Plasma thermal effect on the relativistic current-filamentation and two-stream instabilities in a hot-beam warm-plasma system

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Based on fully kinetic model using drift-Maxwellian distributions and taking into account the transverse electrostatic field (TEF), it is shown that the current-filamentation instability (CFI) grows unexpectedly with the plasma temperature. The growth is attributed to the decreasing of the TEF as the plasma becomes hot. In the low-temperature plasma regime where the TEF is strong, it is identified that the TEF can dominate over the thermal pressure in suppressing the CFI. Since the TEF originates from the temperature difference between the beam and the plasma, the plasma temperature plays a significant role for the development of the CFI and the quasistatic magnetic fields in a hot-beam warm-plasma system. Particle-in-cell simulations verify the above results.

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I. INTRODUCTION

As a fundamental issue, the current-filamentation instability (CFI) or Weible instability as sometimes called [1–4], in a beam-plasma system has been the subject of research interest for half a century. When a hot electron beam propagates in the background plasma, it will induce a return current in the plasma to keep current neutralization of the beam-plasma system [5], resulting in the CFI. The CFI can lead to filaments formation in the transverse plane perpendicular to the beam propagation direction, and generate strong quasistatic magnetic fields. The filaments formation and magnetic-fields generation are relevant to many problems both in laboratory and astrophysical environments such as the fast ignition scenario of laser inertial confinement fusion [6], laser-plasma interactions [7], laser particle accelerator [8], gamma-ray burst sources [9], active galactic nuclei [10], astrophysical shock acceleration [11], electronic pulsar winds [12], and even the contribution to the dark energy [13]. At the linear stage, it has been pointed out that the collisionless CFI can be suppressed by thermal spread [14–20]. Since the ambient plasma temperature is usually low in reality, it was considered to be irrelevant and negligible. Instead, attentions were usually focused on the beam temperature, which can be as high as thousands of keV. And it was concluded that the CFI is mainly suppressed by the beam temperature. Recently, it was pointed out that a temperature difference between the beam and the plasma will result in a self-consistent electrostatic field [21], which is in the transverse plane perpendicular to the beam drift direction. The transverse electrostatic field (TEF) tends to raise the instability threshold and suppress the CFI. For the basic configuration of a hot-beam warm-plasma system, the temperature difference between the beam and the plasma is large and the self-consistent TEF can be very strong. As the plasma temperature rises, the usual thermal suppression will also increase. However, the electrostatic suppression will decrease since the temperature difference between the beam and the plasma becomes small. So whether the CFI will be amplified or suppressed as the plasma temperature rises still remains a question. Recently it has been shown in the nonrelativistic case that the growth rate of the CFI may be suppressed more than one order of magnitude by TEF [22]. However, due to the difficulty in handling the plasma dispersion function, the present relativistic theoretical models for the CFI have either ignored the TEF [9,18,20], or been based upon water-bag distributions [18,19]. They are only appropriate for the fluid limit and may not treat the plasma thermal effect properly. Therefore, the plasma thermal effect on the CFI is still unclear.

In this paper, we present fully kinetic investigation of the plasma thermal effect on the collisionless relativistic CFI. Our model is based on Vlasov-Maxwell equations and drift-Maxwellian distribution. By using the most accurate Padé approximation Z_{53} [23] for the plasma dispersion function, we can obtain the numerical CFI growth rates at arbitrary plasma electron temperature for the first time. We show both analytically and numerically that the plasma temperature effect is essential for the CFI and the quasistatic magnetic fields generation for a hot-beam warm-plasma system. A low plasma temperature can result in large TEF, which can suppress the CFI more than the thermal inhibition. Therefore, in contrast to the beam thermal effect, the CFI can grow anomalously with the plasma temperature since the TEF becomes weaker as the plasma temperature gets higher. The water-bag distribution model at this point is not proper for the system since the CFI is fully stabilized [18,19]. Although the beam-plasma interaction can also result in two-stream and oblique instabilities [24,25], the long-existent strong quasistatic magnetic fields resulted from the beam-plasma interaction are due to the CFI, regardless of whether it grows faster or slower than the electrostatic instability

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[11,12,25,26]. Therefore, in this paper we mainly focus on the plasma thermal effect on the CFI due to its importance in fast ignition and astrophysics. For comparison, we also include the plasma thermal effect on the two-stream instability (TSI). It is found that the electrostatic instabilities are suppressed or even stabilized as the plasma temperature rises. This is distinctly different from the behavior of the CFI.

II. THEORETICAL MODEL

We assume the plasma is homogeneous, spatially infinite, and unmagnetized, where the ions are fixed to form a charge neutralized background. The beam drifts along the $\hat{e}_z$ direction, and the perturbation wave vector is in the direction $\hat{e}_x$ with $\vec{k}=k_x\hat{e}_x$. The perturbed fields have electromagnetic components with $\vec{E}_{EM}=E_z\hat{e}_z$ and $\vec{B}_{EM}=B_y\hat{e}_y$, and an electrostatic component with $\vec{E}_{ES}=E_x\hat{e}_x$. Assuming the perturbations are in the form of $\exp[i(k_x x-\omega t)]$, we can obtain the well-established linearized kinetic dispersion relation

$$(\varepsilon_{zz}-k_x^2c^2)\varepsilon_{xx}=\varepsilon_{xz}\varepsilon_{zx}, \tag{1}$$

where the dielectric tensor is defined as

$$\varepsilon_{ij}=\omega^2\delta_{ij}+\sum_\alpha \frac{\omega_{p\alpha}^2}{n_\alpha}\int d^3\vec{p}_\alpha \frac{p_{\alpha i}p_{\alpha j}}{\gamma_\alpha^*}\frac{\vec{k}\cdot\partial f_\alpha(\vec{p}_\alpha)/\partial\vec{p}_\alpha}{m_\alpha\gamma_\alpha^*\omega-\vec{k}\cdot\vec{p}_\alpha} + \sum_\alpha \frac{\omega_{p\alpha}^2}{n_\alpha}\int d^3\vec{p}_\alpha \frac{p_{\alpha i}}{\gamma_\alpha^*}\frac{\partial f_\alpha(\vec{p}_\alpha)}{\partial p_{\alpha j}}. \tag{2}$$

Here $\gamma_\alpha^*=(1+\vec{P}_\alpha^2/m_\alpha^2c^2)^{1/2}$ is the Lorentz factor, n_α is the electron density, $\omega_\alpha=(4\pi n_\alpha e^2/m_\alpha)^{1/2}$ is the α-type electron plasma frequency, m_α is the electron mass, and $\alpha=b,p$ represent the beam and the plasma components, respectively. When $\varepsilon_{xz}=0$, $\varepsilon_{zz}-k_x^2c^2=0$ represents the CFI branch and $\varepsilon_{xx}=0$ represents the Langmuir oscillation or ion acoustic branch (where we have ignored the ion oscillation term since it is typically much smaller than the thermal motion term of the ambient electrons). We can see that the CFI mode is coupled to the Langmuir mode or ion acoustic mode due to ε_{xz} and ε_{zx} on the right hand side of Eq. (1). The TEF is related to the electromagnetic part by $E_x=-\varepsilon_{xz}/\varepsilon_{xx}E_z$. If $\varepsilon_{xz}\neq 0$, it will grow exponentially at the linear stage, which tends to inhibit filamentation. For the TSI, we assume the perturbation is in the $\hat{e}_z$ direction with an electrostatic component $\vec{E}_{ES}=E_z\hat{e}_z$. The dispersion relation can be obtained in a similar way:

$$\varepsilon_{zz}=0. \tag{3}$$

Generally, the relativistic beam should be described by the drift Jüttner distribution [25,27]. However, it prohibits an analytical treatment since different momentum components become coupled via the Lorentz factor. In order to obtain analytical results, we use a simplified relativistic drift-Maxwellian distribution [28]

$$f_{0b}=\frac{n_b}{2\pi m\gamma_b^2(T_{tb}T_{lb})^{1/2}}\exp\left[-\frac{P_x^2}{2m\gamma_bT_{tb}}-\frac{(P_z-P_{db})^2}{2m\gamma_b^3T_{lb}}\right], \tag{4}$$

where $\gamma_b=(1+P_{db}^2/m^2c^2)^{1/2}$ and T_{tb}, T_{lb}, and P_{db} correspond to the transverse thermal temperature, the longitudinal thermal temperature, and the drift momentum of the beam electrons, respectively. In the limit of $c\to\infty$, Eq. (4) can be reduced to the nonrelativistic drift-Maxwellian distribution. In the present model, the dense ambient electron component is approximately treated nonrelativistically. Substituting Eq. (4) into Eq. (2), keeping to the first-order expansion of the Lorentz factor γ_b^*, setting $m_b=m_p=m$ and $k_x=k$, we can obtain the dielectric tensor as follows:

$$\varepsilon_{zz}=\omega^2-\frac{\omega_b^2}{\gamma_b^3}-\frac{\omega_b^2}{\gamma_b}\frac{p_{db}^2+m\gamma_b^3T_{lb}}{2m\gamma_bT_{tb}}Z'(\xi_b)-\omega_p^2 - \omega_p^2\frac{p_{dp}^2+mT_{lp}}{2mT_{tp}}Z'(\xi_p), \tag{5}$$

$$\varepsilon_{xx}=\omega^2\left[1-\frac{\omega_b^2}{\gamma_b}\frac{Z'(\xi_b)}{k^2\zeta_b^2}-\omega_p^2\frac{Z'(\xi_p)}{k^2\zeta_p^2}\right], \tag{6}$$

$$\varepsilon_{xz}=\varepsilon_{zx}=-\frac{\omega_b^2}{\gamma_b}\frac{p_{db}}{\sqrt{2m\gamma_bT_{tb}}}\left(Z(\xi_b)+\frac{Z''(\xi_b)}{2}\right) - \omega_p^2\frac{p_{dp}}{\sqrt{2mT_{tp}}}\left(Z(\xi_p)+\frac{Z''(\xi_p)}{2}\right), \tag{7}$$

where $Z(\xi_\alpha)$ is the plasma dispersion function, $\zeta_\alpha=\sqrt{2T_{t\alpha}/\gamma_\alpha m}$, and $\xi_\alpha=\omega/k\zeta_\alpha$. The number of primes on Z denotes the differentiation order. From the third term of Eq. (5) we can see that the beam thermal spread can only suppress the drift-anisotropy [29] of the beam directly. The plasma drift-anisotropy, as we will show later, is mainly suppressed by the TEF since the plasma electron temperature is usually low. The dielectric tensor for the TSI can be obtained similarly by setting $k_z=k$:

$$\varepsilon_{zz}=\omega^2-\frac{\omega_p^2Z'(\chi_p)}{2k^2T_{lp}/m}+\frac{\omega_b^2}{\gamma_b}\frac{\omega^2}{k^2c^2}+\frac{\omega_b^2}{\gamma_b}\frac{\omega^2}{k^2c^2}\frac{Z'(\chi_b)}{2k^2\gamma_bT_{lb}/m} + \frac{\omega_b^2}{\gamma_b}\frac{\omega^2}{k^2c^2}\frac{2\omega Z(\chi_b)}{k\sqrt{2\gamma_bT_{lb}/m}}-\frac{\omega_b^2}{\gamma_b}\frac{Z'(\chi_b)}{2k^2\gamma_bT_{lb}/m}, \tag{8}$$

where $\chi_\alpha=(\gamma_\alpha m\omega-kp_{d\alpha})/\sqrt{2k^2\gamma_\alpha^3mT_{l\alpha}}$. In the limit of $|\chi_\alpha|\gg 1$, Eq. (8) can return to the fluid limit form:

$$\varepsilon_{zz}=\omega^2\left(1-\frac{\omega_b^2/\gamma_b^3}{(\omega-kv_{db})^2}-\frac{\omega_p^2}{(\omega-kv_{dp})^2}\right), \tag{9}$$

where $v_{d\alpha}$ is the drift velocity of the α type electron.

III. PLASMA THERMAL EFFECT ON THE CFI AND THE TSI

If the beam-plasma system is weak in drift-anisotropy, the system falls in the kinetic region [29]. In this case we have

$|\xi_\alpha|\ll 1$, i.e., $\delta \ll k_x v_{t\alpha}$ if we assume $\omega = i\delta$. The plasma dispersion function can be expanded as $Z(\xi_\alpha)=-2\xi_\alpha+\frac{4}{3}\xi_\alpha^3-\frac{8}{15}\xi_\alpha^5\cdots+i\sqrt{\pi}\exp(-\xi_\alpha^2)$. The CFI is coupled to the ion acoustic branch, and the growth rate is

$$\delta \simeq \frac{\omega_b^4 v_{db}^2\left(\sqrt{\frac{n_b}{n_p}}+\sqrt{\frac{n_p}{n_b}}\right)^2}{k^2 T_{tp} T_{tb} BD/m^2}+\frac{\omega_b^2(\gamma_b-1/\gamma_b^3)-k^2c^2}{D}, \tag{10}$$

where the current neutralization condition $v_{dp}=-v_{db}n_b/n_p$ [5] is used, and $B=\omega_b^2 m/k^2T_{tb}+\omega_p^2 m/k^2T_{tp}$, $D=\frac{\sqrt{\pi}}{\sqrt{2}k}\left[\omega_b^2\frac{v_{db}^2+\gamma_b T_{lb}/m}{\gamma_b(T_{tb}/\gamma_b m)^{1.5}}+\omega_p^2\frac{v_{dp}^2+T_{lp}/m}{(T_{tp}/m)^{1.5}}\right]$. Assuming both the beam and the background electrons are isotropic with $T_\alpha=T_{t\alpha}=T_{l\alpha}$, and the beam is much hotter than the plasma, i.e., $T_b \gg T_p$, Eq. (10) can be simplified to

$$\delta \sim \sqrt{\frac{2}{\pi}}k\frac{\omega_b^2 v_{db}^2 m/T_b+\omega_b^2(\gamma_b-1/\gamma_b^3)-k^2c^2}{\omega_p^2(v_{dp}^2+T_p/m)}\left(\frac{T_p}{m}\right)^{3/2}. \tag{11}$$

The drift term of the plasma electron v_{dp} does not appear in the numerator of Eq. (11) due to inclusion of the TEF effect. Thus, it is the electrostatic field that mainly suppresses the plasma's contribution to the CFI. The maximal growth rate is obtained as

$$\delta_m \sim \frac{2\sqrt{6}}{9\sqrt{\pi}}\frac{[\omega_b^2 v_{db}^2 m/T_b+\omega_b^2(\gamma_b-1/\gamma_b^3)]^{3/2}}{\omega_p^2(v_{dp}^2+T_p/m)c}\left(\frac{T_p}{m}\right)^{3/2}, \tag{12}$$

which shows that the instability growth rate is proportional to the temperature of the background plasma electrons. Thus, the CFI in fact grows with the plasma temperature. Compared with the beam electron temperature effect which would suppress the CFI growth rate [16–20], the growth of the CFI with the plasma temperature is anomalous.

Due to the complexity of the plasma dispersion function, there are no general analytic solutions to Eq. (1) for all cases. Therefore we mainly focus on numerical solutions next. Figure 1 shows the maximum growth rate of the CFI numerically obtained from Eq. (1), where we have kept to the second order expansion of the Lorentz factor. In order to obtain the CFI growth rates at arbitrary plasma electron temperature, the most accurate Padé approximation Z_{53} [23] for the plasma dispersion function is used. It is seen from Fig. 1 that in the case without the TEF ($\varepsilon_{xz}=0$), the maximum growth rate is suppressed by the plasma thermal spread due to the usual thermal suppression, which is in agreement with the fluid theory. However, when the TEF is included, the growth rate turned to increase with the plasma temperature, which is in agreement with our analytic prediction given above. Moreover, the growth rate is drastically reduced compared to the case without the TEF. The reduction can be more than one order of magnitude in the low-temperature region. The above results can be understood physically as follows. Generally, the forward beam current and the returning plasma current will pinch at different rates in the resultant magnetic fields since the beam and the plasma usually have different thermal spread, resulting in space charge separation and the TEF [21]. The electrostatic fields tend to balance the magnetic pinch force and inhibit further filamentation. As the plasma temperature grows, the thermal suppression becomes stronger but the electrostatic suppression becomes weaker. Since the electrostatic suppression dominates over the thermal suppression, the CFI is enhanced as the plasma temperature grows. This can be identified more clearly by the comparison of the CFI growth rate with and without the TEF in Fig. 2, where the growth rate is plotted as a function of the unstable wave numbers. It can be seen from Figs. 2(a) and 2(b) that for a typical relativistic beam-plasma system, the CFI growth rate is reduced by a factor of 3 when the beam

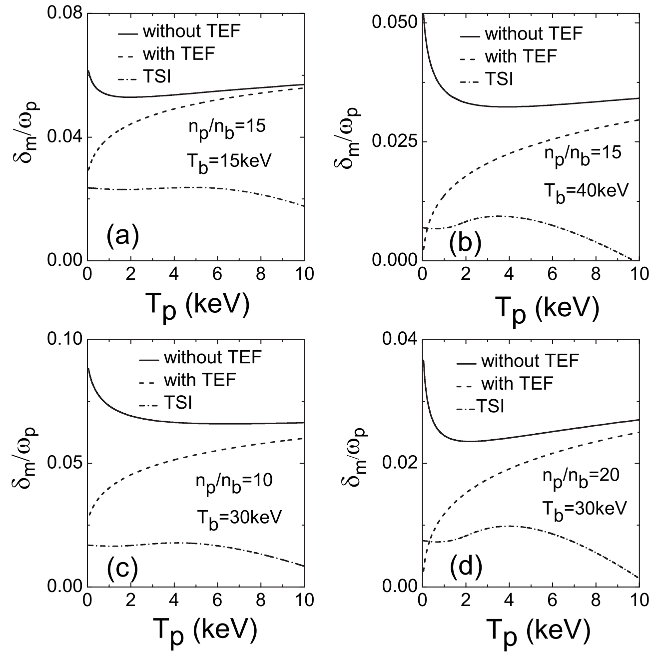

FIG. 1. The maximum linear growth rate of the CFI and TSI versus the plasma temperature T_p obtained by solving Eq. (1), where T_p is normalized to 1 keV. The drift velocity of the beam is set to be $v_{db}=0.9c$, where c is the light speed in the vacuum.

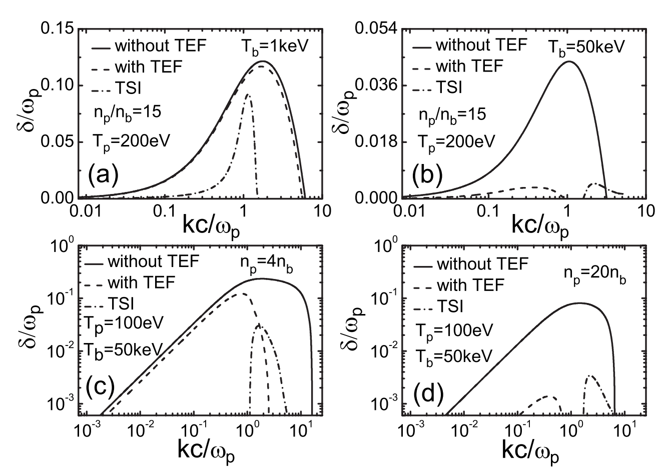

FIG. 2. The typical linear growth rate of the CFI and TSI versus the wave number k obtained by solving Eq. (1). The drift velocity of the beam is set to be $v_{db}=0.9c$.

temperature is changed from 1 to 50 keV in the case without the TEF. However, the growth rate is reduced nearly as much as two orders of magnitude if the TEF is taken into consideration. Thus, it is the TEF instead of the thermal spread [16–20] that mainly suppresses the CFI of a hot-beam warm-plasma system. Besides the growth rate, the unstable range is also suppressed.

For the CFI to be enhanced by the plasma temperature, the beam should be hotter than the plasma and the beam-plasma system should fall in the kinetic domain. We now explore if the TEF effect is significant in the hydrodynamic region where $|\xi_\alpha| \gg 1$. In this case the plasma dispersion function could be expanded as $Z(\xi_\alpha) = -\xi_\alpha^{-1} - (2\xi_\alpha^3)^{-1} + \cdots$ In the nonrelativistic case, to the first order ξ_α^{-1} in the asymptotic expansion, ε_{xz} is zero so the TEF effect is only higher order correction to the CFI [21]. However, in the present relativistic case, the beam electron mass is changed from m to $\gamma_b m$. Thus, the first order correction is nonzero, and the CFI is coupled to the electron Langmuir oscillation mode. In the assumption of current neutralization, the growth rate can be obtained if we expand the plasma dispersion function to the first order,

$$\delta = \left(\delta_0^2 - \frac{A^2}{\omega_b^2/\gamma_b + \omega_p^2} \frac{1}{\omega_b^2/\gamma_b^3 + \omega_p^2 + k^2c^2} \right)^{1/2}, \qquad (13)$$

where $\delta_0 = k[\frac{(v_{db}^2 + v_{tb}^2)\omega_b^2/\gamma_b + (v_{db}^2 + v_{tb}^2)\omega_p^2}{\omega_b^2/\gamma_b^3 + \omega_p^2 + k^2c^2}]^{1/2}$ is the CFI growth rate without the TEF, $A = (1/\gamma_b - 1) k v_{db} \omega_b^2/\gamma_b$ is the term caused by TEF. Since $\omega_b^2/\gamma_b \ll (\omega_b^2/\gamma_b + \omega_p^2)$, the TEF correction is usually small in the hydrodynamic region. It is seen from Figs. 2(c) and 2(d) that the electrostatic suppression becomes significant by increasing the density of the plasma. A careful check finds that in the low plasma density case, the beam-plasma system mainly falls in the hydrodynamic region. As the density of the plasma increases, the CFI growth rate gets smaller, thus the beam-plasma system falls in the hybrid domain with $|\xi_b| < 1$ and $|\xi_p| > 1$, or falls in the kinetic domain with $|\xi_\alpha| \ll 1$. Since the instability growth time δ^{-1} is much longer than the thermal response time $(kv_\alpha)^{-1}$, the thermal motions become important. Therefore, the coherent electrostatic suppression becomes significant. It is also seen from Fig. 2(c) that the growth rate and the unstable range are greatly suppressed in the short wavelength region for the low plasma density case when TEF is included, since the short wavelength region has fallen in the hybrid domain or the kinetic domain. For the high plasma density case, all the unstable waves fall in the hybrid or kinetic domain, therefore the TEF is so strong that the maximum growth rate is reduced by nearly two orders of magnitude. This will make the magnetic fields grow more slowly.

For comparison, we have also plotted in Fig. 1 the maximum growth rates of the electrostatic TSI. It can be seen from Fig. 1 that the TSI is generally suppressed by the plasma thermal spread, especially in the high temperature region, where we can see from case (b) that the TSI is even fully stabilized by the plasma thermal spread. Compared to the case of TSI, the CFI's growth with the plasma temperature is anomalous. Moreover, Figs. 1 and 2 show that the growth rates of the CFI in most cases are larger than those of the electrostatic TSI. Therefore, the anomalous growth of the CFI with the plasma temperature is an important property for the hot-beam warm-plasma system. It is seen both from Figs. 1(b) and 1(d) that in the region where the plasma temperature is ultra low, the TSI can dominate over the CFI if the TEF is included. However, the TSI growth rates are still smaller than the CFI growth rates without the TEF effect. Thus, it is the TEF suppression rather than the thermal suppression that prevents the CFI from becoming the dominant mode in this region. This can be seen more clearly from Fig. 2, where all the unstable waves are shown. By comparing case (a) with case (b) and case (c) with (d), respectively, we can find it is due to the TEF that the CFI is suppressed to have a lower growth rate and a smaller unstable region than the electrostatic TSI.

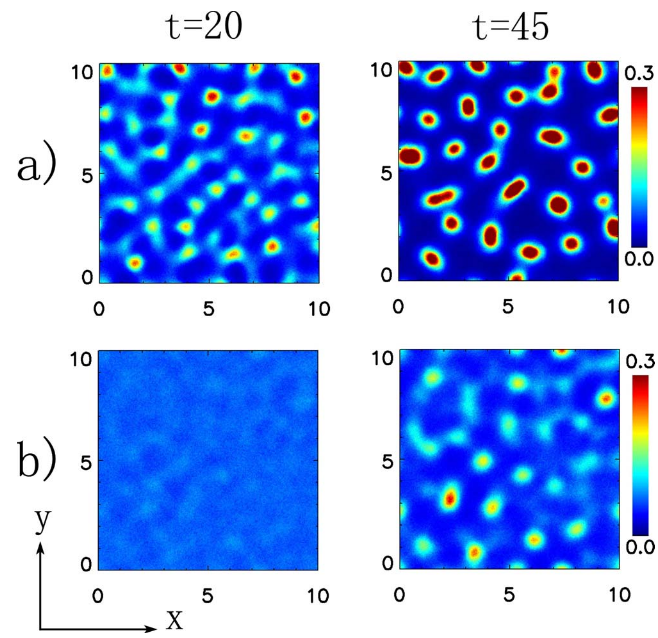

FIG. 3. (Color online) Snapshots of the evolution of the beam filament densities during the linear and the nonlinear stage for hot plasma $T_p = 9$ keV (case a) and cold plasma $T_p = 100$ eV (case b). Here x and y are normalized by c/ω_p.

To compare with our analytic and numerical results, we perform two-dimensional (2D) particle-in-cell simulations. The beam is assumed to propagate along $\hat{z}$ with the initial distribution as Eq. (4). The plasma ions are fixed with the density $n_i = n_b + n_p$, where $n_p = 15 n_b$. The simulation domain is $X \times Y = (20.48\lambda \times 20.48\lambda)$ with 512 cells in each direction, where $\lambda = 2\pi c/\omega_p$. There are 60 particles per cell per species for all the simulations. The drift velocities are set to $v_{db} = 0.9c$ and $v_{dp} = -0.06c$. Figure 3 shows snapshots of the structure of the beam filaments at the time $t = 20(2\pi/\omega_{pe})$ and $t = 45(2\pi/\omega_{pe})$ for high-temperature plasma case with $T_p = 9$ keV and low-temperature plasma case with $T_p = 100$ eV. In Fig. 3(a) for the hot plasma case, the filaments are clearly seen at $t = 20$ (the maximum electron density scale is set to $0.3 n_p$). However, for the low temperature plasma case, the CFI grows much slower that the filaments are not clearly seen until $t = 45$. The electron density in the filaments in case (a) is twice larger than that of case (b). We also plot

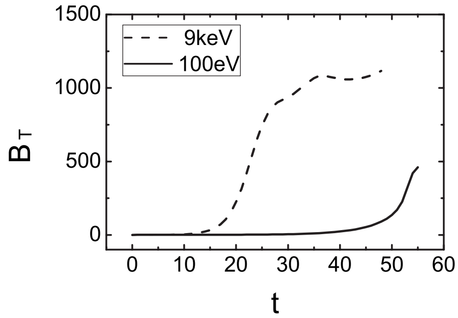

FIG. 4. Time evolution of the transverse magnetic field energies. The time is normalized to $2\pi/\omega_p$.

the evolution of the transverse magnetic field energies for the above simulations in Fig. 4. It is seen that the magnetic field grows much slower at the linear stage for $T_p=100$ eV, which is in good agreement with our analytic prediction. Later at the early quasilinear stage, the fields begin to saturate at a much lower level for the 100 eV case compared to the 9 keV case. Thus, the anomalous thermal effect should have an important effect on the nonlinear evolution of the CFI and the quasistatic magnetic fields.

IV. CONCLUSION

In conclusion, we have presented a kinetic calculation showing that the current-filamentation instability and the quasistatic magnetic fields grow faster while the two-stream instability grows slower as the plasma temperature increases in a hot-beam warm-cold plasma system. It is also identified that the self-consistent transverse electric fields can dominate over the plasma thermal spread in suppressing the CFI. Detailed 2D particle-in-cell simulations confirm the analytic predictions.

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[1] E. S. Weibel, Phys. Rev. Lett. **2**, 83 (1959).
[2] F. Califano, F. Pegoraro, and S. V. Bulanov, Phys. Rev. E **56**, 963 (1997).
[3] T. Taguchi, T. M. Antonsen, C. S. Liu, and K. Mima, Phys. Rev. Lett. **86**, 5055 (2001).
[4] O. Polomarov, I. Kaganovich, and G. Shvets, Phys. Rev. Lett. **101**, 175001 (2008).
[5] A. R. Bell, J. R. Davies, S. Guerin, and H. Ruhl, Plasma Phys. Contr. Fusion **39**, 653 (1997).
[6] M. Tabak, J. Hammer, M. E. Glinsky, W. L. Kruer, S. C. Wilks, J. Woodworth, E. M. Campbell, M. D. Perry, and R. J. Mason, Phys. Plasmas **1**, 1626 (1994).
[7] R. Jung, J. Osterholz, K. Löwenbrück, S. Kiselev, G. Pretzler, A. Pukhov, O. Willi, S. Kar, M. Borghesi, W. Nazarov, S. Karsch, R. Clarke, and D. Neely, Phys. Rev. Lett. **94**, 195001 (2005).
[8] R. Bingham, Nature (London) **368**, 496 (1994).
[9] M. V. Medvedev and A. Loeb, Astrophys. J. **526**, 697 (1999).
[10] J. A. Zensus, Annu. Rev. Astron. Astrophys. **35**, 607 (1997).
[11] K. I. Nishikawa, P. Hardee, G. Richardson, R. Preece, H. Sol, and G. J. Fishman, Astrophys. J. **622**, 927 (2005).
[12] Y. Kazimura, J. I. Sakai, T. Neubert, and S. V. Bulanov, Astrophys. J. **498**, L183 (1998).
[13] I. Contopoulos and S. Basilakos, Astron. Astrophys. **471**, 59 (2007).
[14] M. Bornatici and K. F. Lee, Phys. Fluids **13**, 3007 (1970).
[15] R. Davidson, D. A. Hammer, I. Haber, and C. E. Wagner, Phys. Fluids **15**, 317 (1972).
[16] R. Lee and M. Lampe, Phys. Rev. Lett. **31**, 1390 (1973).
[17] K. Molvig, Phys. Rev. Lett. **35**, 1504 (1975).
[18] L. O. Silva, R. A. Fonseca, J. W. Tonge, W. B. Mori, and J. M. Dawson, Phys. Plasmas **9**, 2458 (2002).
[19] A. Bret, M. C. Firpo, and C. Deutsch, Phys. Rev. E **72**, 016403 (2005).
[20] A. Karmakar, N. Kumar, G. Shvets, O. Polomarov, and A. Pukhov, Phys. Rev. Lett. **101**, 255001 (2008).
[21] M. Tzoufras, C. Ren, F. S. Tsung, J. W. Tonge, W. B. Mori, M. Fiore, R. A. Fonseca, and L. O. Silva, Phys. Rev. Lett. **96**, 105002 (2006).
[22] B. Hao, Z.-M. Sheng, and J. Zhang, Phys. Plasmas **15**, 082112 (2008).
[23] P. Martín, G. Donoso, and J. Z. Cristi, J. Math. Phys. **21**, 280 (1980).
[24] S. A. Bludman, K. M. Watson, and M. N. Rosenbluth, Phys. Fluids **3**, 747 (1960).
[25] A. Bret, L. Gremillet, D. Bénisti, and E. Lefebvre, Phys. Rev. Lett. **100**, 205008 (2008).
[26] X. Kong, J. Park, C. Ren, Z.-M. Sheng, and J. W. Tong, Phys. Plasmas **16**, 032107 (2009).
[27] T. P. Wright and G. R. Hadley, Phys. Rev. A **12**, 686 (1975).
[28] K. M. Watson, S. A. Bludman, and M. N. Rosenbluth, Phys. Fluids **3**, 741 (1960).
[29] B. Hao, Z.-M. Sheng, C. Ren, and J. Zhang, Phys. Rev. E **79**, 046409 (2009).