into the plasma, but can also appear because of different mechanisms in some experiments. The existence of heavy and highly ionized dust grains gives some special and extraordinary properties to the dusty plasma, providing great motivations to investigate it theoretically and experimentally. One of these interests is the study of interaction of laser with dusty plasmas and its related linear and nonlinear effects. These effects include wave dissipation [18], modulation and filamentation instabilities [19–22], linear and nonlinear wave propagation [18, 23–31], parametric instabilities [32], self-focusing [18, 33], etc. Moreover, interaction of laser with dusty plasmas has some important industrial applications. For instance, by interaction of high power lasers with molecular or atomic clusters, during which dusty plasma is created, high-energy electrons can be produced by three processes, i.e., inner ionization, outer ionization, and Coulomb explosion [34–36]. In some experiments, lasers are used in order to study the dynamics of different phenomena in dusty plasmas which some recent experiments about investigation of different exotic phenomena can be found in [37–42].

Here, we focus on the MI of intense lasers in magnetized dusty plasmas. The MI represents a fundamental subject in the theory of nonlinear waves. MI exists due to the interplay between the nonlinearity and dispersion/diffraction effects. The ponderomotive force created by the electromagnetic wave (EMW) stimulates low-frequency perturbations of the electrons density; then, they interact with the primary high-frequency EMW in which the amplitude of the pump wave becomes modulated and the MI of the EMW occurs. The MI of laser beams in plasmas and dielectrics has been the subject of several publications [43–45]. The MI of strong EMWs in plasmas with arbitrary large amplitude was

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studied by Shukla et al. in 1987 [46]. Most of the early publications about MI considered one-dimensional models in which the laser beam was represented as a plane wave [47, 48]. The MI of a laser pulse in the cold nonmagnetized plasma has been considered by several authors [46, 49, 50]. The MI of a linearly polarized laser pulse propagating in the cold magnetized plasma was studied by Jha et al. in 2005 [51]. The MI of the right-hand elliptically laser pulse in cold magnetized plasma has been investigated by Chen et al. in 2011 [52]. Recently, the MI of an intense circularly polarized laser beam in the hot magnetized electron–positron and electron–ion (e–i) plasmas as studied by Sepehri Javan [53, 54]. Our recent work [55] has extended the MI of the circularly polarized laser beam propagating along an external magnetic field in the non-Maxwellian plasma. In this article we study the MI of an intense laser beam in the magnetized hot dusty plasma. In the quasi-neutral approximation and by using a relativistic fluid model, we consider the presence of both negative and positive dust grains and investigate the effect of such grains on the MI. The organization of the paper is as follows. In Sect. 2, the basic assumptions are presented and a nonlinear wave equation is derived for the laser amplitude evolutions. An analytic expression for the growth rate of MI is obtained in Sect. 3. In Sect. 4, a numerical study of the MI of circularly polarized laser beam in the magnetized electron–ion–positive dust–negative dust (e–i–d+–d–) plasma is presented. The concluding remarks are made in Sect. 5.

Deriving a nonlinear wave equation

Let us consider the propagation of a circularly polarized EMW in a hot magnetized four-component plasma which contains electron, ion, and positive and negative dust grains. Each type of plasma particle may have its own specific temperature. To determine the quantities related to the electrons, ions, and positive and negative dust grains, we use indices e, i, d+ and d–, respectively. We take the external magnetic field parallel to the z axis, i.e., $\mathbf{B}_0 = B_0 \hat{\mathbf{e}}_z$. To describe the nonlinear dynamics of the interaction of EMW with the dusty plasma, we define the electric and magnetic fields \mathbf{E} and \mathbf{B} through the vector and scalar potentials \mathbf{A} , φ as:

$$\mathbf{E} = -\frac{1}{c} \frac{\partial \mathbf{A}}{\partial t} - \nabla \varphi, \quad \mathbf{B} = \nabla \times \mathbf{A}, \quad (1)$$

where c is the speed of light.

Using Eq. (1) in Maxwell equations, one can easily obtain:

$$\frac{1}{c^2} \frac{\partial^2 \mathbf{A}}{\partial t^2} - \nabla^2 \mathbf{A} = \frac{4\pi}{c} \mathbf{J}, \quad (2)$$

where $\mathbf{J} = -n_e e \mathbf{v}_e$ is the current density of electrons, e , \mathbf{v}_e and n_e are the density, velocity and charge of the electron, respectively. We ignore the translational velocity of the heavy ions and dust grains. Now, we can write the relativistic fluid momentum equation for electrons as:

$$\frac{\partial \mathbf{p}_e}{\partial t} + (\mathbf{v}_e \cdot \nabla) \mathbf{p}_e = -e \left[\mathbf{E} + \frac{1}{c} \mathbf{v}_e \times (\mathbf{B} + \mathbf{B}_0) \right] - \frac{1}{n_e} \nabla \Pi_e, \quad (3)$$

where \mathbf{p}_e and Π_e are the momentum and pressure of the electron, respectively. Substituting Eq. (1) into Eq. (3) leads to

$$\begin{aligned} \frac{\partial \mathbf{p}_e}{\partial t} + \frac{1}{\gamma_e m_0} (\mathbf{p}_e \cdot \nabla) \mathbf{p}_e &= \frac{e}{c} \frac{\partial \mathbf{A}}{\partial t} + e \nabla \varphi - \frac{e}{\gamma_e m_0 c} \mathbf{p}_e \\ &\times (\nabla \times \mathbf{A}) - \frac{\omega_c}{\gamma_e} \mathbf{p}_e \times \hat{\mathbf{e}}_z - k_B T_e \nabla \ln(n_e), \end{aligned} \quad (4)$$

where T_e is the temperature of electrons, m_0 the electron rest mass, $\gamma_e = \sqrt{1 + p_e^2/m_0^2 c^2}$ the relativistic Lorentz factor, $\omega_c = eB_0/m_0 c$ the electron cyclotron frequency and k_B the Boltzmann constant.

We consider the propagation of circularly polarized wave along the external magnetic field and write the vector potential of this wave as:

$$\mathbf{A} = \frac{1}{2} \tilde{A} (\hat{\mathbf{e}}_x + i\sigma \hat{\mathbf{e}}_y) \exp(-i\omega_0 t + ik_0 z) + c.c., \quad (5)$$

where ω_0, k_0 are the frequency and wave number, respectively. $\sigma = +1, -1$ denotes the right- and left-hand circularly polarized wave, respectively, and also $\tilde{A}(z, t)$ is the slowly varying amplitude that satisfies the following condition:

$$\left| \frac{1}{\omega_0} \frac{\partial \tilde{A}}{\partial t} \right| \ll |\tilde{A}|. \quad (6)$$

Inserting Eq. (5) into Eq. (4), we can find that Eq. (4) is satisfied by [56–58]:

$$\bar{\mathbf{p}}_e = \frac{\bar{\mathbf{A}}}{1 - \sigma\alpha/\gamma_e}, \quad (7)$$

and together with

$$n_e = n_{0e} \exp \left[\frac{e\varphi}{k_B T_e} - \beta_e \left(\gamma_e - 1 - \frac{\sigma\alpha}{2\gamma_e^2} |\bar{\mathbf{p}}_e|^2 \right) \right], \quad (8)$$

where n_{0e} is the unperturbed density of electrons, $\bar{\mathbf{p}}_e = \mathbf{p}_e/m_0 c$ is the normalized electron momentum, $\bar{\mathbf{A}} = e\mathbf{A}/m_0 c^2$ is the normalized vector potential, $\alpha = \omega_c/\omega_0$, $\beta_e = c^2/v_{Te}^2$ and $v_{Te}^2 = k_B T_e/m_0$ is the electron thermal velocity. For weakly relativistic laser intensity, when $|\bar{\mathbf{A}}|^2, |\bar{\mathbf{p}}_e|^2 \ll 1$ and $\gamma_e \approx 1 + \frac{1}{2} |\bar{\mathbf{p}}_e|^2$, we can simplify density of electrons as follows:

$$n_e = n_{0e} \exp \left[\frac{e\varphi}{k_B T_e} - \frac{\beta_e}{2} \frac{|\bar{\mathbf{A}}|^2}{(1 - \sigma\alpha)} \right]. \tag{9}$$

We suppose that the ion and dust grains slow motions are non-relativistic and, by assuming an isothermal equation of state for these heavy particles, obtain the following expressions for number densities:

$$n_i = n_{0i} \exp \left(-\frac{e\varphi}{k_B T_i} \right), \tag{10}$$

$$n_{d+} = n_{0d+} \exp \left(-\frac{z_+ e\varphi}{k_B T_{d+}} \right), \tag{11}$$

$$n_{d-} = n_{0d-} \exp \left(\frac{z_- e\varphi}{k_B T_{d-}} \right), \tag{12}$$

where T_j , n_{0j} , z_+ and z_- are the temperature and unperturbed density of j -type particle, and order of ionization of positive and negative dust grains, respectively.

Expanding Eqs. (9)–(12) and using them in the quasi-neutral condition, i.e., $n_i + z_+ n_{d+} - n_e - z_- n_{d-} = 0$, yield the following result:

$$\Phi = \frac{e\varphi}{k_B T_e} = \frac{\mu\beta_e |\bar{\mathbf{A}}|^2}{2(1 - \sigma\alpha)}, \tag{13}$$

where

$$\mu = \frac{1}{1 + \zeta_i \delta_i^{-1} + z_+^2 \zeta_{d+} \delta_{d+}^{-1} + z_-^2 \zeta_{d-} \delta_{d-}^{-1}}, \tag{14}$$

$$\zeta_j = \frac{n_{0j}}{n_{0e}}, \quad \delta_j = \frac{T_j}{T_e}, \quad j = i, d+, d-. \tag{15}$$

Substituting Eq. (13) into Eq. (9) results in the following expression for the electron density:

$$n_e = n_{0e} \exp \left[\frac{(\mu - 1)\beta_e}{2(1 - \sigma\alpha)} |\bar{\mathbf{A}}|^2 \right]. \tag{16}$$

In physical units, from Eq. (7) we can obtain the following for the velocity of electrons:

$$\mathbf{v}_e = \frac{e}{m_{0e} c} \frac{\mathbf{A}}{\gamma_e - \sigma\alpha}. \tag{17}$$

Also, for electron Lorentz factor, we can approximately write:

$$\gamma_e \approx \sqrt{1 + \frac{|\bar{\mathbf{A}}|^2}{(1 - \sigma\alpha)^2}}. \tag{18}$$

Now, taking Eqs. (16) and (17) into consideration, we derive the nonlinear current density as follows:

$$-\frac{4\pi}{c} \mathbf{J} = \frac{\omega_p^2}{c^2} \frac{\mathbf{A}}{(\gamma_e - \sigma\alpha)} \exp \left[\frac{(\mu - 1)\beta_e}{2(1 - \sigma\alpha)} |\bar{\mathbf{A}}|^2 \right], \tag{19}$$

where $\omega_p = \sqrt{4\pi n_{0e} e^2 / m_0}$ is the electron Langmuir frequency.

In the weakly relativistic regime of laser intensity, we can expand the nonlinear current density of Eq. (19) with respect to the normalized vector potential amplitude and save only the second orders of amplitude. In this case, substituting simplified current density, together with the vector potential in the form of Eq. (5) into Eq. (2) leads to the following equation for the EMW envelope evolutions:

$$\left(\nabla^2 - \frac{1}{c^2} \frac{\partial^2}{\partial t^2} \right) a e^{i(k_0 z - \omega_0 t)} = k_p^2 \left[\frac{1}{1 - \sigma\alpha} - |a|^2 N \right] a e^{i(k_0 z - \omega_0 t)}, \tag{20}$$

where

$$N = \frac{\omega_0^4}{2(\omega_0 - \sigma\omega_c)^4} + \frac{\omega_0^2}{2(\omega_0 - \sigma\omega_c)^2} (1 - \mu)\beta_e, \tag{21}$$

and $a = e\tilde{A}/m_0 c^2$, $k_p = \omega_p/c$ are the normalized amplitude of vector potential, and wave number of the plasma wave, respectively.

For e–i plasma when $\zeta_i = 1$ and $\zeta_{d+} = \zeta_{d-} = 0$, the nonlinear term reduces to

$$N = \frac{\omega_0^4}{2(\omega_0 - \sigma\omega_c)^4} + \frac{\omega_0^2}{2(\omega_0 - \sigma\omega_c)^2} \frac{\beta_e}{1 + T_e/T_i}, \tag{22}$$

which agrees with the results of Sepehri Javan and Nasirzadeh [56].

Derivation of nonlinear dispersion relation and MI

To derive the nonlinear dispersion relation, Eq. (18) is simplified in the new form:

$$\begin{aligned} & \frac{\partial^2 a}{\partial t^2} - c^2 \frac{\partial^2 a}{\partial z^2} - 2i\omega_0 \frac{\partial a}{\partial t} - 2ik_0 c^2 \frac{\partial a}{\partial z} \\ & + \left[-\omega_0^2 + c^2 k_0^2 + \omega_p^2 \left(\frac{1}{1 - \sigma\alpha} - N|a|^2 \right) \right] a \\ & = 0. \end{aligned} \tag{23}$$

In the last term of Eq. (23), the coefficient of a is the nonlinear dispersion relation. In the absence of interaction between EMW and plasma, when amplitude is a real constant ($a = a_0$), we can derive the nonlinear dispersion relation for magnetoplasma with negative and positive dust grains as follows:

$$\begin{aligned} & c^2 k_0^2 - \omega_0^2 \\ & + \omega_p^2 \left[\frac{1}{1 - \sigma\alpha} - a_0^2 \left(\frac{\omega_0^4}{2(\omega_0 - \sigma\omega_c)^4} + \frac{\omega_0^2}{2(\omega_0 - \sigma\omega_c)^2} \frac{(\mu - 1)}{(1 - \sigma\alpha)} \beta_e \right) \right] = 0. \end{aligned} \tag{24}$$

In the linear limit (when $a^2 \rightarrow 0$), Eq. (24) can be reduced to the well-known linear dispersion relation of circularly polarized EMW in the magnetized plasma:

$$k_0 = \frac{\omega_0}{c} \left(1 - \frac{\omega_p^2}{\omega_0(\omega_0 - \sigma\omega_c)} \right)^{1/2}. \quad (25)$$

It is worth mentioning that in the linear approximation, there is no contribution for dust grains on the dispersion Eq. (25) because we have investigated the evolution of high-frequency EMWs where heavy ions and dust particles cannot respond to this high frequency. However, traces of dust particles can be found in the nonlinear dispersion Eq. (24) through bipolar diffusion caused by slow motion of particles under the influence of ponderomotive laser force and thermal collision force.

By considering the condition of slowly varying amplitude (Eq. 6) and assuming that ω_0 and k_0 satisfy the linear dispersion of Eqs. (23), (25) can be modified as:

$$i \left(\frac{\partial a}{\partial t} + v_g \frac{\partial a}{\partial z} \right) + \frac{c^2}{2\omega_0} \frac{\partial^2 a}{\partial z^2} + \frac{\omega_p^2}{2\omega_0} N |a|^2 \quad a = 0 \quad (26)$$

where $v_g = \frac{k_0 c^2}{\omega_0}$ is the group velocity. Using the following dimensionless variables $\tau = \frac{\omega_p^2}{\omega_0^2} t$, $U_g = \frac{\omega_0 v_g}{\omega_p c}$ and $\zeta = \frac{\omega_p}{c} z + U_g \tau$, Eq. (26) can be written as:

$$i \frac{\partial a}{\partial \tau} + \frac{1}{2} \frac{\partial^2 a}{\partial \zeta^2} + D_{\text{NL}} a = 0, \quad (27)$$

where $D_{\text{NL}} = N |a|^2 / 2$. Equation (27) is the well-known nonlinear Schrödinger equation (NLSE). This equation is frequently met in different areas of theoretical physics, especially in nonlinear optics. The NLSE describes the propagation of waves in nonlinear media taking into account both the group velocity dispersion (second term) and the nonlinearity (third term). It is a classical field equation whose important applications are in the propagation of EMWs in nonlinear optical fibers and planar waveguides [59] and to Bose–Einstein condensates confined to highly anisotropic cigar-shaped traps, in the mean-field regime [60]. Additionally, it can be revealed in the studies of small-amplitude gravity waves on the surface of deep zero-viscosity water [59], Langmuir waves in hot plasmas [59], propagation of plane-diffracted wave beams in the focusing areas of the ionosphere [61] and propagation of Davydov's alpha-helix solitons, which are responsible for energy transport along molecular chains [62].

The MI for the right- and left-hand circularly polarized EMW can be obtained using the usual method introduced by Shukla et al. [46]. In this approach, we suppose:

$$a = (a_0 + a_1) \exp(iA\tau), \quad (28)$$

where a_0 is a real constant and $a_0 \gg |a_1|$,

$$A \equiv D_{\text{NL}}(a = a_0) = \frac{1}{2} a_0^2 N. \quad (29)$$

By substituting Eq. (28) into Eq. (27) and linearizing obtained equation with respect to a_1 , we can achieve:

$$i \frac{\partial a_1}{\partial \tau} + \frac{1}{2} \frac{\partial^2 a_1}{\partial \zeta^2} + \frac{1}{2} a_0^2 N (a_1 + a_1^*) = 0. \quad (30)$$

Introducing $a_1 = X + iY$, inserting it into Eq. (30) and separating the real and imaginary parts of this equation yield:

$$\begin{cases} \frac{\partial X}{\partial \tau} + \frac{1}{2} \frac{\partial^2 Y}{\partial \zeta^2} = 0, \\ -\frac{\partial Y}{\partial \tau} + \frac{1}{2} \frac{\partial^2 X}{\partial \zeta^2} + a_0^2 N = 0. \end{cases} \quad (31)$$

We consider the following oscillational form for X and Y :

$$\begin{pmatrix} X \\ Y \end{pmatrix} = \begin{pmatrix} \tilde{X} \\ \tilde{Y} \end{pmatrix} \exp(-i\Omega\tau + iK\zeta), \quad (32)$$

where \tilde{X} and \tilde{Y} are real amplitudes, Ω is the modulation frequency normalized by ω_p^2/ω_0 and K is the modulation wave number normalized by ω_p/c . By substituting Eq. (32) into the set of Eq. (31), we can obtain the following nonlinear dispersion relation of MI:

$$\Omega^2 = -\frac{K^2}{2} \left[a_0^2 N - \frac{K^2}{2} \right]. \quad (33)$$

The temporal growth rate $\Gamma = -i\Omega$ can be extracted from Eq. (33) as below:

$$\Gamma = \frac{K}{\sqrt{2}} \left[a_0^2 N - \frac{K^2}{2} \right]^{1/2}. \quad (34)$$

The maximum growth rate of MI that occurs at $K = K_m = a_0 N^{1/2}$ is

$$\Gamma_{\text{max}} = \frac{a_0^2}{2} N. \quad (35)$$

It may be useful to note that for e-i plasma, when $\zeta_i = 1$ and $\zeta_{d+} = \zeta_{d-} = 0$, Eq. (33) reduces to,

$$\Gamma = \frac{K}{\sqrt{2}} \left[\frac{a_0^2}{2(1 - \sigma\alpha)^2} \left(\frac{\beta_e}{(1 + T_e/T_i)} + \frac{1}{(1 - \sigma\alpha)^2} \right) - \frac{K^2}{2} \right]^{1/2}. \quad (36)$$

Numerical discussions

For numerical studies, in all the investigated cases, we suppose an Nd:YAG laser with frequency $\omega_0 = 1.88 \times 10^{15} \text{ s}^{-1}$ (that corresponds to the laser wave length $\lambda \approx 1 \mu\text{m}$) and $a_0 = 0.271$ (laser intensity $I \approx 10^{17} \text{ W/cm}^2$); also, we

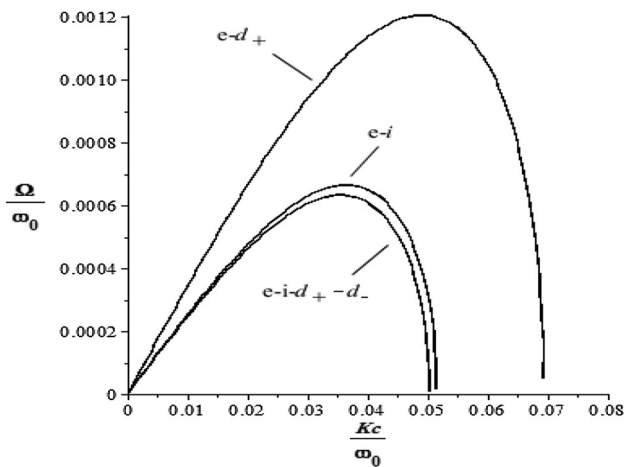


Fig. 1 Variations of the normalized MI growth rate with respect to the normalized modulation frequency for three different cases, e–i plasma with $\eta = \zeta = 0$, e–d+ plasma with $\eta = 10^{-1}$, $\zeta = 0$ and e–i–d+–d– plasma with $\eta = \zeta = 5 \times 10^{-2}$, when $z_+ = z_- = 10$

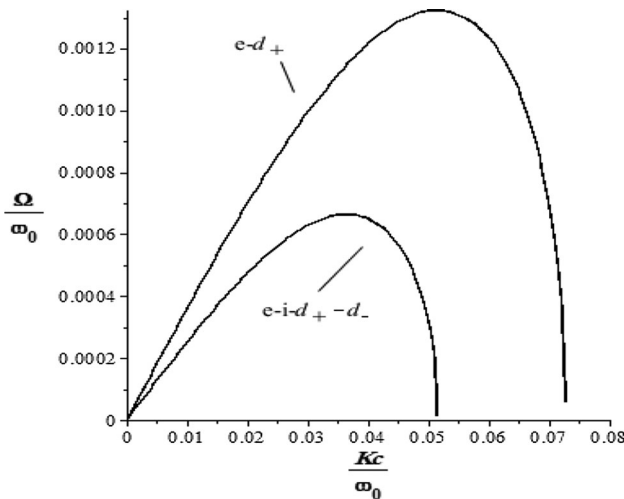


Fig. 2 Variations of the normalized MI growth rate with respect to the normalized modulation frequency for two different cases, e–d+ plasma with $\eta = 10^{-3}$, $\zeta = 0$ and e–i–d+–d– plasma with $\eta = \zeta = 5 \times 10^{-4}$, when $z_+ = z_- = 1000$

consider only the right-hand polarization laser in magnetized medium with $\alpha = 0.2$ and fix the temperature $T_j = 1$ keV for all the plasma components. For more clarification, we introduce two new parameters η and ζ as:

$$n_{0d+} = \eta n_0, \quad n_{0d-} = \zeta n_0, \tag{37}$$

where we supposed $n_{0i} + z_+ n_{0d+} = n_{0e} + z_- n_{0d-} = n_0$ and set $n_0 = 10^{17} \text{ cm}^{-3}$.

Figure 1 shows variations of the normalized modulation growth rate Ω/ω_0 with respect to the normalized modulation wave number Kc/ω_0 , when $z_+ = z_- = 10$. We have three different cases, in which $\eta = \zeta = 0$ corresponds to the e–i

plasma, $\eta = 10^{-1}$, $\zeta = 0$ to the e–d+ plasma and $\eta = \zeta = 5 \times 10^{-2}$ to the e–i–d+–d– plasma. We can see that adding positive dust grains instead of ions substantially increases the modulation growth rate, because localization of positive charges on the dust grains improves the ambipolar potential, which in turn leads to the sharpness of the density profile and consequently to more modulation. In the case of e–i–d+–d– plasma, decreasing the density of electrons and adding equivalent negative dust grains to the plasma results in the decrease in the laser modulation growth rate, because the nonlinear current is created by the motion of the electrons and decrease in the population of electrons leads to the decrease in the nonlinearity of the medium and consequently to the decrease in the growth rate. To investigate the effect of the order of ionization of dust grains on the spot size, in Fig. 2, we choose $z_+ = z_- = 1000$ for two different cases, the e–d+ plasma with $\eta = 10^{-3}$, $\zeta = 0$ and the e–i–d+–d– plasma with $\eta = \zeta = 5 \times 10^{-5}$. We can see that the increase in the ionization order causes a small increase in the MI growth rate. It is worth mentioning that our numerical experiments show that an increase in the temperature causes a decrease in the modulation growth rate. In addition, magnetization of plasma enhances the modulation growth rate for the right-hand polarization and inversely reduces it for the left-hand one. These results are not new and have been investigated earlier [53, 54]; for brevity we do not provide them.

Conclusions

In this paper, we investigated the MI of a weakly relativistic laser propagating along an external magnetic field in the hot plasma containing positive and negative dust grains. The MI growth rate of the circularly polarized laser beam in the dusty plasma was obtained. It was found that adding the positive dust grains to plasma enhances the MI, but existence of the negative dust grains reduces it. Furthermore, the effect of the order of dust grain ionization on the MI was investigated and it was observed that its increase leads to the increase in the MI growth rate.

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