discussions with P. A. Wolff, C. K. N. Patel, C. H. Henry, and R. Dingle. They thank D. G. Thomas and J. K. Galt for comments on the manuscript, and R. Epworth and R. P. Reeves for their excellent technical assistance.

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[12]This observation of large gain, g_ν, in a spectral region where the observed low-intensity absorption, $\alpha_\nu^{\rm low}$, is small, i.e., $g_\nu > |\alpha_\nu^{\rm low}|$, is not unique in semiconductor physics to GaP. Indeed, stimulated emission and laser action with $g_\nu > |\alpha_\nu^{\rm low}|$ occur in low-absorption regions well below the band gap in many semiconductors. GaAs {Ref. 4 and N. G. Basov, O. V. Bogdankevich, V. A. Gancharov, B. M. Lavrushin, and V. Uy. Sudzilovskii, Dokl. Akad. Nauk SSSR 168, 1283 (1966) [Sov. Phys. Dokl. 11, 522 (1966)]} is an excellent example where $g_\nu > |\alpha_\nu^{\rm low}|$. In addition to inferences drawn in the text, the observation $g_\nu > |\alpha_\nu^{\rm low}|$ implies rapid depopulation of the terminal state for the radiative transition; i.e., this transition does not correspond to a simple two-level system.

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Parametric Excitation of Plasma Instabilities in Semiconductors

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The parametric excitation of density waves in semiconducting plasmas is considered. A new nonlinear mechanism for the direct conversion of photons into plasmons is presented, resulting from the nonparabolic momentum-energy relations for the single electron. Applications to InSb yield threshold intensities typically in the range 10^4–10^6 W/cm^2.

The parametric excitation of density waves in gaseous and solid-state plasmas has been of considerable interest recently.[1] In this Letter we discuss a new nonlinear mechanism for the direct conversion of photons into plasmons which, to the best of our knowledge, has not been considered before. It results from the nonparabolic momentum-energy relation for the single electrons in semiconductors and becomes important in the case of an intense field, e.g., laser radiation. Next we show that for InSb under typical conditions, modest radiation intensities can give rise to growth rates of about 10% per period. Thus, a strong instability can be generated with current state-of-the-art laser fields. The large-amplitude plasma waves obtained in the instability can be monitored in light-scattering experiments where the strength of the anti-Stokes line is proportional to the induced plasma-wave amplitude. The efficient conversion of radiant energy into plasmons provides us with a useful tool for studying large-amplitude plasma oscillations as well as such processes as stimulated Raman scattering.

The starting point of our discussion is based upon the effective single-particle Hamiltonian[2]

$$H=[(m^*c^{*2})^2+(c^*p)^2]^{1/2},$$

where $c^{*2}=E_g/2m^*$. Here E_g is the gap energy and m^* is the electronic effective mass at the bottom of the conduction band. It follows that

$$\vec{v}=c^*\vec{p}[(m^*c^*)^2+p^2]^{-1/2}.$$

In our case the electric field can be treated as homogeneous and the ac Lorentz force may be neglected, since $v/c\ll 1$. The dynamics of the conduction electrons are governed by the Vlasov equation,

$$\frac{\partial f}{\partial t}+\vec{v}\cdot\frac{\partial f}{\partial\vec{x}}\cdot e\frac{\partial\varphi}{\partial\vec{x}}\cdot\frac{\partial f}{\partial\vec{p}}-e\vec{E}(t)\cdot\frac{\partial f}{\partial\vec{p}}=0, \tag{1}$$

and Poisson's equation,

$$\epsilon_L\nabla^2\varphi=4\pi n_0 e[-1+\int f\,d\vec{p}]. \tag{2}$$

Here $f(\vec{x},\vec{p},t)$ represents the electronic distribution function in phase space, and $\vec{E}(t)$ is the electric field strength of the electromagnetic wave. The lattice dielectric constant has been denoted by ϵ_L and the equilibrium density by n_0.

The instability of a plasma wave is characterized by two parameters—a threshold field and a growth rate. To calculate them we solve Eqs. (1) and (2) for f in an arbitrary external field. It will be assumed that the plasma-wave excitations may be treated linearly, which is adequate for the calculation of the threshold field. Our solution will thus predict an exponential growth of the plasma waves in time, but not the saturation amplitude of these oscillations.

We now express the distribution function as $f=f_0+F$, where f_0 is the response of the electrons to the homogeneous field, and F is the spatially dependent part describing plasma-wave excitations. The function f_0 obeys the kinetic equation

$$\partial f_0/\partial t-e\vec{E}(t)\cdot\partial f_0/\partial\vec{p}=0. \tag{3}$$

It is convenient to transform to an oscillatory frame of reference by letting

$$\vec{\eta}=\vec{p}+\int_{-\infty}^{t}\vec{E}(t')e^{\epsilon t'}dt',\quad \vec{\xi}=\vec{x}+(e/m^*)\int_{-\infty}^{t}dt'\int_{-\infty}^{t'}dt''\,\vec{E}(t'')e^{\epsilon t''},$$

with $\epsilon\to 0^+$, and $\tau=t$. The zeroth-order distribution function, f_0, could be chosen to be an arbitrary function of $\vec{\eta}$ although collisions would cause it to become an equilibrium (e.g., Fermi-Dirac) distribution function. The zeroth-order response of the electrons, apart from thermal motion, is to oscillate coherently with the electric field. Equation (1) becomes

$$\frac{\partial F}{\partial\tau}+\left(\vec{u}+\frac{\vec{\eta}}{m^*}\right)\cdot\frac{\partial F}{\partial\vec{\xi}}+e\frac{\partial\varphi}{\partial\vec{\xi}}\cdot\frac{\partial f_0}{\partial\vec{\eta}}=0, \tag{4}$$

and Eq. (2) is

$$\epsilon_L\nabla_\xi^2\varphi=4\pi\eta_0 e\int F\,d\vec{\eta}. \tag{5}$$

We have introduced a quantity $\vec{u}$ defined by

$$\vec{u}=c^*[\vec{\eta}-e\int_{-\infty}^{\tau}\vec{E}e^{\epsilon t}dt]\{[(m^*c^*)^2+(\vec{\eta}-e\int_{-\infty}^{\tau}\vec{E}e^{\epsilon t}dt)^2]^{-1/2}-1/m^*c^*\}. \tag{6}$$

It represents the nonlinearity in the plasma arising from the "relativistic" velocity-momentum relation (thus as $c^*\to\infty$, $\vec{u}\to 0$). It is worthwhile mentioning that our nonlinearity, for a one-component plasma, is similar to the nonlinearity due to the convective term that would appear in a hydrodynamical description. This is seen by simply taking the velocity moment of Eq. (4). If c^* were infinite, the nonlinearity would be a true relativistic effect induced by the ac magnetic field (i.e., a finite wave-number effect) and would be of the order $\langle v\rangle/c$. In our case the nonlinearity is of order $\langle v\rangle/c^*$ which is roughly two orders of magnitude larger than $\langle v\rangle/c$. Here $\langle v\rangle$ is the velocity induced by the driving electric field.

The solution to Eqs. (4)–(6) can be obtained in powers of the field strength E. We shall retain terms through order E^2. Our attention will be limited to running-wave solutions, i.e., to waves propagating as $\exp i\vec{k}\cdot\vec{\xi}$. We now take $\vec{E}(t)$ to be $\vec{E}=\vec{E}_1\cos(\omega_1 t-\eta_1)+\vec{E}_2\cos(\omega_2 t-\eta_2)$ with $\vec{E}_1\parallel\vec{E}_2$. Then Eqs. (4)–(6) yield

$$-i\frac{\partial F}{\partial\tau}+\{\alpha+\textstyle\sum_i[\beta_i\sin(\omega_i\tau-\eta_i)+\gamma_i\cos 2(\omega_i\tau-\eta_i)]+\delta[\cos((\omega_1-\omega_2)\tau-\eta_1+\eta_2)-\cos((\omega_1+\omega_2)\tau-\eta_1-\eta_2)]\}F$$
$$-\frac{m^*\omega p^2}{k^2}\vec{k}\cdot\frac{\partial f_0}{\partial\vec{\eta}}\int F\,d\vec{\eta}=0, \tag{7}$$

where $\omega_p^2 = 4\pi n_0 e^2/m^*\epsilon_L$ and letting

$$\alpha = c^*\{\vec{k}\cdot\eta/\Delta - (1/2\Delta^3)\textstyle\sum_i[\vec{k}\cdot\vec{\lambda}_i\vec{\eta}\cdot\vec{\lambda}_i + \frac{1}{2}\lambda_i^2\vec{k}\cdot\vec{\eta} - \frac{3}{2}\vec{k}\cdot\vec{\eta}(\vec{\eta}\cdot\vec{\lambda}_i/\Delta)^2]\}, \tag{8a}$$

$$\beta_i = (c^*/\Delta^3)[\vec{k}\cdot\vec{\eta}\vec{\eta}\cdot\vec{\lambda}_i - \vec{k}\cdot\vec{\lambda}_i(\Delta^2 - \Delta^3/m^*c^*)], \tag{8b}$$

$$\gamma_i = (c^*/4\Delta^3)[\vec{k}\cdot\vec{\eta}\lambda_i^2 + 2\vec{k}\cdot\vec{\lambda}_i\vec{\eta}\cdot\vec{\lambda}_i - (3/\Delta^2)\vec{k}\cdot\vec{\eta}(\vec{\eta}\cdot\vec{\lambda}_i)^2], \tag{8c}$$

$$\delta = -(c^*/2\Delta^3)[\vec{k}\cdot\vec{\eta}\vec{\lambda}_1\cdot\vec{\lambda}_2 + \vec{k}\cdot\vec{\lambda}_1\vec{\eta}\cdot\vec{\lambda}_2 + \vec{k}\cdot\vec{\lambda}_2\vec{\eta}\cdot\vec{\lambda}_1 - (3/\Delta^2)\vec{k}\cdot\vec{\eta}\vec{\lambda}_1\cdot\vec{\eta}\vec{\lambda}_2\cdot\vec{\eta}], \tag{8d}$$

where $\vec{\lambda}_i = e\vec{E}_i/\omega_i$, $i = 1, 2$; $\Delta = [\eta^2 + (m^*c^*)^2]^{1/2}$.

The only modes of excitation which can propagate in our system are the plasma waves. An instability can occur when the radiation field resonantly couples to two such waves. In our analysis we therefore discard all nonresonant terms. We take F to be of the form

$$F = A_1(\vec{\eta})e^{i(\omega - i\mu)\tau} + A_2(\vec{\eta})e^{-i(\omega + i\mu)\tau} + \text{c.c.}, \tag{9}$$

where ω, the plasmon frequency, is near ω_p. There are two possibilities for an instability to occur to this order in E. The first instability [see Fig. 1(a)] describes the direct conversion of two identical photons into two plasmons. Here we set $E_2 = 0$ and $\omega_1 = \omega$. The second instability [see Fig. 1(b)] is generated by the simultaneous interaction of two beams in the plasma. The photons at frequency ω_1 are down-converted to frequency $\omega_2 = \omega_1 - 2\omega$ with the emission of two plasmons. The beam at frequency ω_2 stimulates this process.

Inserting Eq. (9) into Eq. (7) leads to a pair of coupled integral equations for A_1 and A_2. In the long-wavelength limit ($\vec{k} \to 0$) an approximate solution can readily be found. Thus, we find an equation of the form[3]

$$|\epsilon(\vec{k}, \omega + i\mu)|^2 = \chi^2, \tag{10}$$

where we have introduced the dielectric function

$$\epsilon(\vec{k}, \omega) = \epsilon_L\left[1 - m^*\left(\frac{\omega_p}{k}\right)^2\int\frac{\vec{k}\cdot\partial f_0/\partial\vec{\eta}}{\alpha + \omega}d\vec{\eta}\right]. \tag{11}$$

The term χ, representing the nonlinear susceptibility, is given by

$$\chi = -\tfrac{1}{2}m^*\epsilon_L\left(\frac{\omega_p}{k}\right)^2\int d\vec{\eta}\, f_0\vec{k}\cdot\frac{\partial}{\partial\vec{\eta}}\frac{\gamma_1}{(\omega + \alpha - i\mu)(\omega - \alpha + i\mu)}, \tag{12a}$$

for the first instability, and for the second instability by

$$\chi = -\tfrac{1}{2}m^*\epsilon_L\left(\frac{\omega_p}{k}\right)^2\int d\vec{\eta}\, f_0\vec{k}\cdot\frac{\partial}{\partial\vec{\eta}}\frac{\delta}{(\omega + \alpha - i\mu)(\omega - \alpha + i\mu)}. \tag{12b}$$

We have taken $\hat{k} \parallel \vec{E}_1$ to guarantee a maximum growth rate. Aside from relativistic corrections, Eq. (11) reduces in the long-wavelength limit to the simple free-electron–like dielectric function $\epsilon_L[1 - (\omega_p/\omega)^2]$. Thus, we find for case (a), in

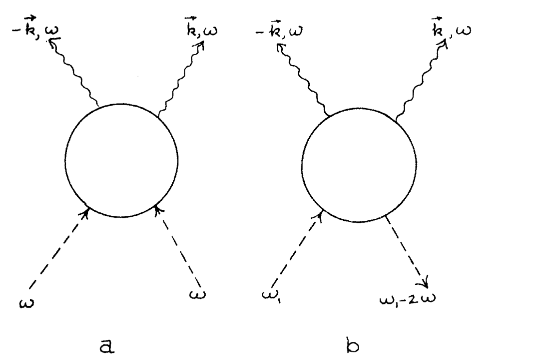

FIG. 1. (a) Conversion of two photons (dashed lines) into two plasmons (wavy lines). (b) Down-conversion of a photon stimulated by another photon.

the nonrelativistic limit,

$$\mu = \tfrac{3}{16}\omega(eE_1/m^*c^*\omega)^2, \tag{13a}$$

and for case (b)

$$\mu = \tfrac{3}{8}\omega(e/m^*c^*)^2(E_1E_2/\omega_1\omega_2). \tag{13b}$$

The effect of a finite collision time can be introduced simply by replacing μ by $\mu - 1/2\tau_0$, where τ_0 is related to the carrier mobility by $\tau_0 = m^*\mu_n/e$. Thus the net growth rate, Γ, can be written as

$$\Gamma = \mu - 1/2\tau_0. \tag{14}$$

The condition $\Gamma > 0$ is used to define the threshold field.

We now apply the theory to the case of InSb.[4] At $T = 77°$K, $E_g = 0.234$ eV, $m^* = m_e/60$, $\epsilon_L = 16$, and $c^* = 1.11\times10^8$ cm/sec. In determining the field in the crystal we use the relation $E = (8\pi I/c)^{1/2}(2/|1 + \sqrt{\epsilon}|)$, where $\epsilon = \epsilon_L[1 - \omega_p^2/\omega(\omega + i\tau^{-1})]$. Here I is the incident intensity. In Table I, part

TABLE I. Intensity for growth rate $\Gamma = 0.1$ period.

λ (μm)	Possible source	n (cm^{-3})	τ (sec) [a]	I_{thresh} (W/cm^2)	I (W/cm^2)
		(a)			
100	H_2O laser	3.0×10^{16}	7.6×10^{-13}	8.9×10^4	1.3×10^5
311	HCN laser	3.1×10^{15}	2.1×10^{-12}	1.1×10^4	1.5×10^4
1000	Microwaves	3.0×10^{14}	4.9×10^{-12}	1.7×10^3	2.2×10^3
5000	Microwaves	1.2×10^{13}	9.4×10^{-12}	3.3×10^2	3.7×10^2
		(b)			
10.81, 13.15	CO_2 laser	2.0×10^{16}	2.3×10^{-12}	8.4×10^6	$1.8\times10^7=\bar{I}$
10.81, 28	CO_2, H_2O	2.4×10^{17}	4.4×10^{-13}	6.3×10^6	$1.1\times10^7=\bar{I}$
28, 33	H_2O	2.2×10^{15}	2.3×10^{-12}	4.0×10^6	$5.5\times10^6=\bar{I}$

[a] C. Hilsum and A. C. Rose-Innes, *Semiconducting III-V Compounds* (Pergamon, New York, 1961).

(a), we calculate what intensity would be required for the growth rate to be a tenth of the plasma period. The various frequencies may be studied by employing tunable lasers, gas and vapor lasers, or even conventional microwave sources.

In the microwave case it should be pointed out that the nonrelativistic assumption, i.e., $v < c^*$, is violated. The growth rate quoted is therefore only an estimate. In the microwave case a self-sustained plasma can be produced in a very pure sample of InSb by impact ionization.

In the type (a) instability the plasma excitation occurs within roughly a wavelength of the surface. While this distance is fairly small, it is still much larger than typical plasmon wavelengths. Furthermore, it affords us the opportunity to probe the surface region selectively.[5] The type (b) instability is not restricted to the surface region, however, when ω_1 and ω_2 are above the plasma frequency. In Table I, part (b), we present several possible situations for observing the (b) instability. Again we calculate the intensity $\bar{I}$ for a growth rate which is one tenth of the inverse period. Here $\bar{I}$ is defined as $(I_1 I_2)^{1/2}$, I_1 and I_2 being the intensity in each beam.

In summary, our theory discusses two new nonlinear mechanisms by which photons are converted to plasmons. It is well known that for weak field strengths, the direct conversion of a photon to a plasmon is forbidden since a single transverse photon cannot excite a longitudinal plasmon. Only in the presence of surface or density gradients, for example, would the breaking of translational invariance permit such a process to proceed. The effect is therefore normally limited to several Fermi-Thomas lengths from the surface. Our mechanism, however, is operative through the bulk of the sample—the type (a) instability through a photon wavelength, and the type (b) instability through the whole sample—and is consequently dominant. Since it is a nonlinear process it becomes important when the field becomes sufficiently intense—as in the case of laser fields. We finally point out that the mechanism for the nonlinear coupling between photons and plasmons is an efficient way to produce large-amplitude plasma excitations. It appears to provide a suitable diagnostic method for studying nonlinearities in solid-state plasmas, especially in materials such as InSb.[6]

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Doorway-State Effects in the $M1$ Radiative Excitation of p-Wave Resonances in ^{57}Fe†

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The reaction ^{57}Fe(γ,n) has been studied near threshold. Intermediate structure observed in a strong p-wave resonance structure excited by $M1$ transitions is interpreted in terms of a doorway state resulting from a particle-hole excitation of ^{57}Fe.

Special reaction mechanisms such as doorway states can produce strong local concentrations of strength in the radiative cross sections for highly excited nuclear states. The well-known doorway state discovered by Farrell *et al.*[1] in studies of the reaction ^{206}Pb(n, n') is an excellent example. Baglan, Bowman, and Berman[2] observed a corresponding concentration of strength in studies of the reaction ^{207}Pb(γ, n) near threshold. Both the envelope and the fine structure of the strength of $E1$ transitions in the region of photoneutron energies of 200–600 keV correlate with the neutron data of Farrell *et al.* This note presents evidence for very sharp concentrations of strength in the cross section for $M1$ radiative excitation of p-wave levels in ^{57}Fe. These concentrations might be associated with a doorway consisting of an $(f_{5/2})(f_{7/2})^{-1}$ particle-hole pair coupled to the ^{57}Fe ground state. The data, obtained from high-resolution studies of the photoneutron cross section near threshold for ^{57}Fe, constitute a large sampling of ground-state radiation widths for p-wave resonances which heretofore could be obtained from neutron-induced reactions only with the greatest difficulty.

The measurements were performed on the Argonne threshold photoneutron facility[3] at the high-current electron linac. A 40-g target of ^{57}Fe was irradiated by a pulsed bremsstrahlung beam with the end-point energy adjusted[4] so that the nuclear states excited by photon absorption can decay by neutron emission only via a transition to the ground state of ^{56}Fe. Neutron resonance groups corresponding to each of the states excited were observed by time-of-flight measurements, with the detector set to observe neutrons emitted at 90° and 135° relative to the photon beam. Data taken at 90° are shown in Fig. 1. Observations covering the low-energy region [Fig. 1(a)] were taken with a ^{6}Li-glass neutron detector, while those covering the higher-energy region [Fig. 1(b)] were made with a proton-recoil detector. From the extensive data available on s-wave resonances in the total cross section of ^{56}Fe,[5] a complete list of expected s-wave ($j=\frac{1}{2}^+$) levels can be obtained. Where the resonance structure was completely resolved, i.e., below 225 keV, a measurable yield in the (γ, n) spectrum was obtained for all but one of the known s-wave levels. However, the most prominent feature of the results for ^{57}Fe as well as for other nuclei studied in this region is the strength of resonances with $l>0$. In Fig. 1(a) the integrated strength of the high-angular-momentum component is greater than that of the s-wave component.

To clarify this aspect of the data, the spins of strong neutron groups were determined by studying the angular distribution of the photoneutrons. Dipole absorption by the $\frac{1}{2}^-$ ground state of ^{57}Fe excites $\frac{1}{2}$ and $\frac{3}{2}$ states which then decay by neutron emission to the 0^+ ground state of ^{56}Fe. For the spin sequence $\frac{1}{2}\to\frac{1}{2}\to 0$, the photoneutron angular distribution will be isotropic; for $\frac{1}{2}\to\frac{3}{2}\to 0$, the ratio $d\sigma(90°)/d\sigma(135°)$ will be 1.43. Spin assignments were made by normalizing the data for 90° and 135° so that the relative yields gave isotropy for the strong s-wave level at 212 keV, and calculating the corresponding ratio for the other resonances. The intensity ratios observed were consistent with the assumption that only dipole photon absorption is important in the excitation process. The parities of resonances with $j=\frac{1}{2}$ were assigned by comparing the photoneutron results with the total neutron cross-section data on ^{56}Fe. The s-wave resonances ($j=\frac{1}{2}^+$) observed in this ^{56}Fe cross section were identified in the ^{57}Fe photoneutron spectra, and the remaining $j=\frac{1}{2}$ levels were assigned spin and parity $\frac{1}{2}^-$. Parity assignments for levels with $j=\frac{3}{2}$ could be complicated by the presence of d-wave photoneutrons.