Addison-Wesley, 1963].
[4] Ya. B. Zel'dovich and I. D. Novikov, Relyativistskaya astrofizika (Relativistic Astrophysics), Nauka, 1967.
[5] J. Gunn and J. Ostriker, Nature, 221, 454 (1969).

HYBRID SURFACE PLASMONS ON A METAL-SEMICONDUCTOR BOUNDARY

Yu. I. Balkarei and V. G. Baru
Institute of Radio Engineering and Electronics, USSR Academy of Sciences
Submitted 1 August 1969; resubmitted 21 August 1969
ZhETF Pis. Red. 10, No. 7, 324 - 327 (5 October 1969)

Interband transitions in metals have a strong influence on the character of the plasma oscillations [1]. In some cases, allowance for the dispersion of the electronic part of $\varepsilon(\omega)$, connected with the interband transitions, leads to the occurrence of "hybrid" volume plasmons, the frequency of which is much lower than that of the plasma frequency of the free electrons. Such a "hybrid" plasmon was observed experimentally in Ag [2]. A similar situation arises on the boundary between a metal and a semiconductor, where allowance for the dispersion of Re $\varepsilon(\omega)$ near the singularities of the interband density of states of the semiconductor leads to the appearance of surface waves of a new type.

Let us consider the contact between a metal and a semiconductor. Assume for simplicity that their work functions are equal or are such that the region where the bending of the bands plays a role is small compared with the wavelengths that are of importance.

Proceeding in the usual manner [3], without allowance for the retardation effects, we obtain the solutions of Maxwell's equations of the surface type in the form:

$$E_0^{(1)} \exp[i(k_1x + k_2y + i\kappa_1 z) \quad (z > 0 - \text{semiconductor})$$

$$E_0^{(2)} \exp[i(k_1x + k_2y - i\kappa_2 z) \quad (z < 0 - \text{metal})$$

Using the ordinary boundary conditions, neglecting spatial dispersion, and assuming the semiconductor and the metal to be isotropic, we obtain

$$\kappa_1 = \kappa_2 = \sqrt{k_1^2 + k_2^2}, \tag{1}$$

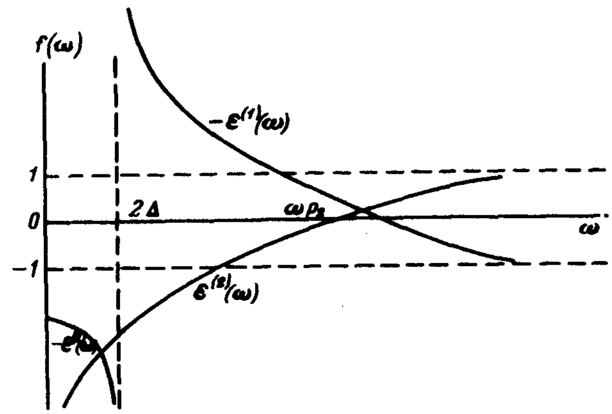
$$\varepsilon^{(1)}(\omega) + \varepsilon^{(2)}(\omega) = 0, \tag{2}$$

where $\varepsilon^{(1)}(\omega)$ and $\varepsilon^{(2)}(\omega)$ are respectively the dielectric constants of the semiconductor and of the metal. Equation (2) determines the possible frequencies of the surface oscillations. Taking for $\varepsilon^{(2)}(\omega)$ the simple plasma expression

$$\varepsilon^{(2)}(\omega) = 1 - \frac{\omega_{p2}^2}{\omega^2}, \tag{3}$$

where ω_{p2} is the plasma frequency of the electrons in the metal, and disregarding the frequency dispersion of $\varepsilon^{(1)}(\omega)$, we obtain from (2) the well known solution for the frequency of the surface plasmon, $\omega = \omega_{p2}[1 + \varepsilon^{(1)}(\infty)]^{-1/2}$. Actually, $\varepsilon^{(1)}(\omega)$ can have a strong frequency dependence near certain singular points of the interband density of states [2, 4]. In these cases, Eq. (2) should contain additional new solutions at frequencies $\tilde{\omega} \ll \omega_{p2}[1 + \varepsilon^{(1)}(\infty)]^{-1/2}$.

We shall demonstrate this using a very simple semiconductor model, which makes it possible to obtain $\varepsilon^{(1)}(\omega)$ in the entire frequency interval, without resorting to experimental data. We shall then discuss certain real cases, when the existence of new solutions should be expected.



We obtain an expression for $\varepsilon^{(1)}(\omega)$ of a semiconductor in the weak-coupling approximation using a model [5] in which the semiconductor is obtained from a metal having a Fermi surface of radius p_{F1} as a result of the appearance of a spherically symmetrical dielectric gap 2Δ on the Fermi surface. In the calculation of $\varepsilon^{(1)}(\omega)$, we take into account only transitions from the valence to the conduction band. The imaginary part $\mathrm{Im}\,\varepsilon^{(1)}(\omega)$ has in this model the form

$$\mathrm{Im}\,\varepsilon^{(1)}(\omega) = \begin{cases} 0 & (\omega < 2\Delta) \\ \dfrac{16}{3a_0 p_{F1}}\left(\dfrac{\epsilon_{F1}}{2\Delta}\right)\left(\dfrac{2\Delta}{\omega}\right)^4 \dfrac{\omega}{\sqrt{\omega^2 - 4\Delta^2}}\left\{1 - \left[\left(\dfrac{\omega}{4\epsilon_{F1}}\right)^2 - \left(\dfrac{\Delta}{2\epsilon_{F1}}\right)^2\right]^{1/2}\right\} & (\omega > 2\Delta). \end{cases} \tag{4}$$

Here a_0 is the Bohr radius and ε_{F1} is the Fermi energy. The dispersion relation yields $\mathrm{Re}\,\varepsilon^{(1)}(\omega)$ in the region $\omega < 2\Delta$. Neglecting corrections containing powers of the small parameter $2\Delta/\varepsilon_{F1}$, we obtain

$$\mathrm{Re}\,\epsilon^{(1)}(\omega) \simeq \frac{16}{3a_0 p_{F1}} \frac{2\Delta}{\sqrt{(2\Delta)^2 - \omega^2}} \operatorname{arctg} \frac{2\Delta}{\sqrt{(2\Delta)^2 - \omega^2}} \quad (\omega < 2\Delta). \tag{5}$$

The form of $\mathrm{Re}\,\varepsilon^{(1)}(\omega)$ when $\omega > 2\Delta$ can be found in [5]. The solution of (2) obtained by using (5) is shown schematically in the figure. The renormalization of the usual plasma root is of no importance to us now. Principal interest attaches to the root at the frequency $\omega < 2\Delta$. The position of the new solution, assuming it is to be close to 2Δ, is determined by the expression

$$\tilde{\omega} \simeq 2\Delta\sqrt{1 - \left(\frac{16}{3a_0 p_{F1}}\right)^2 \left(\frac{\epsilon_{F1}}{\omega_{p2}}\right)^2 \left(\frac{2\Delta}{\omega_{p2}}\right)^2}. \tag{6}$$

At typical parameter values $a_0 p_{F1} \sim 1$, $\varepsilon_{F1} \sim \omega_{p2}$, and $2\Delta/\omega_{p2} \sim 10^{-1}$, the value of $\tilde{\omega}$ is actually near 2Δ. The attenuation of the new mode is determined by the imaginary parts of $\varepsilon^{(1)}(\omega)$ and $\varepsilon^{(2)}(\omega)$ at $\omega \sim \tilde{\omega}$. In this frequency region $\mathrm{Im}\,\varepsilon^{(1)}(\omega) = 0$, and $\mathrm{Im}\,\varepsilon^{(2)}(\omega)$ has in the random-phase approximation the following form when $T = 0$ [1]:

$$\mathrm{Im}\,\epsilon^{(2)}(\omega) = \frac{\pi}{2}\frac{\tilde{\omega}}{kv_{F2}}\frac{\kappa_{D2}}{k^2}\theta(kv_{F2} - |\tilde{\omega}|), \quad (7)$$

where κ_D is the reciprocal Debye radius, v_F is the velocity on the Fermi surface, and

$$\theta(x) = \begin{cases} 1 & x > 0 \\ 0 & x < 0 \end{cases}$$

Thus, just as for ordinary plasmons, there is no attenuation of the new surface plasmons in the random-phase approximation when $k < \tilde{\omega}/v_{F2}$. The attenuation appears in this region of k when account is taken of the electron scattering in the metal plasma.

As seen from the foregoing analysis, the appearance of a new solution at frequencies $\omega \ll \omega_{p2}$, where Re $\epsilon^{(2)}(\omega)$ is negative and large, is possible only in the region where Re $\epsilon^{(1)}(\omega)$ is large and positive, i.e., in the region of strong frequency dispersion. The concrete character of the dispersion plays no role.

In our model, strong dispersion appears near the threshold of the interband absorption. Actually, Re $\epsilon^{(1)}(\omega)$ experiences usually a strong growth not at the threshold, but near the first crest of the interband density of states. In this region, further, Im $\epsilon^{(1)}(\omega) \ll$ Re $\epsilon^{(1)}(\omega)$. This character of $\epsilon^{(1)}(\omega)$ was obtained experimentally for Ge, Si, InSb, InAs, GaAs, GaP, Ag, and Cu [2, 4]. A particularly strong dispersion is observed, for example, in Si, where Re $\epsilon^{(1)}(\omega)$ increases by approximately 4 times in the interval 3.2 - 3.4 eV. In Ag, the strong dispersion of Re $\epsilon^{(1)}(\omega)$, connected with the start of interband transitions at $\hbar\omega \sim 3.9$ eV, leads to the appearance of the experimentally observed "hybrid" volume plasmon [2].

In the case when the strong dispersion of Re $\epsilon^{(1)}(\omega)$ is connected with the exciton levels, new surface solutions should arise ahead of the intrinsic absorption edge.

The plasmons considered here can make an appreciable contribution to the onset of superconductivity in layered metal-semiconductor structures[1)].

They can appear experimentally in the spectrum of the characteristic losses, in electron tunneling, and may also be excited by an electromagnetic field under definite conditions [3, 7].

We are grateful to V. B. Sandomirskii and L. N. Bulaevskii for a discussion and useful remarks.

[1] D. Pines, Elementary Excitations in Solids, Benjamin, 1963, (Russ. transl., Mir. 1965, p. 272).
[2] H. Ehrenreich and H. R. Philipp, Phys. Rev. 128, 1622 (1962).
[3] V. M. Agranovich, Teoriya eksitonov (Exciton Theory), Nauka, 1968, p. 246.
[4] H. R. Philipp and H. Ehrenreich, Phys. Rev. 129, 1550 (1963).
[5] A. Bardasis and D. Hone, Phys. Rev. 153, 849 (1967).
[6] E. A. Pashitskii, Zh. Eksp. Teor. Fiz. 56, 662 (1969) [Sov. Phys.-JETP 29, 362 (1969)].
[7] D. Beaglehole, Phys. Rev. Lett. 22, 708 (1969).

[1)] The possible occurrence of high-temperature superconductivity in a heterojunction, as a result of surface plasmons, is discussed, for example, in [6].